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UNIVERSITY OF CALIFORNIA, SAN DIEGO

Theoretical and Observational Analysis of
Particle Acceleration Mechanisms at Astrophysical Shocks

A dissertation submitted in partial satisfaction of the
requirements for the degree Doctor of Philosophy

in

Physics

by

Edward Lawrence Lever

Committee in charge:

Kim Griest, Chair
Kevin B. Quest, Co-Chair
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2011

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University of California, San Diego

2011

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“Shock Surfing vs. Shock Drift Acceleration”, Lever, E.L., Quest, K.B., Shapiro, V.D.,
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Studies in Computational Astrophysics
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Studies in Astrophysics and Compact Objects
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Studies in Non-Neutral Plasmas
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ABSTRACT OF THE DISSERTATION

Theoretical and Observational Analysis of
Particle Acceleration Mechanisms at Astrophysical Shocks

by

Edward Lawrence Lever

Doctor of Philosophy in Physics

University of California, San Diego, 2011

Professor Kim Griest, Chair

Professor Kevin B. Quest, Co-Chair

We analytically and numerically investigate the viability of Shock Surfing as a pre-injection mechanism for Diffusive Shock Acceleration, believed to be responsible for the production of Cosmic Rays. We demonstrate mathematically and from computer simulations that four critical conditions must be satisfied for Shock Surfing to function; the shock ramp must be narrow, the shock front must be smooth, the magnetic field angle must be very nearly perpendicular and, finally, these conditions must persist without interruption over substantial time periods and spatial scales. We quantify these necessary conditions, exhibit predictive functions for velocity maxima and accelerated ion fluxes based on observable shock parameters, and show unequivocally from current observational evidence that all of these necessary conditions are violated at shocks within the heliosphere, at the heliospheric Termination Shock, and also at Supernovae.

Chapter 1 Introduction

We numerically investigate Shock Surfing, suggested as a pre-injection mechanism for the production in the heliosphere of Anomalous Cosmic Rays (ACR) by diffusive shock acceleration (DSA). We demonstrate analytically and computationally that Shock Surfing (SS) is not a viable mechanism, and show that these conclusions are consistent with the latest observational data.

While SS has been shown to function well under very stringent conditions at quasi-perpendicular shocks, we demonstrate several factors each of which can severely limit not only the maximum velocity attainable by ions under this mechanism but also the flux of ions promoted above the velocity ceiling of Shock Drift Acceleration (SD). One conundrum is that DSA requires magnetic turbulence near the shock, which may disturb it. Whether perturbations along the shock are induced by this turbulence or result from the manifold waves generated by other processes, we show that such corrugations suppress the surfing mechanism, a coherent process that requires a smooth front.

We also demonstrate that the maximum attainable surfing velocity is inversely proportional to ramp width, and the minimum potential gradient required for surfing velocities to exceed those of the SD in turn places limits on this width. Further, since DSA is an inherently slow process, time and length scales place constraints on this mechanism, particularly at convex planetary bow shocks where the probability of escape along the flanks is high due to rapid advection by the solar wind. This constraint is also detrimental to SS, and limits its viability for both DSA pre-injection and as a standalone mechanism. This is especially true at the terrestrial bow shock, yet observational data

have confirmed that gas giant shocks, although much larger, have ramps too broad to support surfing.

Although the heliospheric Termination Shock (TS) was long believed to be a viable possibility, recent Voyager observations have confirmed that the ramp there is also too broad to support surfing, and that the flux of ions accelerated by other means is outward-bound. Based on these observations, recent consensus is that ACR are not created in the quasi-perpendicular ecliptic plane, if at all, at the heliopause.

Finally, we show that, perpendicular shock reformation time scales observed both *in situ* and in simulations, limit the maximum attainable shock surfing velocity to that of SD, rendering the mechanism's contribution to pre-injection insignificant at best. If in fact this periodic reformation (on the order of an ion cyclotron period) is a common feature of perpendicular shocks, we claim that shock surfing cannot be a viable acceleration mechanism anywhere at all.

Our work is motivated by the preponderance of research in this field over the past several years, despite publication of our early results nearly a decade ago (Lever et al, 2001). There, we demonstrated that a maximum shock ramp width existed beyond which SS was completely ineffective and, incidentally, that a particle's shock-incident location in velocity space completely determined which acceleration mechanism acted upon it. This was followed by the presentation of our research on surfing at turbulent shocks with the important result that both surfing velocity maxima and flux above that of the shock drift mechanism were severely attenuated, quite possibly beyond the point of viability.

Nevertheless, shock surfing research has continued, e.g., Ucer & Shapiro (2001), Shapiro et al (2001), and Ucer (2005), with variations including "Gyroresonant" surfing, Kuramitsu & Krasnoselskikh (2005) and "Stochastic" surfing acceleration, Chandran &

Bessho (2004). Remarkably, interest in SS at the Termination Shock has continued even after *in situ* analysis from Voyager conclusively established that it was not the source of ACR – Burrows et al (2009), Kichigin (2009), and Yang Zhong-Wei et al (2008). One such example may be found in Chalov (2005) who asserts, “It may not be doubted anymore that [ACRs] are produced in the heliosphere from interplanetary pick-up ions through their acceleration at the solar wind termination shock” while adding that “... there is no general agreement in the community of heliospheric researchers concerning the mechanism of injection of the pick-up ions into the shock acceleration.”

As a passing caveat to the reader, sustainable shock surfing by ions should not be confused with a recently proposed mechanism called **electron** shock surfing. Articles on this subject have appeared which include Amano & Hoshino (2009), Amano & Shinohara (2009), Umeda et al (2009), Umeda et al (2008), Ohira & Takahada (2007), Sircombe et al (2006), Dieckmann et al (2005), Dieckmann et al (2004), Hoshino & Shimada (2002), and Hoshino (2001). The process they promote may be briefly described as follows: waves grow due to the Buneman instability, and electrons are accelerated on these via the ‘surfatron’ mechanism. Matsukiyo & Scholer (2006) present a review of this mechanism, along with discussion of the associated Modified Two Stream Instability (MTSI).

Finally, we question whether DSA is itself viable and, if not, what other mechanisms may be responsible for CR production. Perhaps the most interesting so far suggested are Dieckmann et al. (2006) who propose acceleration at jets, and Lazzati et al (2002) supernova “Fireball Surfing”. The original model of colliding shocks first proposed by Fermi has recently been revisited. Roth & Bale (2006) study merging shocks of Coronal Mass Ejections (CME), and Dudkin et al (2000) investigate Deuterium

acceleration at colliding shocks. Magnetic reconnection has been explored by Hoshino (2002), by Scholer (2004) as an alternate injection mechanism at the termination shock, as a beam source by Mukai (2000) and as an ACR source by Drake (2010).

Some work has appeared which explored previously overlooked difficulties. Scholer et al (2003) asserted that surfing should not be viable at the TS due to instabilities in the electrostatic potential, Dieckmann & Shukla (2006) illustrated instabilities at relativistic colliding plasmas, Gonzalez-Esparza et al (2003) argued that surfing is not viable at CME due to their relative speed to the solar wind, and Anagnostopoulos et al (2000) argued in favor of Shock Drift over Shock Surfing at the Jovian bow shock.

Fortuitously, we now have the ability to validate our own predictions with observational data gathered by the Cluster Mission at the terrestrial bow shock, the Cassini Mission, and by the two Voyager Missions, which explored the outer planets and the Termination Shock. The surprises delivered by the second Voyager mission reveal the difficulty of this subject. Indeed, many theoretical treatments have been revised with respect to the structure of the TS, while some previously published analyses of data suggesting the viability of shock surfing there have reappeared in the literature with the phrases "... a re-evaluation" and "... a re-interpretation" appended to their titles. According to Lembege et al (2004), the physics of collisionless shocks is a vastly broad topic, well studied for many decades with a number of important but unresolved issues remaining. Indeed, many new findings have "cast doubt on well-established ideas". On the topic of DSA itself, Zank (2004) asserts that the "problem of perpendicular diffusion by a particle in a turbulent plasma is a problem of enduring interest ... yet to be fully solved. Analytic models do not agree with either observations or numerical simulation."

Among the surprises from the Voyager mission were contradictions to some of the most fundamental assumptions made regarding the solar wind. Theoretical estimates base upon the Parker Solar Wind Model failed to correctly predict features as simple as the location of the shock itself. The model requires closure by an equation of state but, as it happens, adiabatic and isothermal models both appear from observations to fail at the TS. Also, data show that the outgoing (downstream) fluid remains supersonic, begging the question as to whether the TS is, in fact, a true shock at all.

Our aim is to put further investigation of shock surfing to rest. We present conditions each of which individually render the surfing mechanism untenable, and hope to convince the reader that, taken together, the probability is vanishingly small that at least one of these is not present at any given time.

In this chapter, we present a brief overview of problem and, in Chapter 2, develop a theoretical framework from which we make estimates and predictions of the mechanism behavior as functions of parameters such as ion mass, shock Mach number, etc. In Chapter 3, we recap the results of our 2001 GRL publication, establishing baseline data over a range of parameters for test particles at a smooth, quiescent shock front. In Chapter 4, we present the results of our investigation into the effect on SS of various perturbations including waves and physical corrugation of the front, portions of which were previously given in late 2001 at APS and AGU conference meetings.

In Chapter 5, we seek a tractable function to express the DSA injection threshold velocity, investigate the foundations of turbulent velocity diffusion and connect the turbulence to front ripples. In Chapter 6, we exhibit data from Cluster Mission observations and show that SS is not a viable at the terrestrial bow shock. We continue our investigation with other shocks in the solar system, examining case studies from

planetary missions, and conclude a 'Grand Tour' of the heliosphere in Chapter 7 with the Termination Shock. Finally in Chapter 8, we seek to discover whether shock surfing may be possible at supernovae.

History and Motivation

The source of cosmic rays has been an intriguing problem that has captured the imagination of great minds for over half a century. Interest has waxed and waned decade by decade as new legions of theorists and observers turn their attention towards solving the many questions that still remain. Particle acceleration is a common feature of astrophysical shocks and Cosmic Rays are believed to originate at those of supernovae and other violent processes associated with compact objects. Diffusive Shock Acceleration (DSA) is a theoretical model developed to describe the mechanism by which shocks promote particles to extremely high cosmic ray energies.

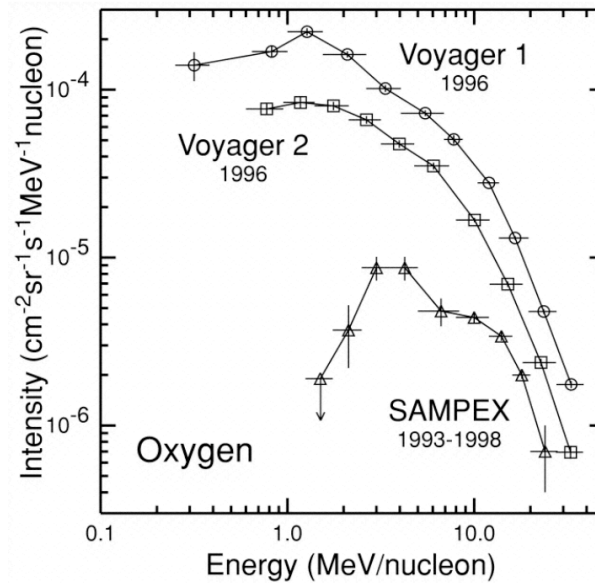


Figure 1-1 Oxygen Component of the Anomalous Cosmic Ray Spectrum. After Hill et al (2002)

But an “anomalous component” in the energy spectrum shown in Figure 1-1 was believed to originate within the solar system itself, perhaps at planetary bow-shocks or the heliospheric termination shock. Quasi-perpendicular shocks have been of particular interest since they are believed to accelerate ions more quickly & efficiently than quasi-parallel ones (Ellison et al. 1995). The termination shock was believed to be quasi-perpendicular in the ecliptic plane, due to the winding up of the magnetic field at that distance, and so became a subject of intense scrutiny for theorists and observers alike.

The problem is that, according to many theorists, DSA requires a ‘seed’ population of particles of sufficient energy to begin the process, and the thermal component of the solar wind lies below this threshold. It was thought that the Shock Drift (SD) Mechanism might prove a suitable candidate (Hudson 1965, Winske & Leroy 1983, Armstrong et al 1985, Forman & Webb 1985, Decker 1988 and Giacalone et al 1991).

However, this mechanism has a sharp velocity ceiling, which limits its application in this case.

Shock Surfing was first proposed by Sagdeev (1966) to explain how charged particles might be accelerated to high velocities at these perpendicular shocks where there is no easy path to allow multiple transits across the front, a necessary condition for DSA to promote them to cosmic ray energies. Theoretical models were further explored and refined by Sagdeev and Shapiro (1973), Lee et al (1996), Zank (1996), Lee et al (1997), Lipatov et al (1998) and Lipatov & Zank (1999). The process was believed to be a viable pre-injection mechanism for DSA, both theoretically and through numerical simulations, but questions remain: does this process actually take place and, if so, does it significantly contribute to observed cosmic ray fluxes?

It is impossible to prove the effectiveness or even presence of Shock Surfing as a pre-injection mechanism through forensic analysis of observed cosmic ray fluxes, since DSA is a stochastic process, which erases all evidence of the injection mechanisms that feed it. It may, however, be possible to show under what conditions Shock Surfing does **not** work and thereby more critically evaluate its possible contribution and effectiveness.

This is the primary purpose of our investigation.

Diffusive Shock Acceleration

Enrico Fermi first proposed in his 1949 paper, *On the Origin of the Cosmic Radiation*, “A theory of the origin of cosmic radiation is proposed according to which cosmic rays are originated and accelerated primarily in the interstellar space of the galaxy by collisions against moving magnetic fields... The chief difficulty is that it fails to explain in a straight-forward way the heavy nuclei observed in the primary radiation.”

This paper launched a body of research, which has continued for over a half century and has covered not only cosmic ray production, but also the more fundamental nature of astrophysical plasmas including collisionless shocks. Note Fermi’s caveat regarding the difficulties and contradictions evident even at this early stage – a prophetic harbinger of controversies that followed.

The study of these processes requires profound understanding across a broad scope of disciplines; for example, the study of DSA is inextricably intertwined with that of collisionless shocks, an outgrowth of the work of J. M. Burgers in the 1930’s. The tutorial review “A quarter century of collisionless shock research” (Kennel, 1985) is aptly titled. A veritable industry developed during that period with analytic and numerical works by Hudson (1965), Kadomsev & Karpman (1971), Krymsy (1977), Axford, Leer & Skadron (1977), Bell (1978), Blandford & Ostriker (1978), Lagage & Cesarsky (1983), Blandford & Eichler (1987), Forman & Webb (1985), and numerous other works by Armstrong, Decker, Drury, Ellison, Galeev, Lee, Ostrowski, Quest, Sagdeev, Scudder, Volk and Winske (1983 – 1988). Based upon this work, DSA theory advanced substantially with publication of a seminal paper by Webb et al. (1996) and additional references therein.

The Injection Problem

One feature of the Webb model is a minimum injection threshold required for ions to begin the acceleration process. According to Ellison et al (1995) "... only modest gains in acceleration rates can be expected for strong oblique shocks over parallel ones if the only source of seed particles is the thermal background." It is essential, then, for some pre-injection mechanism to accelerate particles up to this 'threshold' velocity. The question, which has remained controversial, is what exactly is that threshold? Estimates have ranged from many times the upstream fluid velocity, to a few times (Giacalone 2005), to none at all (Dworsky & Fahr 2000). Zank et al (2005) summarized the problem by asserting, "... perpendicular diffusion by a particle in a turbulent plasma is a problem of enduring interest and one that has yet to be fully solved. Analytic models do not agree with either observations or numerical simulation."

Clearly, if Shock Surfing can be shown to be viable, establishing this threshold is of paramount importance. We begin our investigation by defining the 'arena' in which the game is to be played.

The Parker Solar Wind Model

In 1958 Eugene Parker published his hypothesis of a supersonic solar wind in The Astrophysical Journal which, although initially met with considerable skepticism, was directly observed the following year by the Soviet satellite Luna One. The Solar Wind (SW) is a plasma, predominantly ionized hydrogen (H^+ and e^-) which flows radially outwards from the sun's corona through the solar system until it encounters the Interstellar Medium (ISM). In so doing, it carries the solar magnetic field with it, creating an immense 'bubble' of magnetized plasma called the heliosphere.

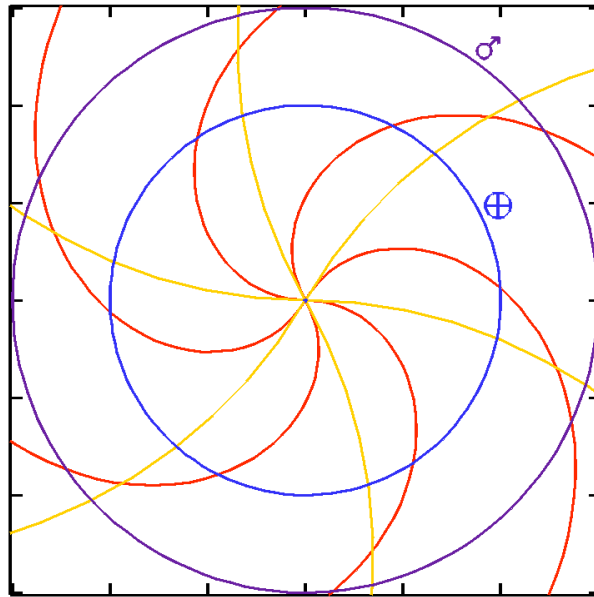


Figure 1-2 Magnetic field structure of the solar wind in the ecliptic
The red curves represent the 'slow' wind and yellow, the 'fast'. Orbits of the Earth and Mars are shown to illustrate the difference in Parker angle.

Due to the sun's rotation, the dipole field becomes twisted and stretched out in the ecliptic plane, as shown in Figure 1-2, creating a vast current sheet or disk in the ecliptic. The interplanetary magnetic field (IMF) within the heliosphere takes on a spiral shape in the ecliptic and the Parker Angle that field lines make with the radial becomes more perpendicular with solar distance. Along the Earth's orbit, this angle is approximately 45° , and approaches 90° at the edge of the heliosphere.

Shocks in the Heliosphere

Since the Solar Wind is supersonic with respect to the Alfvén speed, shocks form when it encounters obstacles. The plasma is tenuous and, with extremely long mean

free paths even for Coulomb scattering, such shocks are collisionless. Bow shocks form at the magnetospheres of magnetized planets, as well as at non-magnetic planets where the SW induces magnetospheres. Shocks also form at the noses of comets, Coronal Mass Ejections (CME) and co-rotating interactive regions (CIR).

Although most planetary shocks are quasi-stationary, they are responsive to changes in the solar wind pressure, which cause secular displacement of the entire shock structure in time. They are also pummeled by turbulence created during periods of violent solar activity such as flares and CME. Shocks at the gas giants also exhibit variability – even during periods of solar quiescence – due to strong coupling of the magnetospheres with rapidly rotating planetary atmospheres as well as effects from geologically active satellites.

Shock Features

All shocks are characterized by a transition from supersonic to subsonic flow and an increase in density to conserve momentum and temperature to conserve energy. Since the solar wind is magnetized, there is also a corresponding increase in magnetic field strength. Perpendicular shocks also feature an electrostatic potential step due in part to the electron-ion mass difference. Simulation examples may be seen in the plots of Winske and Leroy (1988) Figure 1-3, and a plot of *in situ* data at the Bow Shock is shown from Bale et al. (2003) Figure 1-4. The compression ratio is a function of the Alfvén Mach number and plasma parameters (beta and gamma), but fine structure can vary substantially, depending on the magnetic field angle with respect to the front.

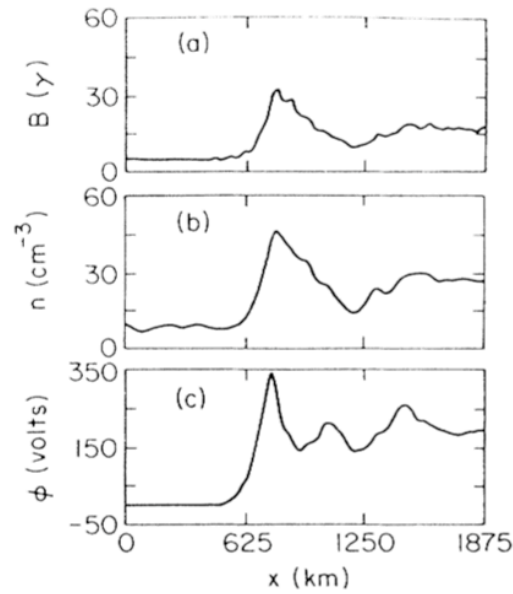


Figure 1-3 Density and Field profiles at a quasi-perpendicular shock
Fluid flows rightwards through ramp of width 200 km.
After Winske and Leroy (1988)

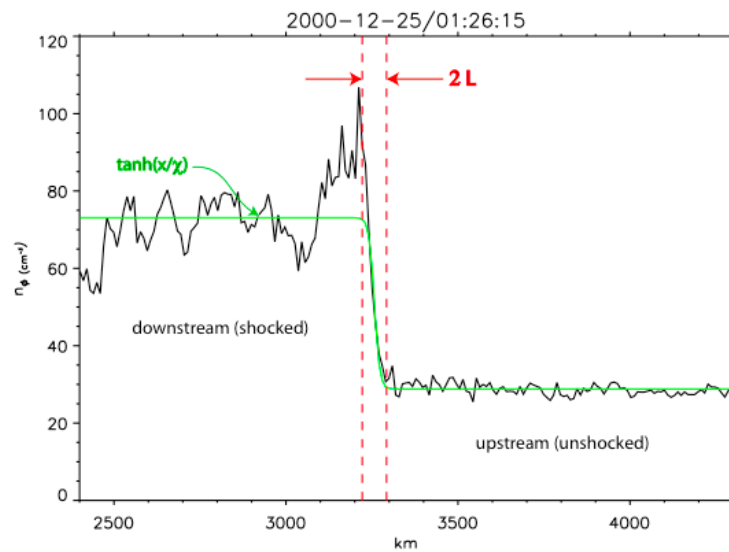


Figure 1-4 Density profile at the Earth's Bow Shock
Ramp width is roughly 100 km, an ion inertial length.
Cluster mission observations, after Bale et al (2003)

Parallel Shocks

Quasi-parallel shock regions are generally characterized by substantial turbulence, broad ramps and poorly defined transitions. They do not support shock drift since the magnetic gradient is parallel to the ambient field. Due to field lines penetrating the front and the absence of an electrostatic potential gradient, shock surfing is not supported there either. They theoretically have long DSA acceleration times, but can readily inject thermal ions and are thus expected to be sources of energetic particles. However, they may also be effective in accelerating ions by stochastic scattering. *In situ* observations have revealed the presence of SLAMS (Small Large Amplitude Magnetic Structures), extreme variability in coherence length, and other features which suggest that, under favorable conditions, it may be possible for surfing ions from a perpendicular region to reach a parallel one and so be further accelerated there.

Perpendicular Shocks

At a perpendicular shock, the magnetic field is perpendicular to the shock normal, i.e., it lies in the shock tangent plane. As such, there is a convective (motional) electric field in the shock frame as the fluid carrying the magnetic field passes through. There is also an electric field normal to the front, which results from charge separation due to the mass difference between electrons and ions as they slow during transit. A jump in electrostatic potential follows from this field. Additional features include a foreshock region or foot, a ramp overshoot, and decaying downstream oscillations in magnetic field strength. Perpendicular shocks are intrinsically more stable than parallel ones, and DSA theory predicts that these shocks accelerate ions more rapidly, but can't inject thermal

ions. Hence it is here that an efficient pre-injection mechanism is not only most needed, but would also be most effective.

Pickup Ions

Since solar wind thermal ions have been ruled out as pre-injection candidates, we now consider a more suitable population. Pickup ions are created from background neutrals by photoionization or collisions. The sources are not from the solar wind itself, but the interplanetary medium, planetary atmospheres or influx from the interstellar medium. Since these newborn ions are essentially at rest in the heliosphere, they are pickup up and carried along by the fields in the solar wind, evolving in time first into a ring distribution in velocity space by their cyclotron orbits and ultimately by pitch-angle scattering into a shell as shown in Figure 1-5.

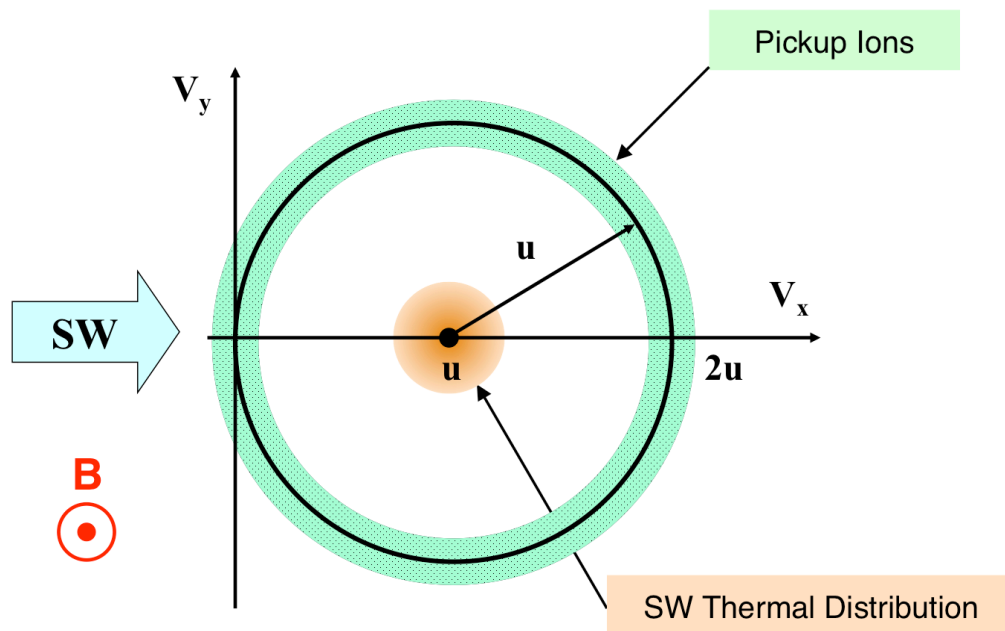


Figure 1-5 Pickup ion distribution in velocity space

The cross-section shown lies in the shock stationary reference frame

Time and Length Scales

Spatio-temporal scales play an important role in determining the viability of pickup ions as candidates for acceleration. Saul et al. (2004) have estimated the time scales for various processes affecting pickup ions in the heliosphere, to which we have added distance scales based roughly on a mean solar wind speed of 500 km/s as shown below in Table 1-1.

Table 1-1. Pickup Ion time and length scale evolution.

R_E is the Earth's radius, $R_H \approx 15 R_E$ is the radius of the Magnetosphere, i.e., the 'standoff' distance to the nose, and AU is an Astronomical Unit – the Earth's mean solar orbital radius.

Lorenz Force	Forms gyrotropic distribution	10 – 30 s	$\sim 1 R_E$
Pitch Angle Diffusion	Forms isotropic distribution	10 – 100 hr	$\sim 10^3 R_H$
Adiabatic Cooling	Expanding SW and decreasing IMF cools Pickup Ions	100 hr	$\sim 1 \text{ AU}$
Energy Diffusion	Statistical acceleration by waves Scales with Alfvén Mach number	1000 hr	$\sim 10 \text{ AU}$
Coulomb Scattering	SW density decreases as r^{-2} $mfp \rightarrow$ order of the heliosphere	50 days	$\sim 15 \text{ AU}$

Note that Pickup Ion 'age' has consequences on its distribution when encountering a shock. If 'born' anywhere near the magnetosphere, for example, they quickly assume a ring distribution, while development of a fully isotropic shell requires three orders of magnitude larger distance than the standoff itself. Since they are generally accepted as being created by photoionization near the sun, this provides ample time for the shell to form at a distance of 1AU.

Shock Drift Acceleration

We now introduce two complementary mechanisms, each of which is particularly well suited to accelerate pickup ions. At a purely perpendicular shock, the upstream magnetized solar wind plasma flows in the shock normal direction $\mathbf{u}_1 = u_1 \hat{\mathbf{x}}$ while the magnetic field is parallel to the shock surface $\mathbf{B}_1 = B_1 \hat{\mathbf{z}}$ so that, in the shock stationary frame, a convective electric field $\mathbf{E}_{\parallel} = -\mathbf{u} / c \times \mathbf{B}$ is induced in the $\mathbf{y}$ -direction along the shock surface. Both the Shock Drift (SD) and Shock Surfing (SS) mechanisms produce energy gain by displacing the ion guiding centers along the direction of this convective electric field $\Delta\epsilon = qE_{\parallel}\Delta y$.

SD displacement is achieved through the magnetic gradient drift mechanism, proportional to $\mathbf{B} \times \nabla \mathbf{B}$. Since the field increases through the ramp $\nabla \mathbf{B} = \nabla B \hat{\mathbf{x}}$, the drift velocity is along the front $\mathbf{v}_{\nabla B} = v_{\nabla B} \hat{\mathbf{y}}$, and ions are displaced in the direction of the convective electric field a distance $\Delta y = v_{\nabla B} \Delta \tau$ as they pass through the ramp gradient in time $\Delta \tau$. Early models – see Giacalone (1991) and references therein – were based on the assumption that the ion cyclotron radius was small compared to the magnetic field gradient scale (shock ramp width), and calculations based on this assumption were thus consistent with multiple shock crossings as shown by the idealized ion trajectory in Figure 1-6. However, as discussed in Kallenrode (2004) and validated in our own as well as other simulations, SD ions return to the front only once, promoting Pickup Ions by at most several times the piston energy, so the final downstream velocities for these particles range from between two to three times the upstream SW velocity. This is the

main drawback for SD since, should this hard velocity ceiling lie below the minimum DSA threshold, the mechanism is rendered completely inconsequential for injection.

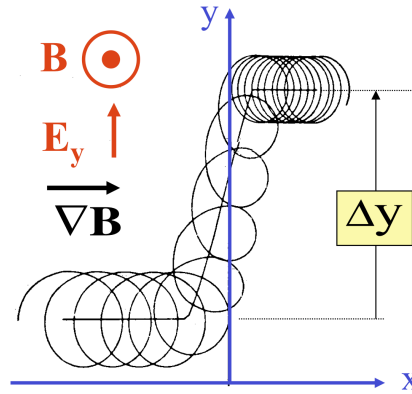


Figure 1-6 Idealized Shock Drift trajectory

Ions shown in this early model must satisfy $v_{\perp} \gg u_{sw}$.

Shock Surfing

Thermal solar wind ions are slowed by the potential gradient at the front, and proceed downstream. Some pickup ions, however, may be reflected if they encounter the front with insufficient normal velocity to overcome the potential barrier, as shown below in Figure 1-7. If they still have insufficient normal velocity on subsequent encounters, they will continue their reflections – much like a stone skipping on a pond – until they finally breach the potential barrier and escape downstream. As long as they remain constrained upstream, these ions will be accelerated by the convective electric field in the shock-stationary frame, theoretically to much higher velocities than those attainable through Shock Drift.

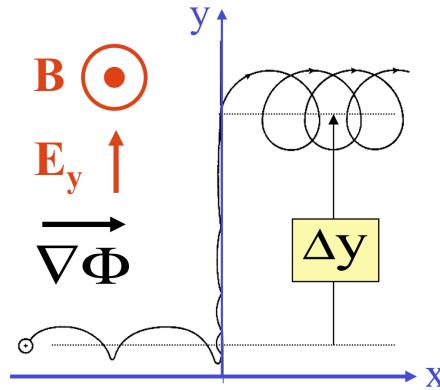


Figure 1-7 Idealized Shock Surfing Trajectory

An electrostatic potential step is shown. After Lee et al (1996)

Idealized Goals

In order to pursue a comprehensive investigation of Shock Surfing and its possible contribution to various observed phenomena, we must be able to answer the following questions:

- Is there a maximum velocity to which SS can accelerate ions?
- What dependence if any does SS have on various parameters such as ion mass, shock Mach number and magnetic field angle?
- Can ions accelerated by this mechanism be directly detected as such?
- Is it possible to either measure or infer SS as the DSA precursor mechanism by analysis of the cosmic ray spectrum – as can be done for astrophysical processes by virtue of detected photon polarization?
- Can SS explain the fluxes in the anomalous CR spectrum, specifically the heavier ion components such as oxygen?
- How robust is the SS mechanism in the presence of turbulence, shock structure anomalies and the possible spatial and temporal constraints these may impose?

The most ideal scheme in pursuing a genuine scientific proof would be to apply, in the true spirit of the Scientific Method, the following sequence of steps: Observation, Analysis, Hypothesis, Prediction and, finally, Validation. In our area of investigation, these translate into the following specific tasks:

1. Observe and measure the actual flux of each PI species incident on the shock of interest
2. Derive a shock surfing ‘throughput’ function to produce a resultant distribution function for each of the accelerated species of interest
3. Based on observed shock parameters and those of the regional plasma, calculate a specific set of DSA injection threshold velocities (possibly one for each species, if necessary)
4. Applying currently accepted ‘state-of-the-art’ theory, calculate the DSA throughput function and thereby predict observable fluxes, i.e., complete distribution functions
5. Validate the theory and process by comparing the predictions with the observed spectra

If we could accomplish all that and match prediction to data, then we will have forged an unbreakable chain of reasoning – a proof.

Chapter 2 Theory & Predictions

In the previous chapter, we introduced two mechanisms – Shock Surfing (SS) and Shock Drift Acceleration (SD) – as pre-injection candidates for Diffusive Shock Acceleration (DSA). In this chapter we lay the theoretical foundations for these mechanisms and shall investigate them numerically thereafter. We do this in order to establish a set of baseline parameters that will guide the direction of our simulations, and also provide a means of validating our numerical code in various limiting cases.

As seen from the observational data above, real shocks are not discontinuities, but exhibit finite features that necessarily complicate the mathematical models used to describe the mechanisms of interest to us. The solutions are not analytically tractable without making simplifying assumptions and must therefore be explored numerically. In order to accomplish this to best advantage we shall adopt a simplified scheme for the field profiles, a schematic representation of which may be seen in Figure 2-1 below.

While the surfing mechanism is primarily dependent upon the shock-normal electric field, some acceleration of those ions must necessarily occur by drift action as they pass through the magnetic field gradient. Similarly, the orbits of shock drift ions are also affected by the retarding electric field. In other words, the two mechanisms are not entirely disjoint.

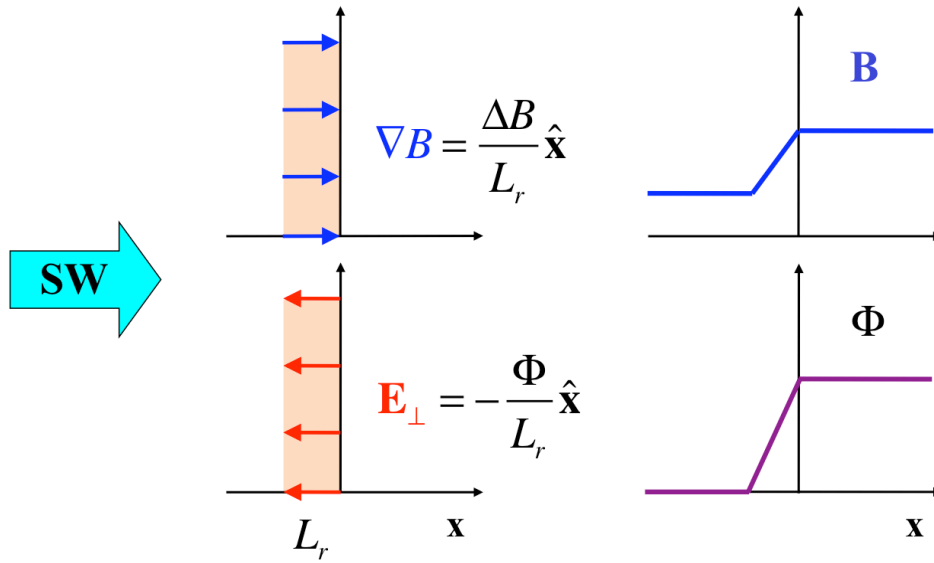


Figure 2-1 Finite Ramp Field Profiles

The schematic shows the magnetic field and electrostatic potential on the right and associated field gradients in the shock ramp on the left.

Theory & Predictions

A mathematical treatment introduced by Lee et al (1996), models an idealized shock front with Heaviside step function of zero width for the ramp. All reflections are elastic, and the criterion for transmission of ions downstream is that their normal kinetic

energy exceeds that of the electrostatic potential barrier $\frac{1}{2}mv_{x0}^2 > q\Phi$. In terms of the

original PI shell distribution, this imposes the restriction that reflected ions must have

incident normal velocity such that $v_{x0} < \sqrt{\frac{2q\Phi}{m}}$. It is immediately apparent that, for singly

ionized species, the fraction of PI reflected on first encounter decreases as the square

root of the ion mass. For sufficiently thin PI shells, this can be approximated by a surface

integral $\frac{1}{4\pi} \int_0^{\theta_0} 2\pi \sin \theta d\theta = \frac{1}{2}(1 - \cos \theta_0)$ where $\cos \theta_0 = 1 - \frac{v_{x0}}{u_1}$, giving a coefficient that

represents the fraction of reflected PI $\mathfrak{R}_{ion} = \frac{v_{x0}}{2u_1}$ as illustrated in the schematic of Figure

2-2. Shown is a cross-section of the entire shell, which is a surface of rotation about the v_y axis, and that portion to the left of the critical normal incident velocity will be reflected at the shock front.

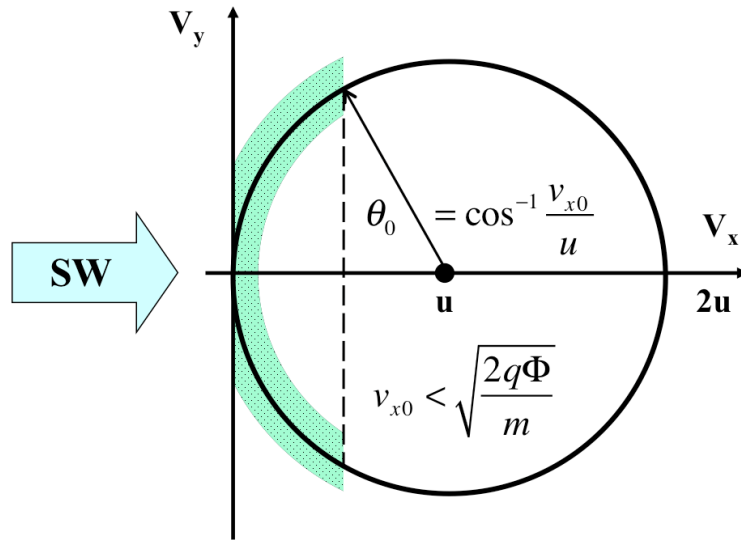


Figure 2-2 Cross section of Pickup Ion shell

Green shaded region shows the population of incident ions that are reflected from a zero-width shock ramp.

At an idealized Mach 5 perpendicular shock, the Alfvén normalized electrostatic potential is approximately $\Phi \approx 10.6$ which for protons gives a maximum normal velocity limit of $v_{x0} = 4.60$ and a reflection coefficient $\mathfrak{R}_{ion} = 0.460$ or about 46%. This simply ensures a single reflection, not multiple bounces required for shock surfing. We show in

Figure 2-3 below a family of curves for reflection coefficients of various masses plotted against the shock Mach number. Note the strong Mach dependence for weak shocks and the asymptotic behavior for moderate to strong shocks. The next diagram Figure 2-4 shows the reflection coefficient curves superimposed upon the incident PI shell. The regions of reflection lie within the intersection of the shell with the region to the left of each curve.

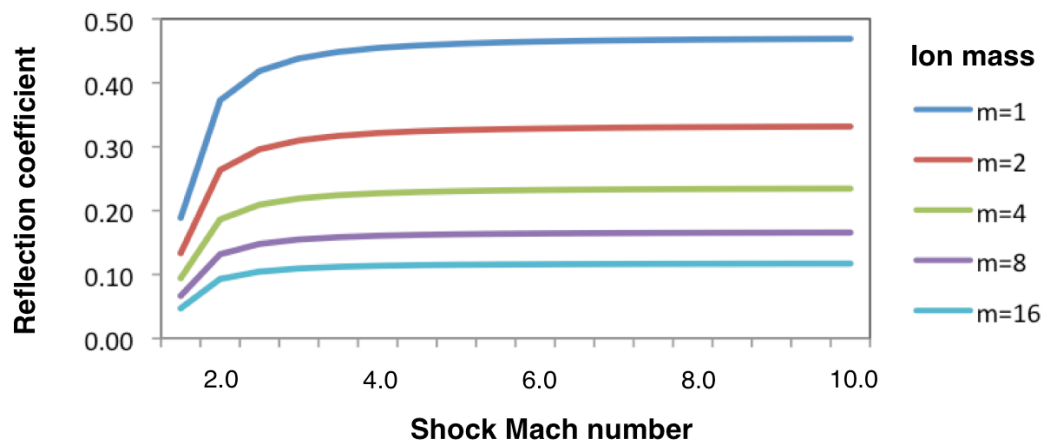


Figure 2-3 Reflection Coefficient vs. Mass

At a step shock of zero ramp width, the coefficient is mass dependent, but insensitive to the shock strength above Mach 5.

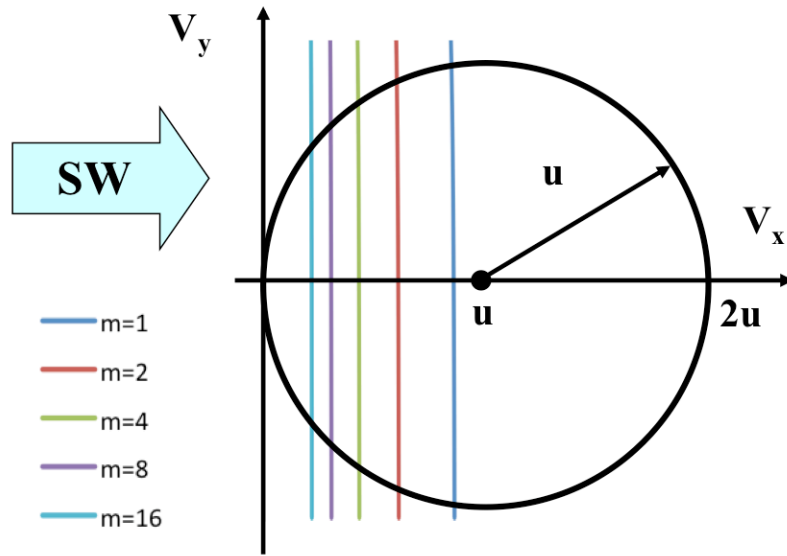


Figure 2-4 Critical velocity boundaries at a step shock.
 Reflected Pickup Ions originate from the left of each vertical line.
 The reflection regions diminish with increasing ion mass.

Finite Ramp

We now extend our analysis to the case of shock ramps of finite extent, and begin with an energy balance equation $q\Phi = \frac{1}{2}mv_x^2 + qv_y B_z L_r$. The left hand term is the energy required to overcome the electrostatic potential across the shock. The first term on the right is the normal incident ion kinetic energy, as we applied previously for the zero-width ramp. The new term represents the work done by the Lorentz force on the ion as it traverses the ramp, more precisely, $\int_0^{L_r} \mathbf{F}_B \cdot d\mathbf{x} = \int_0^{L_r} qv_y B_z(x) dx$. We will assume that ions at this limit traverse the ramp quickly enough that the longitudinal velocity v_y remains nearly constant. We can then replace the magnetic field function in the

integrand with the mean magnetic field strength within the ramp as $\tilde{B} = \frac{1}{2} B_0 (1 + R_c)$

where B_0 is the upstream solar wind field strength and R_c the shock compression ratio.

For an incident ion to avoid transmission through the shock, we must have

$$v_y < \frac{1}{\tilde{B} L_r} \left\{ \Phi - \frac{1}{2} \frac{m}{q} v_x^2 \right\}. \text{ Holding the potential and mean magnetic field fixed for a}$$

particular shock Mach number, we generated a family of inverted parabolae by varying

the ramp width as shown in Figure 2-5 below for H^+ (protons).

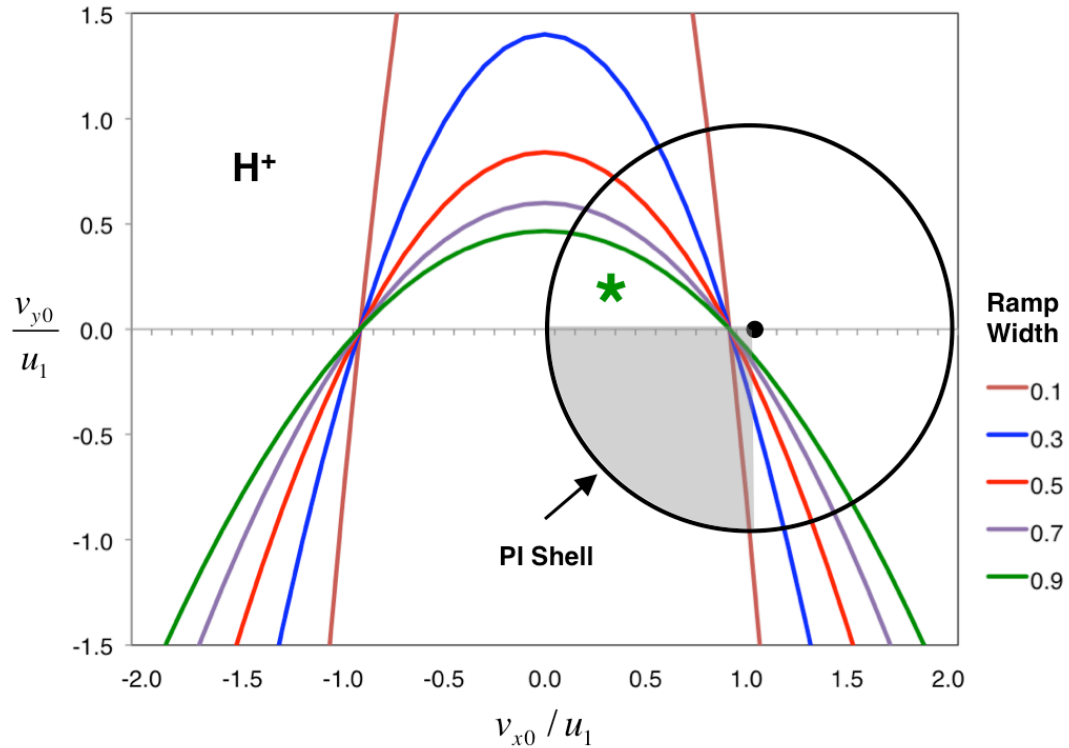


Figure 2-5 Proton Boundaries at Finite Ramp

Inverted parabolae are superimposed on the proton PI shell.

Reflection regions lie under the curve for each ramp width shown as indicated by the green star. The fourth quadrant (shaded region) is forbidden for multiply reflected ions (shock surfing).

Only that portion of the PI population lying below the curves will be reflected. Note that the x-intercept is fixed by the ion mass and charge, and this dependence is strong, as shown by the same family of curves for singly ionized Helium He^+ shown in Figure 2-6. We shall refer back to these dependencies when we analyze our test particle results in the following section.

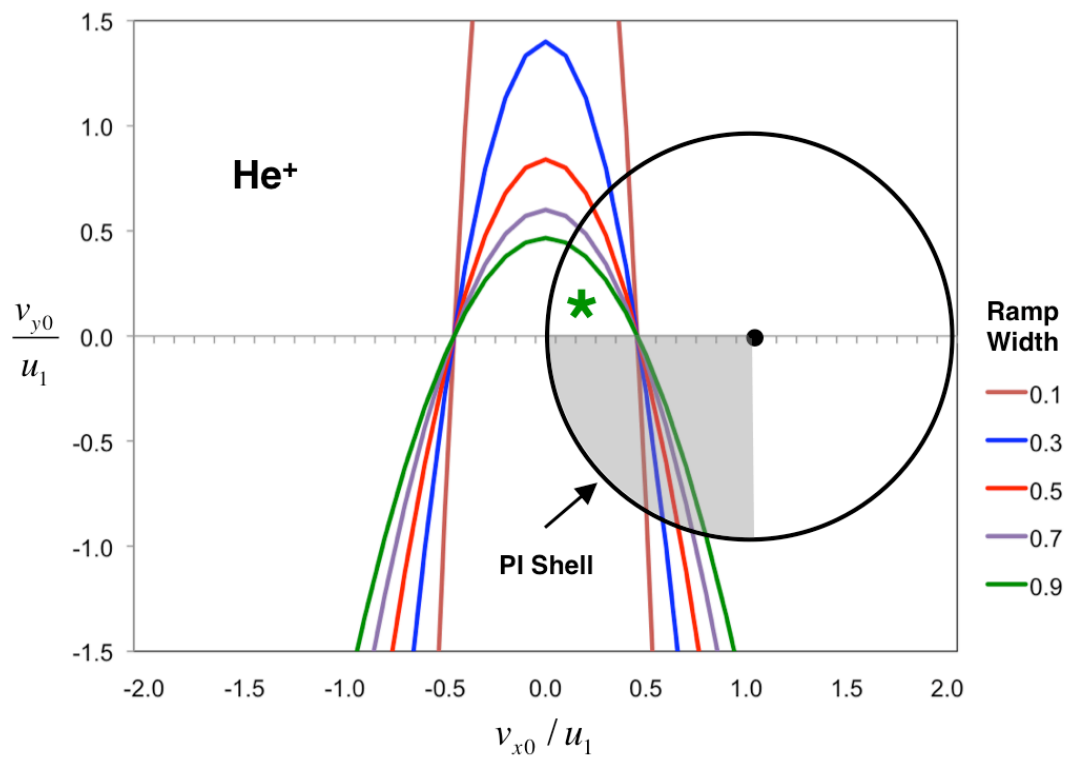


Figure 2-6 Singly Ionized Helium at Finite Ramp

Colored curves show reflection region boundaries. The locus of points at $v_{y0} = 0$ has shifted to the left, and the parabolae are steeper, eliminating a larger fraction from the surfing region in the Pickup Ion shell

It should be mentioned here that these analyses proceed from static models, with no dynamics included, and the results should be taken as qualitative rather than quantitative predictors. For example, it is clear that in undergoing a reflection by virtue of the shock-normal electric field, an ion must both decelerate to zero normal velocity and also penetrate the ramp in doing so. This leads to two results: the latency period for the

fastest incident ions inside the ramp increases with mass, that is $\tau_{lat} = \frac{mL_r}{qE_{\perp}}$ and the magnetic field strength must increase with penetration depth. Together, these imply that, regardless of mass, ions must gain some kinetic energy from the convective electric field while inside the ramp and therefore experience greater Lorentz force by virtue of their acquired velocity. In so doing, the deepest penetrating ions must escape downstream by ‘leaking’ through the shock. Short of a fully dynamic treatment, we may still make an approximation of a correction factor by noting that the penetration depth should scale as

the ion inertial length $\lambda_{pi} = \frac{c}{\omega_{pi}} = c\sqrt{\frac{m_i}{4\pi ne^2}} \sim m^{1/2}$. Note further, that the penetration

depth for any given incident velocity must increase with mass by the same scaling.

Hence, simply from this rough estimate, we may expect that the flux of reflected ions

should scale with mass by not one, but two factors of $m^{-1/2}$, i.e.,

$\mathfrak{R}_{ion}(m) \sim m^{-1/2} m^{-1/2} \sim m^{-1}$ or inversely with mass.

Another point of interest is that, while the fraction of the PI population for a given ion species mass appears insensitive to ramp width, the rotation of the interface boundary impinges upon and ultimately sweeps through the region where Multiply Reflected Ion (MRI) originate. Wider ramps thus deplete shock surfing candidates more dramatically for heavier ions, and we reference this when analyzing arguments in favor

of shock surfing to explain the abundance of heavy ions in the observed Anomalous Cosmic Ray (ACR) spectrum.

Computational notes – Shock Surfing

When setting simulation parameters such as Mach number, ion mass and, in the case of static test particle runs, shock ramp width, one must make reasonably accurate predictions of both time and length scales needed to resolve features of interest. Certainly, the time step (and grid for kinetic hybrid codes) must be fine enough to resolve motion at high velocities or steep field gradients. The simulation box size and duration of the run must also be taken into account.

We can predict the velocity ceiling for ions surfing at a finite-width shock using a force balance approach. The two forces acting on such an ion are the electrostatic due to the retarding electric field within the ramp and the Lorentz force due to the ion's velocity perpendicular to the magnetic field. When the latter is dominant, the ion will be pushed through the front terminating the surfing process. To remain upstream, we require that

$$q \frac{v_y}{c} B < q E_{\perp} \Rightarrow v_y < \frac{c E_{\perp}}{B} = \frac{c \Phi}{B L_r}$$

We can predict the maximum surfing velocity for a given ramp width, once the potential is calculated for a given shock Mach number, and thus choose our time step accordingly. Note that a surfing ion's velocity ceiling is independent of mass, but the time required to attain this maximum velocity is not. The simulation duration can be calculated from the

acceleration due to the convective electric field $\tau = \frac{v_{\max}}{a_{\parallel}} = v_{\max} \frac{m}{q E_{\parallel}} = v_{\max} \frac{m}{q} \frac{c}{u_1 B_1}$ and,

substituting the velocity ceiling expression above, we have $\tau = \frac{m_i}{L_r} \left(\frac{\Phi c^2}{q_i u_1 B_1^2} \right)$. Restricting ourselves to singly ionized species fixes the term within the parenthesis by choice of the shock Mach number and, applying our normalization, we obtain $\tau = \frac{m_i}{L_r} \left(\frac{\Phi}{M_A} \right)$. As an example, a Mach 5 shock gives $\Phi \simeq 10.6$ so that, for protons $m = 1$ and a thin ramp $L_r = 0.05$, the minimum time required for the fastest expected ions to be observed and recorded should be $\tau \sim 42$ but, in fact, we have found that twice this value must be used in order to capture the fastest accelerated ions. The time scale must be extended for adequate statistical resolution based on the flux chosen for PI injection when running steady-state ensembles. From elementary analysis of the theory, one could simply choose to inject a single particle at the base of the ramp with precisely zero velocity, and produce the exact maximum surfing velocity. However, we generate ensembles that evolve through repeated injection of single particles, each of whose initial velocity is calculated to be statistically consistent with that of the PI shell model and, in this way, we ‘fill in’ the velocity-space shell with discrete points. The probability of randomly generating any particular point is precisely zero, but we can come arbitrarily close by choosing sufficient density of points. The desired accuracy of our results then dictates either the minimum injection flux based on the simulation duration or, conversely, the additional time (beyond the minimum required as shown above) in order to achieve the desired resolution.

One last point to consider is our choice of magnetic field strength. While it is clear that the field increases through the ramp to its downstream value $B_2 = R_C B_1$, only those surfing candidates fortunate enough arrive at the foot of the ramp with nearly zero

normal velocity can attain the highest possible velocity. Hence we are justified in taking the upstream field magnitude, since penetration into the ramp cannot occur until the Lorentz force overcomes that due to the constraining electric field. At that moment, transit through the ramp should be rapid enough to ignore further consideration of the increasing magnetic field.

Shock Drift Theory

Method 1. We model an ion that transits the shock normal to the front, makes a half stationary orbit downstream followed by another half orbit upstream as shown in the Figure 2-7 below. For each such complete orbit, there is a net displacement of the guiding center in the direction of the convective electric field

$$\Delta y = 2r_{g1} - 2r_{g2} = 2r_{g1}(1 - R_C^{-1}) = 2\frac{v_{\perp 1}}{\Omega_{g1}}(1 - R_C^{-1})$$

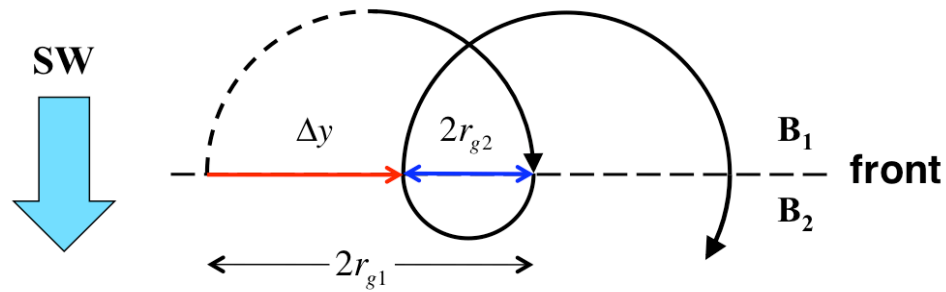


Figure 2-7 Schematic of Shock Drift Trajectory
Ions thread the front once before being swept downstream

Assume that a Pickup Ion is temporarily constrained at the front, that is, ignore the guiding center motion due to the local fluid velocity. If we also ignore the velocity change due to acceleration in the convective electric field, we may roughly estimate that

$v_{\perp 1} \sim u_1$ for the fastest incident PI and obtain the ion's displacement after one initial

encounter cycle $\Delta y = \frac{u_1 m_i c}{q_i B_1} (1 - R_C^{-1})$.

The kinetic energy gain from this displacement along the convective electric field

will be $\Delta \varepsilon = q_i E_{\parallel} \Delta y = q_i \frac{u_1 B_1}{c} \frac{u_1 m_i c}{q_i B_1} (1 - R_C^{-1}) = m_i u_1^2 (1 - R_C^{-1})$. Since the fastest incident

ions in the PI shell have velocity $v_1 = 2u_1$, the final kinetic energy after the gain

associated with this translation is $\varepsilon_f = \varepsilon_i + \Delta \varepsilon$ or $\frac{1}{2} m_i v_f^2 = \frac{1}{2} m_i (2u_1)^2 + m_i u_1^2 (1 - R_C^{-1})$

from which it follows that $\tilde{v}_f \equiv \frac{v_f}{u_1} = [6 - 2R_C^{-1}]^{1/2}$. Applying the compression ratio

$R_C = 2.62$ at a Mach 5 shock gives $\tilde{v}_f \simeq 2.29$ which is low, but reasonably close to that produced by our numerical simulations.

Method 2. Another approach uses the gradient drift, based upon an

assumption that the ion's trajectory threads the ramp many times adiabatically. In order for this to be valid, two conditions must hold: the cyclotron velocity must be greater than the fluid velocity $v_{\perp 1} \gg u_1$, and the guiding center displacement over an ion cyclotron

period must be less than the ramp width $u_1 \frac{2\pi}{\Omega_{ci1}} \ll L_r$. The first condition is never

satisfied by Pickup Ions, and the second is never met at actual physical shocks.

However, this approach has sometimes been put forth, and we offer it as an alternative.

The magnetic gradient drift velocity is given by $v_{\nabla B} = \frac{1}{2} v_{\perp i} r_{gi} \frac{B \times \nabla B}{B^2}$ where $v_{\perp i}$ is

the ion velocity perpendicular to the ambient field, and r_{gi} is the cyclotron radius. Since

the magnetic field strength is proportional to the plasma density, we may write

$$\frac{B \times \nabla B}{B^2} = \frac{R_C - 1}{L_r} \text{ where } R_C \text{ is the shock compression ratio and we assume } B \perp \nabla B$$

(appropriate at purely perpendicular shocks) and a linear profile, that is,

$$\nabla B = \frac{B_2 - B_1}{L_r} = \frac{B}{L_r} (R_C - 1). \text{ Care must be taken in choosing the value of the magnetic}$$

field, which obviously changes through the ramp. However, making this substitution

$$\text{along with } r_{gi} = \frac{v_{\perp i}}{\Omega_{ci}}, \text{ we now have } v_{\nabla B} = \frac{1}{2} \frac{v_{\perp i}^2}{\Omega_{ci}} \frac{R_C - 1}{L_r}. \text{ We can estimate the longitudinal}$$

$$\text{displacement resulting from the drift by } \Delta y = v_{\nabla B} \tau \text{ where } \tau = \frac{L_r}{\langle u \rangle} \text{ is the transit time}$$

$$\text{through the ramp and } \langle u \rangle = \frac{1}{2} (u_1 + u_2) = \frac{u_1}{2} (1 + R_C^{-1}) \text{ the mean guiding center velocity.}$$

$$\text{Substituting } \Omega_{ci} = \frac{m_i c}{q_i B}, \text{ assuming the frozen field condition } uB = u_1 B_1 = \text{constant}, \text{ and}$$

$$\text{applying the mean magnetic field through the ramp } \langle B \rangle = \frac{1}{2} (B_1 + B_2) = \frac{B_1}{2} (1 + R_C), \text{ the}$$

kinetic energy gained by this process is

$$\begin{aligned} \Delta \varepsilon &= q_i E_{\parallel} \Delta y = q_i \left(\frac{u_1 B_1}{c} \right) \left(\frac{1}{2} \frac{m_i v_{\perp}^2 c}{q_i \langle B \rangle} \frac{R_C - 1}{L_r} \right) \frac{L_r}{\langle u \rangle} \\ &= \frac{1}{2} m_i v_{\perp}^2 \frac{4(R_C - 1)}{(1 + R_C)(1 + R_C^{-1})} \\ &= \frac{1}{2} m_i u_1^2 \frac{(R_C - 1)(3 - R_C^{-1})^2}{(1 + R_C)(1 + R_C^{-1})} \end{aligned}$$

For the last equality, we have assumed that the change in ion velocity is small,

and have used the mean cyclotron velocity across the ramp

$$\langle v_{\perp} \rangle = \frac{1}{2} [v_{\perp 1} + v_{\perp 2}] = \frac{1}{2} [(2u_1 - u_1) + (2u_1 - u_2)] = \frac{1}{2} u_1 (3 - R_C^{-1})$$

The energy of the fastest incident ions will be $\varepsilon_{final} = \varepsilon_0 + \Delta\varepsilon$ and again applying

$R_C = 2.62$ for a Mach 5 shock, we obtain the final velocity

$$\tilde{v}_{final} = \left[4 + \frac{(R_C - 1)(3 - R_C^{-1})^2}{(1 + R_C)(1 + R_C^{-1})} \right]^{1/2} \approx 2.49$$

Once again, we remind the reader that this derivation is not strictly warranted at the shocks of interest to us. However, this latter exercise does demonstrate that the energy should be independent of the ramp width, and also independent of mass.

We note that the retarding electric field present at a perpendicular shock has not been taken into account. Nevertheless, we can say that those particles favored by this mechanism are the fastest of the incident PI shell. Not only must they have the greatest cyclotron velocity, but also must be fast enough to overcome the potential barrier and cross downstream. Further, as we have already shown, the velocity loss when passing through the potential gradient is smaller for heavier ions.

Computational notes – Shock Drift

As we have shown, since the acceleration time for ions to reach their maximum surfing velocity $\tau = \frac{v_{max}}{a} = \frac{v_{max}}{q_i E_{\parallel}} m_i$, the simulation runtime scales with mass. However, if we are also interested in comparing relative efficiency or flux contributions of both surfing and shock drift ions, the species' ion cyclotron radius must be also taken into account in setting the box size.

Let us estimate the minimum length of the downstream region of our box, sufficient to guarantee the retention of all those shock drift ions likely to return upstream. As shown by the construction in Figure 2-8, once an ion has traversed the front, the

subsequent evolution of its trajectory is completely governed by guiding center dynamics. The dashed portion of the upstream path is extended as if the magnetic field were everywhere uniform since the shock front marker is merely a graphical construct for computational purposes within one complete cyclotron orbit.

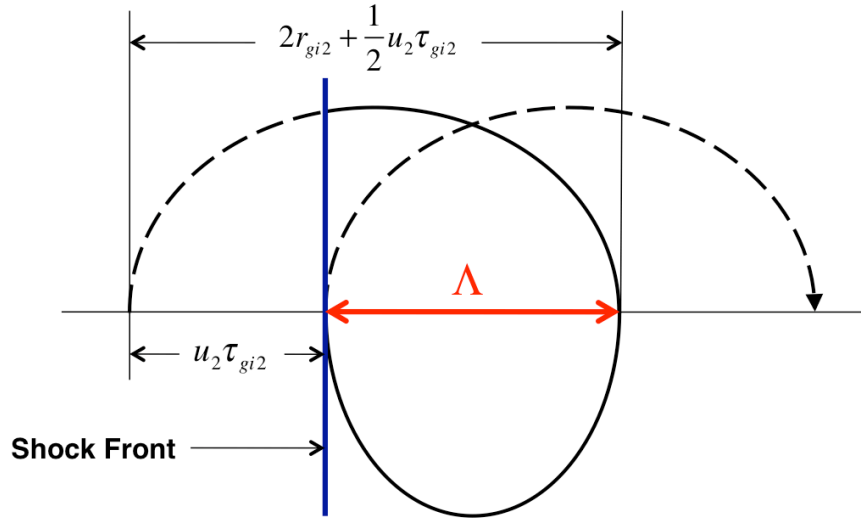


Figure 2-8 Minimum Downstream Box Length
required for 'crossing' ions to be retained and returned to the front

The last possible moment for an ion to return upstream is just prior to the end of its gyrophase when $v_x = 0$. For such an ion, the downstream penetration depth of its

trajectory prior before returning to the front is $\Lambda = 2r_{gi2} + \frac{1}{2}\tau_{g2}u_2 - \tau_{g2}u_2 = 2r_{gi2} - \frac{1}{2}\tau_{g2}u_2$

where the downstream variable u_2 is the fluid velocity, B_2 the magnetic field,

$\tau_{g2} = 2\pi\Omega_{ci2}^{-1}$ the ion gyroperiod and $r_{gi2} = \frac{m_i v_{\perp 2} c}{q_i B_2}$ the ion gyroradius for a downstream

ion with cyclotron velocity $v_{\perp 2}$. We may estimate the downstream cyclotron velocity

using the construction shown in Figure 2-9 below.

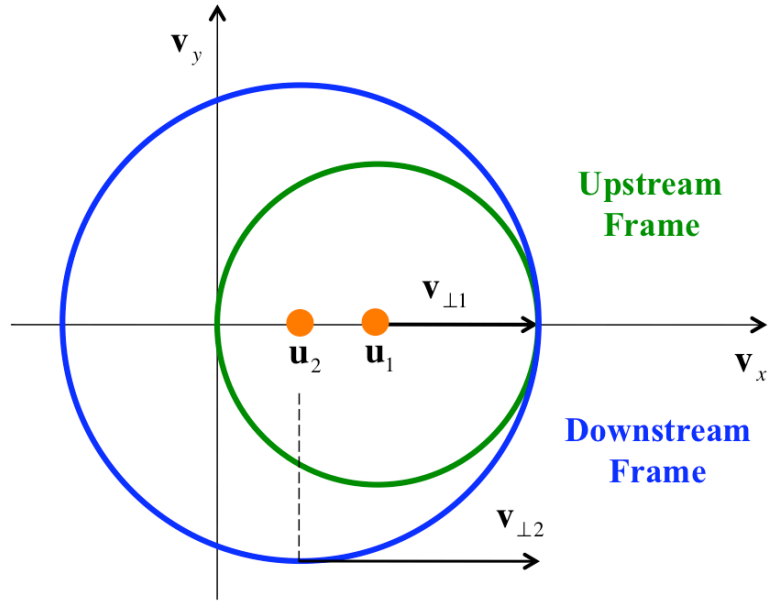


Figure 2-9 Velocity Space Diagrams

The superimposed circles represent two frames of reference

Assuming that the fastest particles encounter the shock with velocity

$\mathbf{v} = (v_x, v_y) \approx (2u_1, 0)$, we see from the graphic that the perpendicular velocity upstream

is simply $v_{\perp 1} = v_x - u_1 = 2u_1 - u_1 = u_1$. After crossing the front, however, the downstream

fluid velocity displaces the velocity shell center so that, holding $v_{2\max} = v_{1\max}$ we have

$v_{\perp 2} = 2u_1 - u_2 = (2 - R_C^{-1})u_1$. However, all transmitted ions must climb the electrostatic

potential 'hill' during their passage through the potential step $\Delta\Phi$ at the ramp, and are

thus subject to a loss of kinetic energy $\Delta E = q\Delta\Phi$. Applying the construction used

previously to calculate the potential, we have $\frac{1}{2}m_i v_1^2 - \frac{1}{2}m_i v_2^2 = \frac{1}{2}m_p u_1^2 (1 - R_C^{-2})$ which

gives $v_2^2 = v_1^2 - \frac{m_p}{m_i} u_1^2 (1 - R_C^{-2})$. Assuming that $v_y \ll v_x$ for the fastest PIs, we

approximate the incident velocity as $v_1 \approx 2u_1$, which gives $v_2^2 = (2u_1)^2 - \frac{m_p}{m_i} u_1^2 (1 - R_C^{-2})$.

Expressed in normalized units, this becomes simply $v_2 = M_A \left[4 - m_i^{-1} (1 - R_C^{-2}) \right]^{1/2}$, the absolute downstream velocity, which gives the corrected downstream cyclotron velocity

$\hat{v}_{\perp 2} = \frac{v_2}{M_A} = \left[4 - m_i^{-1} (1 - R_C^{-2}) \right]^{1/2} - R_C^{-1}$ relative to the upstream wind speed. Note that

this correction is small, as shown in Figure 2-10 below where we have plotted the

variable $\Delta \equiv 1 - \frac{v_2}{2u_1}$, the difference for the fastest incident PIs.

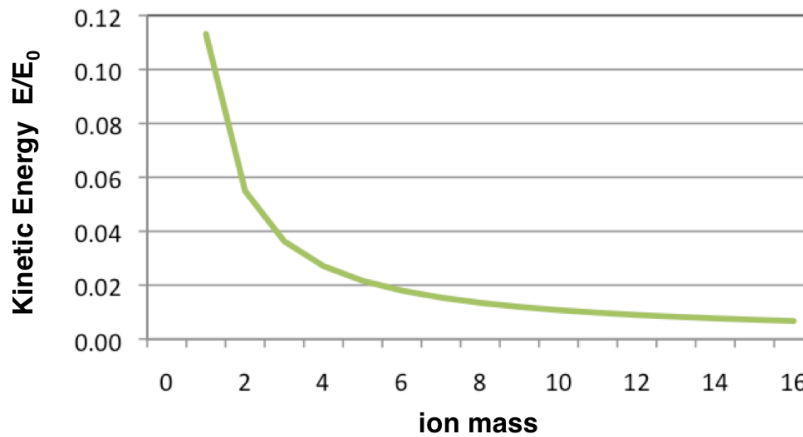


Figure 2-10 Kinetic Energy Loss

Ions lose energy during transit through the electrostatic potential gradient.

The loss is roughly inversely proportional to the ion mass.

The difference is greatest for protons at less than 12% and drops rapidly with increasing mass. Ignoring this correction in any case will merely extend the downstream region

slightly farther than required and, since the extra computation time is relatively small, we will use the approximation $v_2 \approx v_1 = 2u_1$ and reinstate the perpendicular cyclotron velocity estimate $v_{\perp 2} = (2 - R_C^{-1})u_1$. This now fixes the penetration depth since all other variables are uniquely determined by the shock parameters. We can simplify the expression using the substitution $r_{gi2} = v_{\perp} \Omega_{ci2}^{-1}$, exhibit the mass dependence explicitly by $\Omega_{ci2} = m^{-1} \Omega_{c2}$ and convert to upstream parameters via $u_2 = u_1 R_C^{-1}$ and $\Omega_{c2} = R_C \Omega_{c1}$ where once again R_C is the shock compression ratio. Applying the Alfvén normalizations $\Omega_{c1} \rightarrow 1$ and $u_1 \rightarrow \frac{u_1}{v_A} = M_A$ we obtain

$$\tilde{\Lambda} = 2r_{gi2} - \frac{1}{2} \tau_{g2} u_2 = 2 \frac{v_{\perp 2}}{\Omega_{ci2}} - \frac{1}{2} \frac{2\pi}{\Omega_{ci2}} u_2 = \tilde{m}_i M_A R_C^{-1} [4 - (\pi + 2) R_C^{-1}].$$

Hence we see that, to retain all the shock drift ions in the simulation, the minimum downstream box length required scales linearly with the ‘reduced’ ion mass $\tilde{m}_i$, i.e., roughly the integer atomic mass. For a typical Mach 5 shock, the compression ratio is $R_C \approx 2.62$, which gives the minimum downstream excursion length $\tilde{\Lambda} \approx 3.89 \tilde{m}_i$. For protons, we have $\tilde{\Lambda} \approx 3.89$ and for singly ionized Oxygen, $\tilde{\Lambda} \approx 62.2$. Although we have restricted our study to ions of unit charge, we may modify our multiplicative factor to accommodate multiply ionized species by simply replacing the reduced mass by the mass-to-charge ratio $\tilde{m}_i \rightarrow \tilde{m}_i / \tilde{q}_i$.

Runtime considerations

A longer simulation box requires more time for ions to traverse the entire length. But this is a negligible factor for steep shocks in light of the much longer times necessary

for the fastest shock surfing ions to attain their maximum velocity. What is **not** trivial, however, is the effect that the longer box has on the simulation runtimes for heavier species. The population in the box at any given moment is directly proportional to the box length, which in turn scales with mass. We must thus include the mass factor twice – once for simulation time in units of Ω_c^{-1} to accommodate the surfing ion velocity, and once again for the computational load imposed the shock drift return length. Hence the runtime scales as $T_{run} \propto \tilde{m}_i^2$, so that the actual time required to complete a series for O_{16}^+ is $16^2 = 256$ times longer than that for protons. Recall that we have explored a large region of parameter space including variation in ion mass, shock Mach number, magnetic field angle, and ramp width. Holding all other parameters fixed, one baseline series for masses $m_i \in \{1, 2, 4, 8, 16\}$ over a range of ramp angles $0.05 < L_r < 1.00$ in increments $\Delta L_r = 0.05$ takes several days on a typical desktop computer. To complete an initial investigation, the same series must be repeated over a range of shock Mach numbers and field angles, which in all may take several months. And this is just the baseline for simple plane-front shocks. To extend the study to non-ideal shocks, the entire range must be repeated for various types of imposed waves and perturbations. Expediency is required.

We offer one last analytic prediction before proceeding. It is essential when performing simulations to ensure that the code is as accurate as possible, given the physics we wish to include. Test suites exist for this very purpose when using kinetic hybrid and other self-consistent codes, and include tests for energy conservation among many others. However, when running test-particle simulations in the absence of self-consistent active backgrounds, one is free to choose the fields at will, and bizarre results

can ensue without due diligence. Therefore, it is not entirely an academic exercise to make theoretical predictions to which we can compare our results and thereby feel confident in their veracity.

The construction we employed above to compute the maximum downstream SD orbital depth also allow us to calculate the gyrophase angle corresponding to the transit velocity of SD ions so that, working backwards we can estimate the fraction not only the fraction of the total PI population which should become shock drift species, but also their location within the shell.

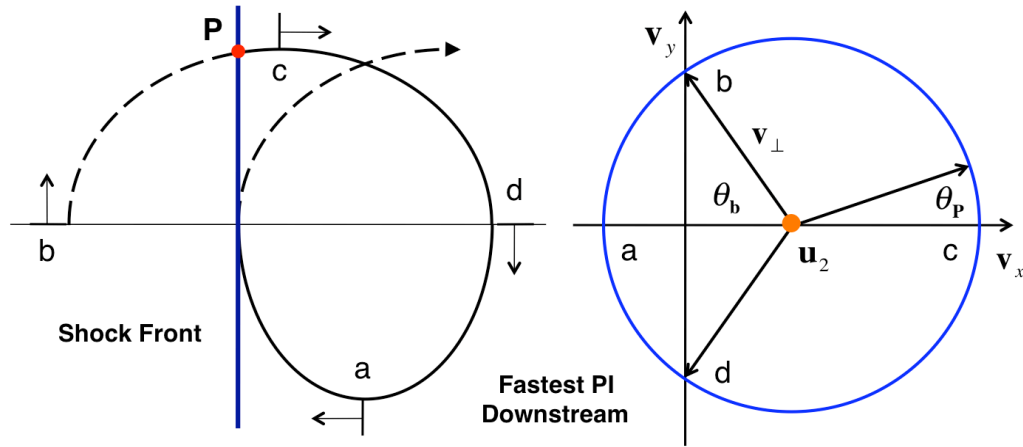


Figure 2-11 Schematic phase space construction
to compute boundary of the downstream shock drift return region

We show in Figure 2-11 above a guiding center trajectory for a downstream ion as a plot in position space along side the corresponding velocity space plot. We label the points **a**, **b**, **c** and **d** to identify the corresponding gyrophase angle $\theta(t) = \Omega_{gi2}t$. The radius of the downstream PI shell ‘remnant’ we will take for the moment to be approximately that of the upstream velocity shell $v_{\perp 2} \approx v_{\perp 1} = u_1$.

To find the angle of attack for these fastest PIs which cross the front on their initial shock encounter, we begin with the downstream velocity equations

$$v_x(t) = -v_{\perp 2} \cos(\Omega_{ci2} t) + u_2 \quad \text{and} \quad v_y(t) = -v_{\perp 2} \sin(\Omega_{ci2} t). \quad \text{Integrating these gives the}$$

$$\text{position equations } x(t) = -\frac{v_{\perp 2}}{\Omega_{ci2}} \sin(\Omega_{ci2} t) + u_2 t \quad \text{and} \quad y(t) = -\frac{v_{\perp 2}}{\Omega_{ci2}} \cos(\Omega_{ci2} t). \quad \text{We note}$$

two conditions that hold: 1) when an ion's trajectory returns to the front at point **b**, its horizontal velocity vanishes, and 2) at the point **P** where the trajectory first crosses the front, the horizontal position $x(\theta_p) = x(\theta_b) + 2\pi$. Expressing these mathematically, we

$$\text{have for the velocity } -v_{\perp 2} \cos(\theta_b) + u_2 = 0 \Rightarrow \cos(\theta_b) = \frac{u_2}{v_{\perp 2}} \quad \text{and, multiplying the}$$

displacement equation by $-\frac{\Omega_{ci2}}{v_{\perp 2}}$, we have the position condition

$$\tilde{v} \sin(\theta_p) - \theta_p = \tilde{v} \sin(\theta_b) - (\theta_b + 2\pi) \quad \text{where} \quad \tilde{v} \equiv \frac{v_{\perp 2}}{u_2}. \quad \text{Assuming from the construction a}$$

shallow angle of attack, we may apply to the gyrophase angle the identity

$$\sin \theta_p = \sin(\pi - \theta_p) \quad \text{and the small angle approximation } \sin(\pi - \theta_p) \simeq \pi - \theta_p, \quad \text{and write}$$

$$\tilde{v}(\pi - \theta_p) - \theta_p = \tilde{v}\sqrt{1 - \cos^2(\theta_b)} - \theta_b = \tilde{v}\sqrt{1 - \tilde{v}^{-2}} - \cos^{-1}(\tilde{v}^{-1}) - 2\pi, \quad \text{or}$$

$$\theta_p = \frac{\tilde{v} \left[\pi - \sqrt{1 - \tilde{v}^{-2}} \right] + \cos^{-1}(\tilde{v}^{-1}) + 2\pi}{\tilde{v} + 1}.$$

Applying the rough estimate $v_{\perp 2} \simeq v_{1\max} - u_2 = u_1(2 - R_C^{-1})$ gives $\tilde{v} \equiv \frac{v_{\perp 2}}{u_2} = 2R_C - 1$. At

Mach 5, the shock compression ratio $R_C \simeq 2.62$ which gives $\tilde{v} \simeq 4.24$ and we find that

$$\theta_p \simeq 3.21. \quad \text{Checking, we find that } \frac{\sin \theta_p}{\pi - \theta_p} \simeq 0.977, \quad \text{justifying our simplifying}$$

approximations. If we wish to be more precise, we must return to our calculation for the downstream velocity corrected by the kinetic energy loss while crossing the potential

step, $\tilde{v} = \frac{v_{\perp 2}}{u_2} = R_C \left[4 - m_i^{-1} (1 - R_C^{-2}) \right]^{1/2} - 1$. This gives $\tilde{v} \approx 3.65$ for protons at Mach 5,

and differs from our previous value by roughly 14%. Applying this value yields the angle

$\theta_p \approx 3.34$, which gives us a small improvement, since we now have $\frac{\sin \theta_p}{\pi - \theta_p} \approx 0.986$.

Based on this derivation, we conclude that the angle $\theta_p > \pi$ and expect that the most energetic shock drift ions should originate in the third quadrant of the incident PI shell (as seen clockwise), since those residing in the second quadrant cannot return to the front. To take the analysis further would require solving the equation with the trigonometric functions exactly by numerical methods, but we believe the accuracy of our rough prediction should prove sufficient for use in validating the simulation results, and therefore satisfactory as one component of our test suite.

Finally, for completeness we calculate the ‘kink’ angle in the trajectory through

the ramp $\Delta\psi = \sin \frac{\Delta v}{v} \approx \frac{\Delta v}{v}$ since the Lorentz force ensures that $\Delta v \perp v$. The

perpendicular velocity acquired is $\Delta v = \frac{F_B}{m} \Delta\tau = \frac{qv \langle B \rangle}{mc} \frac{L_r}{v} = \frac{q \langle B \rangle L_r}{mc}$ where

$\langle B \rangle = \frac{1}{2} (B_1 + B_2) = \frac{B_1}{2} (1 + R_C)$ is the mean field through the ramp. We note incidentally

that Δv is independent of v . Taking the fastest incident ion velocity as $v \sim 2u_1$, we

obtain the kink angle $\Delta\psi \approx \frac{\Delta v}{v} = \frac{(1 + R_C)}{4M_A} \frac{L_r}{m}$ expressed in reduced units. As we have

shown, the width limit for shock surfing $L_r \leq 1$ to be of any consequence so that, at a

typical Mach 5 shock, protons should deviate from their incident upstream trajectory by roughly 10 degrees at the broadest ramp (as we will show in the baseline study), but less than a single degree at the steep shocks of interest to this investigation. Hence, we could explicitly calculate the incident attack angle ψ and obtain the precise region of interest within the PI shell, but for the ramp widths of interest this should be unnecessary.

We stress that the preceding discussions are not intended as an exhaustive exposition of the theoretical foundations for these mechanisms. Rather, our purpose is simply to use the results to guide our choices of parameters, and to check the validity of our numerical results. We shall now see how well our theoretical predictions agree with simulation data.

Chapter 3 Numerical simulations

In this chapter we numerically investigate the performance of Shock Surfing (SS) on pickup ions (PI) at idealized static quasi-perpendicular shocks. By means of test particle simulations, we explore the distribution functions, maximum attainable velocities and fluxes of surfing ions promoted by this mechanism over a range of parameters including ion mass, ramp width, field angle and shock Mach number. For comparison, we do the same for the Shock Drift (SD) mechanism and conclude that SS is only significantly more effective than SD when the field angle is very nearly perpendicular and the ramp extremely narrow. We also show that SS is far less effective in accelerating heavier ions, and quantify all of these parametric dependences. In so doing, we establish a preliminary set of baseline data, which we apply in Chapter 4 to quantify the effect of shock front perturbations. We also demonstrate that, even for smooth and quiescent shocks, spatial and temporal characteristics common at those within the heliosphere place severe constraints upon the functioning of the SS mechanism, and subsequently prove unequivocally from observational data in Chapter 6 that conditions necessary for effective surfing are violated.

We begin with a brief discussion of our initial numerical methods using kinetic hybrid code, and define the variables, scales and normalizations that we adopted and subsequently applied in test particle simulations thereafter.

Kinetic Hybrid Code

Since our simulations ran on inexpensive personal computers with limited power and resources, we were challenged throughout the study by the constant need for efficiency and economy of effort. As such, we took a reconnaissance approach during

the initial stages in order to establish which parameter ranges would reveal essential features of interest. We continued this approach prior to delving deeply into costly runs, which might otherwise offer little promise of meaningful results. This also provided the opportunity to develop run-time diagnostics specifically tailored to explore our observations in greater depth.

We began by revisiting previous self-consistent studies to validate our code structure and diagnostics, by comparing our results with those of our predecessors. This also enabled us to determine the viability of applying different approaches for producing shocks of interest that have actually been observed *in situ*. We employed a proven one-dimensional particle-in-cell kinetic hybrid code with ion macroparticles and massless electron fluid. Such codes have been used successfully for decades to explore a variety of space plasma physics phenomena, as reviewed most recently by Winske et al (2003) “These codes are used to model phenomena on ion inertial and gyro-radius scales, which fall between longer scales obtained by magnetohydrodynamic simulations and shorter scales attainable by full particle simulations.”

In a hybrid code, ions are pushed forward in time by the Lorentz force in a leapfrog scheme. The simulation ‘box’ is partitioned by a grid or ‘mesh’ into cells and, in one-dimensional code, these retain only the x -coordinate. The force on each particle is calculated by applying electric and magnetic fields local to its cell. After all particles are moved, the density, current, pressure and other moments are gathered for each cell in the mesh and fields re-calculated from Maxwell’s equations via a tri-diagonal solver. In this way, the system evolves self-consistently according to the applied physics, in this case, abbreviated Maxwell’s Equations omitting the induction term. Several commonly

used features were included to ensure accuracy and stability of the code, such as tent functions for computing field values more precisely at ion locations within each cell.

Initial runs employed a one-dimensional box closed at the right hand boundary and open at the left. The box was filled with ions flowing toward the closed end and replenished by continuous injection from the left. Each shock so formed moved from the closed boundary towards the left. Although simple to implement, this approach had the disadvantage of limiting the duration of the run, which ended once the shock traversed the entire length of box. A steady state scheme was employed thereafter in order to keep the shock centered in the box for an indefinite time. Downstream parameters were calculated using the Rankine-Hugoniot relations and approximate density and field profiles generated. After initialization, the shock evolved into a completely self-consistent profile that remained quasi-static. Although our model was one-dimensional and thus required retention of only one spatial coordinate, the ions were treated as fully kinetic i.e., all three components of velocity were used by the particle pusher in order to accommodate oblique shocks. Particles were injected from the left side of the box at $x = 0$ and discarded once they reached the right-hand boundary.

Initial ‘exploratory’ runs were performed and, in order to conserve computational resources, the box size was set just large enough to accommodate the essential features of the shock. Time step and cell widths were set an order of magnitude finer than those required to resolve the time and length scales anticipated and verified during initial runs. A comprehensive test suite of diagnostic routines specific to the active background ions was included to monitor conservation of various parameters – such as total system energy, fluxes, particle number and so forth, to ensure validity of the computational results.

Once the shock stabilized, Test Particles were injected at each time step by a random generator which produced velocities statistically consistent with the standard Pickup Ion shell distribution, i.e., a shell in velocity space with center about, and radius equal to the upstream solar wind speed, as previously shown. The shell was ‘thermally’ thickened by $\pm 5\%$, but its contributions to background density, fields, etc., were not taken into account. In other words, the simulation in that sense was not entirely self-consistent. All three spatial coordinates for each test particle were retained for diagnostic purposes in future simulations where waves were introduced. Typically, the simulation box contained 10^5 particles at any given time. The simulations were run for at least $100 \Omega_{ci}^{-1}$ once the shock had settled, which for the stable massless electron code gave a total number of 10^7 particles participating during each run. This was not true later in the case of the finite-mass code, where the simulations terminated prematurely due to loss of shock structure within the box, runaway energy or other divergent parameters.

Normalizations & Parameters

Throughout our simulations, we applied Alfvén normalization and worked entirely in dimensionless variables: time was normalized to the upstream ion cyclotron frequency $t \rightarrow \tilde{t} = t \Omega_{ci}$, displacement to the ion inertial length $x \rightarrow \tilde{x} = x / \lambda_{pi}$ where $\lambda_{pi} = c / \omega_{pi}$, velocity to the upstream Alfvén speed $v \rightarrow \tilde{v} = v / v_A$, mass and charge to the proton mass and charge $m \rightarrow \tilde{m} = m / m_p$, $q \rightarrow \tilde{q} = q / e$ and magnetic field to the upstream value $B \rightarrow \tilde{B} = B / B_1$. In this norm, the upstream bulk fluid velocity $u \rightarrow \tilde{u} = u / v_A$ simply becomes the Alfvén Mach number M_A , the upstream magnetic field is unity, the

downstream field simply becomes the shock compression ratio R_C , and the mass and charge are roughly the atomic mass and integer ion charge state respectively.

Table 3-1 Upstream Solar Wind Parameters

Ion Density	$n_i = 5 \times 10^6 m^{-3}$
Fluid Velocity	$u_1 = 5 \times 10^2 km \cdot s^{-1}$
Magnetic Field	$B_1 = 5 \times 10^{-9} T$
Ion Temperature	$T_i = 10 eV$
Electron Temperature	$T_e = 50 eV$
Plasma Gamma	$\gamma = 2$
Plasma Beta	$\beta_i + \beta_e \sim 2$

Table 3-2 Derived Parameters

Debye Length	$\lambda_D = 24 m$
Ion Plasma Frequency	$\omega_p = 2.9 \times 10^3 rad \cdot s^{-1}$
Cyclotron Frequency	$\Omega_{ci} = 0.48 rad \cdot s^{-1}$
Ion Inertial Length	$102 km$
Ion Cyclotron Radius	$1040 km$

The free parameters at our command were the shock Mach number, magnetic field angle, ion charge and mass, particle density, plasma beta, electron polytropic index, time step and box length. To model the Earth's bow shock, we originally set the polytropic index $\gamma = 2$ and both the electron and ion beta $\beta_e = \beta_i = 0.5$, giving a total plasma beta of unity. The choice of index merely changes the shock compression ratio

R_C given by the Rankine-Hugoniot relations from which we constructed our profiles as shown in Figure 3-1.

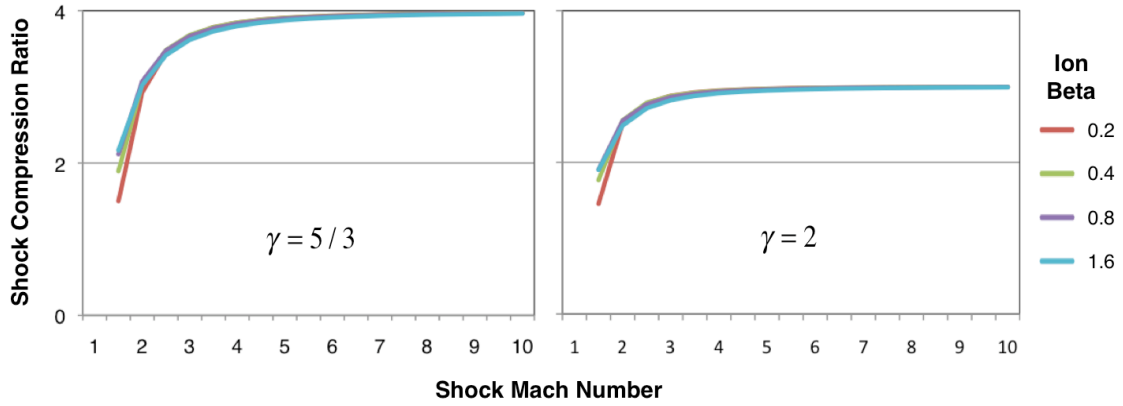


Figure 3-1 Shock Compression Ratio

Shock compression ratio plotted against Mach number for two Polytropic indices $\gamma = 5/3$ (left) and $\gamma = 2$ (right) at a perpendicular shock $\theta = 90^\circ$.

The family of curves demonstrates the weak dependence on plasma beta.

We claim that this has little effect on the shock surfing test particles, since that mechanism is primarily governed by the electrostatic potential Φ and especially the associated retarding ramp-normal electric field $E_\perp = \Phi / L_r$. The normalized potential

function $\Phi = \frac{1}{2} M_A^2 (1 - R_C^{-2})$ clearly illustrates weak compression ratio dependence

compared with very strong Mach dependence. And, since the critical normal velocity also carries the same dependence, so do the reflection coefficient and shock drift flux.

What would be affected is the shock drift velocity ceiling, which depends on the compression ratio. Choosing $\gamma = 2$ imposes a smaller compression ratio, which weakens the SD mechanism and lowers maximum velocity. However, this merely makes

our comparison of relative flux and therefore relative efficiency even more conservative, and we offer additional plots below which illustrate this point more clearly. Figure 3-2 shows the very small effect that the choice of polytropic index makes on the shock electrostatic potential step.

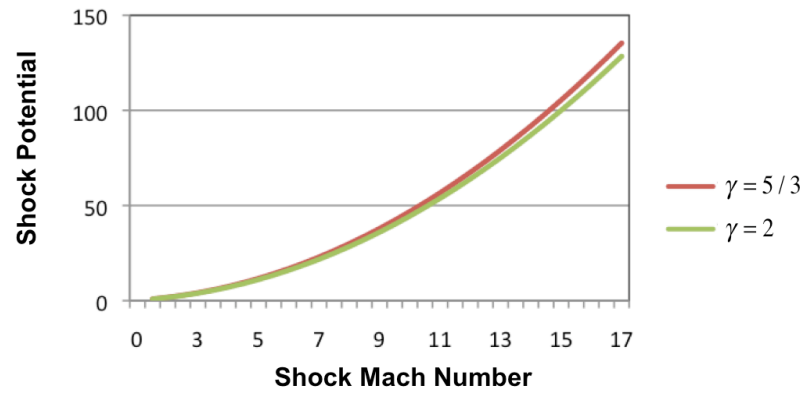


Figure 3-2 Shock Potential vs. Mach Number

Two polytropic indices $\gamma = 2$ and $\gamma = 5/3$ are shown by the curves for a purely perpendicular shock $\theta_{B\hat{n}} = 90^\circ$

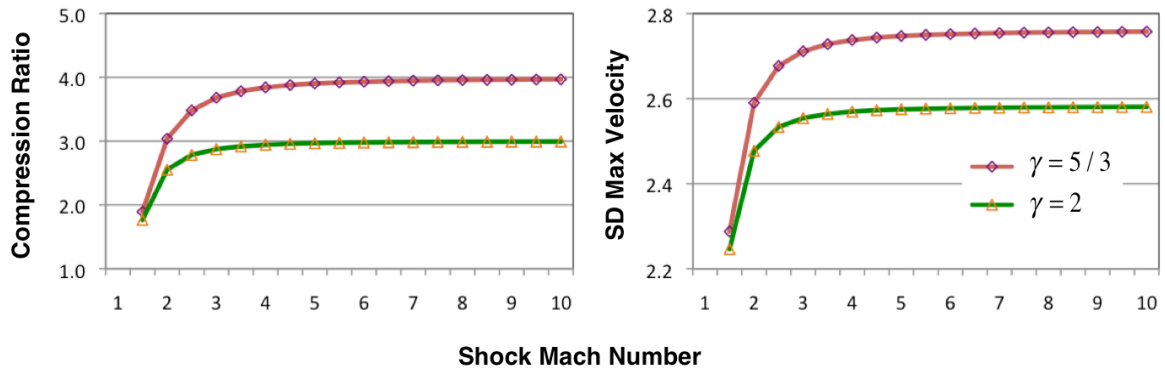


Figure 3-3 Effect of Polytropic Index

Compression ratio (left) and Shock Drift ceiling (right) for two polytropic indices $\gamma = 2$ and $\gamma = 5/3$ with velocity normalized to SW speed.

Diagnostics

Previous computational studies, both self-consistent and test-particle, analyzed the distribution functions for all particles, drawing inferences from the accelerated tails. The best we were able to accomplish with this approach was to verify the work of previous authors – a velocity space plot showing a ‘jet’ of accelerated particles, albeit rather feeble due to the broad ramp, as seen by Lipatov and Zank (1999).

Instead, we chose to ‘tag’ ions at their first encounter with the shock front recording their initial position in velocity space. We followed their trajectories, recording separately for each particle the number of shock crossings and reflections that occurred within the ramp. Multiple shock crossings indicate surfing ions, which we discriminate from those undergoing only single reflections. Multiple shock transits indicate those that have returned upstream before being swept back through the front and these trajectories are perhaps closest to those of early SD models. However, as shown in Figure 3-4, ions we name ‘single-bounce’ follow essentially the same trajectories, threading magnetic field lines within the ramp, while not actually crossing into the downstream region.

Separate distribution functions were accumulated for each category of tagged ion. In this way we were able to identify the behavior and contribution of each specific mechanism, follow individual orbits and determine in which regions of the incident pickup ion distribution the best acceleration candidates reside. These test particles diagnostics were initially applied separately to ions within disjoint regions partitioning the simulation box – far upstream, far downstream and the near-shock or ‘magnetosheath’ region during trial runs.

Of primary interest in light of previous investigations were the test particle distribution functions, but we also kept track of the number of reflections and shock

transits for each particle, binning the results. These were recorded continuously throughout each region for each ‘tag’ category of ion. In order to economize our computational resources, we employed time-averaged integration in place of large test particle number within the box. The validity of this approach is supported by Birkhoff’s theorem, which equates ensemble average with time average in the limiting case. Since the test particle field contributions were not applied, this introduced no anomalous artifacts in the shock structure.

It should be noted that the pickup ion density is estimated to be far below that of the background solar wind population and, since the exact flux has not been established by direct observation, we will proceed under the assumption that the overall shock character is only weakly sensitive to this open parameter. This is actually not true and has been recently addressed in the literature (Winske et al 2009, Wu et al 2009) as we discuss at length in Chapter 7. However, for this study, we seek to isolate the shock from any feedback effects by accelerated ions in order to better study the shock surfing mechanism per se.

Shortcomings of Kinetic Code

An essential element of this tagging scheme is a precise method for defining and locating the shock front. This part of the code is required to identify and distinguish reflected and returning ions. Even for quasi-stationary shocks, changes in ramp profile, overshoot oscillations, and periodic reformation made the task extremely difficult and unpredictable.

While the massless electron code produced a reasonably stable shock profile, and some physically observed bow-shock features were reproduced including foot,

ramp, overshoot and downstream magnetic field oscillations, the shock ramp was quite broad, roughly an ion inertial length. Not surprisingly, this is same length scale produced by applying the KDV equation to model collisionless shocks, as proposed by Kadomsev & Karpman (1971). One can readily show that this scale is necessarily restricted to an ion inertial length while, in fact, electrons play a key role at perpendicular shocks and must therefore be taken into account.

As we have shown, the maximum velocity attainable via shock surfing is inversely proportional to the ramp width $v_{\max} \sim L_r^{-1}$ so that the length scale produced by this code was extremely limiting. Although some acceleration was observed, the maximum surfing ion velocity $v_{\max} \leq 3u_1$ was meager – only slightly greater than the incident shell maximum of $2u_1$, and roughly the same as that produced by shock drift. As a result, and in view of the much steeper shocks that have been physically observed, we abandoned the zero mass code.

Shock Instability

Subsequent attempts to create steeper shocks using a finite-mass code were unsuccessful since tagging became impossible once the shock became uncontrollably active or ambiguously defined. Although the finite mass code produced sharper shock transitions, plasma beta and Mach number still determined the shock ramp width. The shocks produced were highly variable, oscillating in both front position and profile, and attempts to further narrow the shock width by varying the anomalous resistivity merely exacerbated the oscillations, making the shock unstable to the point where it rapidly vanished from our simulation box altogether.

Simulations as far back as 1984 (Quest) have shown the instability of quasi-perpendicular shocks. Indeed, although these shocks are to a great extent self-regulating, oscillations, overturning and reformation have been observed on time scales of the order of the ion-cyclotron period. Matsukiyo & Scholer (2006) performed fully kinetic 2D simulations for a high Mach perpendicular shock and observed six different instabilities develop in less than one gyro-period. More recently, Yang et al (2009) investigated non-stationary supercritical shocks in more detail, re-confirming the same periodicity and extending the analysis to include the responsible mechanisms.

These instabilities may not simply be artifacts restricted to numerical simulations. Recent *in situ* Cluster observations have confirmed their existence at the terrestrial bow shock (Bale 2005). Kucharek (2007) notes from Cluster that foreshock interaction of back streaming ions with the SW causes instabilities that result in wave activity, which then convect back and in turn mediate the shock. Indeed, during two years of observations, the Cluster mission provided only 48 clean, well-defined perpendicular shock crossing examples, suggesting that enduring, long-term stability may not be common. In addressing the problem of DSA at turbulent, rippled shocks, Giacalone & Neugebauer (2008) comment “because of the interaction of the shock with plasma turbulence, the jump in plasma density varies in time and from place to place on the shock front.”

While these observations suggest that real shocks may be too active to support shock surfing, we will proceed under the assumption that shocks may indeed exist somewhere, which are both strong enough and stable enough to allow shock surfing. We now turn to a simpler, more controlled method of inquiry in order to investigate and quantify the potential and limitations of shock surfing under more idealized conditions.

Baseline Studies

By using a static shock model rather than a self-consistent scheme, we were able to vary parameters at will over a wide range of values and thus study the effect of each independently while holding others constant – *Ceteris Paribus* as the Economists say. These constraints are unphysical since, for example, the steepness of the ramp is a function of Mach number, but the model should still prove useful for exploring the behavior and revealing the limitations of shock surfing as functions of these parameters.

We construct an idealized shock profile consistent with the input parameters by applying the Rankine-Hugoniot relations to fix downstream field values. Since these relations do not provide information about the potential step, we approximate this by setting the electrostatic work required for an ion to cross the shock front equal to the

difference between its upstream and downstream kinetic energy, $e\Phi_E = \frac{1}{2}m_p(u_1^2 - u_2^2)$

where m_p is the proton mass and u_1 the upstream fluid velocity. We can express the

estimated potential in reduced units as $\Phi_E = \frac{1}{2}M_A^2(1 - R_C^{-2})$ where R_C is the shock

compression ratio and M_A the Alfvén Mach number. We calculate the normal electric

field in the ramp by assuming a linear profile, $E_\perp = \Phi / L_r$, but tested the effects of

several variations including a sine-squared profile, which we found appealing since it is

$\mathbb{C}_1$ differentiable, e.g., $E_\perp(x) = 2\frac{\Phi}{L_r}\sin^2\left(\frac{\pi x}{2L_r}\right)$, which satisfies $\int_0^{L_r} E_\perp(x)dx = \Phi$.

Differences in the test particle distribution functions resulting from these schemes were negligible, and we therefore adopted the linear profile for the sake of computational efficiency.

All six phase space coordinates were used in the test particle simulations. This was especially important later when artificially induced waves were imposed, increasing the dimensionality of the simulation. In those cases, test particles were injected from the left side of the box, but ‘dithered’ in lateral position over the longest applied wavelength to avoid non-physical artifacts when perturbations were introduced. Because we could define the ramp more sharply and thus expected to attain higher velocities, we replaced the particle pusher with a fourth-order Runge-Kutta solver, and ensured sufficient spatial and velocity resolution when needed by sub-cycling. The fields prior to each push were calculated analytically *de novo* for the particle’s position, including intermediate time steps

Profile and Parameters

Our model is extremely simplified, consisting of only the ramp and upstream and downstream regions and no foot or overshoot were included. Indeed, a ‘real’ shock foot is long and shallow, populated by reflected (gyrating) pickup ions, with negligible normal electric field. We found that the net effect is only an expansion and rotation of the PI shell distribution prior to ions encountering the shock ramp itself.

Hence, since the density in velocity space decreases as the shell expands, our model is conservative in the sense that an extended foot region should, for a given shock step, diminish the surfing efficiency as the ‘support’ (region in velocity space) of the surfing mechanism becomes depopulated. It is, after all, the electrostatic potential gradient and associated retarding field that does the work, and that remains unchanged by foot and overshoot.

We conducted studies of smooth shocks over a range of Mach numbers $M_A = 3, 5, 7, 9$, ramp widths $L_r = 0.05$ through $L_r = 2.0$ and ion masses $m_i = 1, 2, 4, 8, 16$. We applied plasma parameters expected at the magnetosphere, $\beta_e = \beta_i = 0.5$ but with electron polytropic index $\gamma = 2.0$, the effect of which has already been addressed.

Tagging Particles

Each particle in the simulation was “tagged” in several ways. Its position in velocity space was recorded as it first encountered the shock ramp and its progress monitored as it made its way downstream through the ramp and beyond. A bona fide reflection was recorded for a particle if its shock-normal velocity vector changed direction and became negative inside the ramp, and each particle’s ‘bounce number’ bin so incremented. We also counted the number of shock crossings. We categorized particles accordingly; single reflection, multiple reflections and crossing were labeled Single Bounce (B), Shock Surfing (SS) and Crossing (X) respectively. We also distinguished between ions reflected with positive versus negative tangent velocity (in the shock plane), the latter Looping (L) or threading the ramp, but found these to have the same trajectories and distribution functions as the Crossing ions, and have grouped them together as an aggregate mechanism. The rest of the ions, which did none of the above, were simply labeled Shock Drift (D).

Trajectories

The trajectory plots shown in Figure 3-4 below illustrate what appear to be three distinct acceleration mechanisms. On the left we show examples of single particle trajectories while the right represent ensembles. We name the particles shown in the top

row ‘surfing’ as they are tagged as multiple reflections. The next row are discriminated as ‘single bounce’ as they are reflected only once, and the third row are ‘crossing’ as they make a return transit upstream across the front. Not shown are trajectories for the rest of the pickup ions which do none of the above but, nevertheless are accelerated by shock drift displacement as they traverse the front. Note that there is really no difference in principle between the “single bounce” and “crossing” ions as they both make retrograde loops somewhere inside the ramp. The only distinction is that the single-bounce particles do so entirely within the magnetic gradient of the shock ramp.

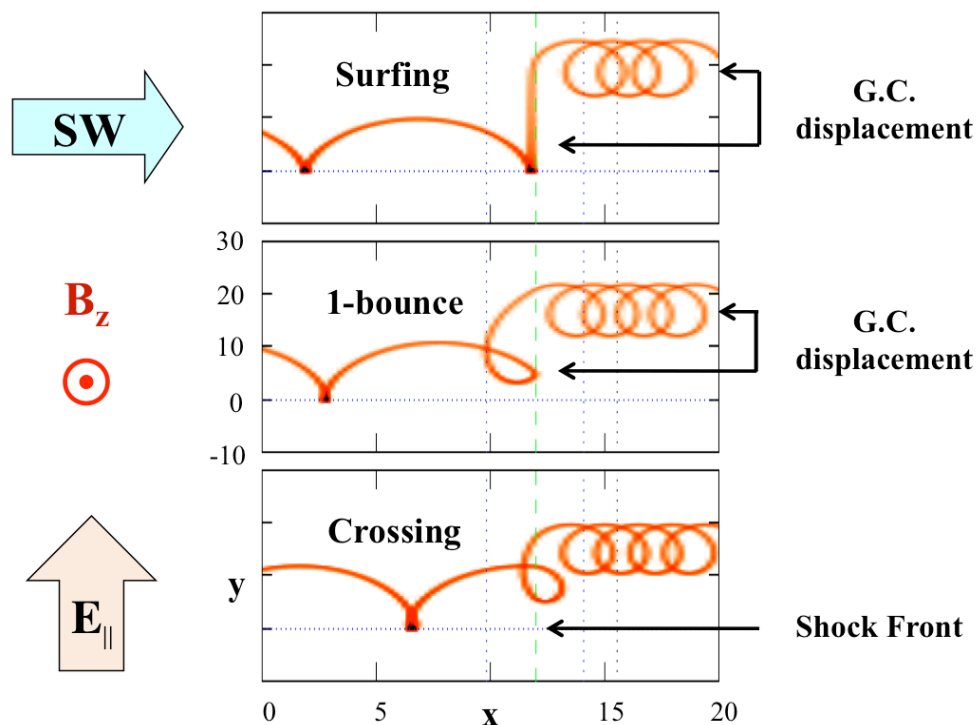


Figure 3-4 Test Particle Trajectories

Three distinct mechanisms – surfing (multiple reflections), single-bounce and ‘crossing’ as discriminated by tagging. The latter two are both shock drift, differing only in the amount of ramp enclosed by the trajectory loop.

Ramp Width Velocity Limit

The best candidates for attaining the highest velocities via shock surfing are those that encounter the front with the smallest normal velocity. Ions that arrive with some normal velocity will ‘bounce’ as they pick up tangential speed, and that oscillation velocity will increase as the normal displacement decreases. Clearly then, the very best candidate would be an ion with zero normal incident velocity which would simply slide along the front and gain energy without bound at an infinitely thin ramp. At a finite width ramp, however, the Lorentz force will push it through. Hence for a given magnetic field and shock potential, the maximum velocity should be inversely proportional to ramp

$$\text{width } qv_y B_z \leq qE_{\perp} \Rightarrow v_y \leq \frac{\Phi}{B_z L_R}.$$

Our initial baseline studies confirm this. We show below in Figure 3-5 a plot of maximum velocity as a function of ramp width, both given in normalized units – the Alfvén velocity v_A and the local ion inertial length $\lambda_{pi} = c / \omega_{pi}$. The test particles are protons $mass = 1$ and the shock Alfvén Mach number $M_A = 5$.

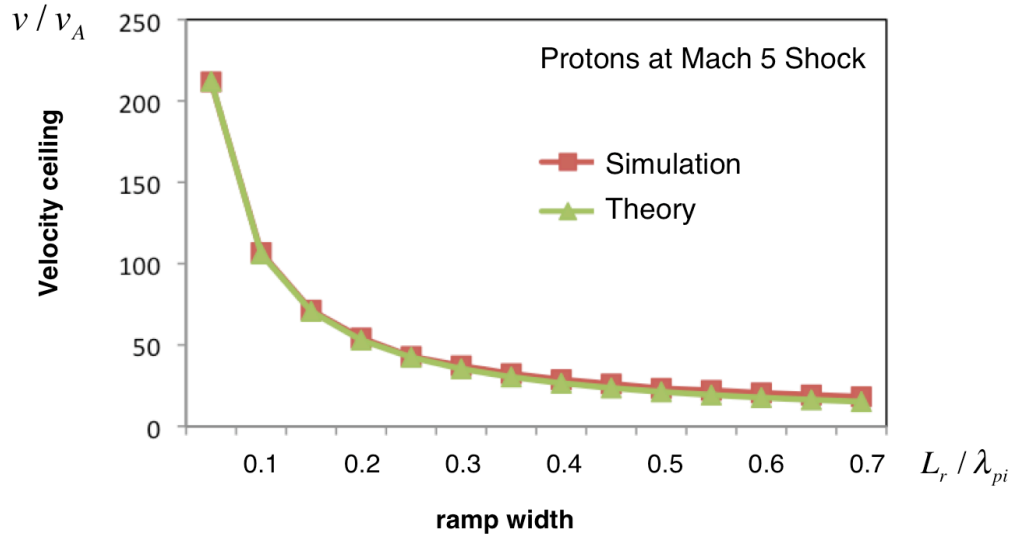


Figure 3-5 Maximum Surfing Velocity vs. Ramp Width

Ramp width L_r is normalized to an ion inertial length c / ω_{pi} and velocity to the Alfvén speed v_A . Data is shown for protons at Mach 5 shock.

Note that the tail of the velocity maximum plot diverges slightly from our theoretical prediction, and we offer two possible explanations. First, note that there is a minimum energy gain of approximately twice the piston energy associated with the first reflection. The first condition sets an absolute floor for all surfing ions. To that we must add any additional energy gained during transit through the ramp while the ion is retarded and therefore still under the influence of the convective electric field. We suggest that the velocities actually attained exceed our predictions simply due to the extra time taken for ions to traverse the wider ramps. Ions that reach the highest velocities via shock surfing begin with zero normal incident velocity at the shock front and we therefore expect them to remain at the leading edge until they reach their force balanced terminal velocity.

A trivial calculation can be made assuming constant normal acceleration through the ramp, and shows that $\tau = \sqrt{2L_r / a} \sim L_r^{1/2}$. The dependence $\Delta v \sim L_r^{1/2}$ follows from the constant longitudinal ion acceleration due to the convective electric field, and was verified by the simulation data shown in Figure 3-6.

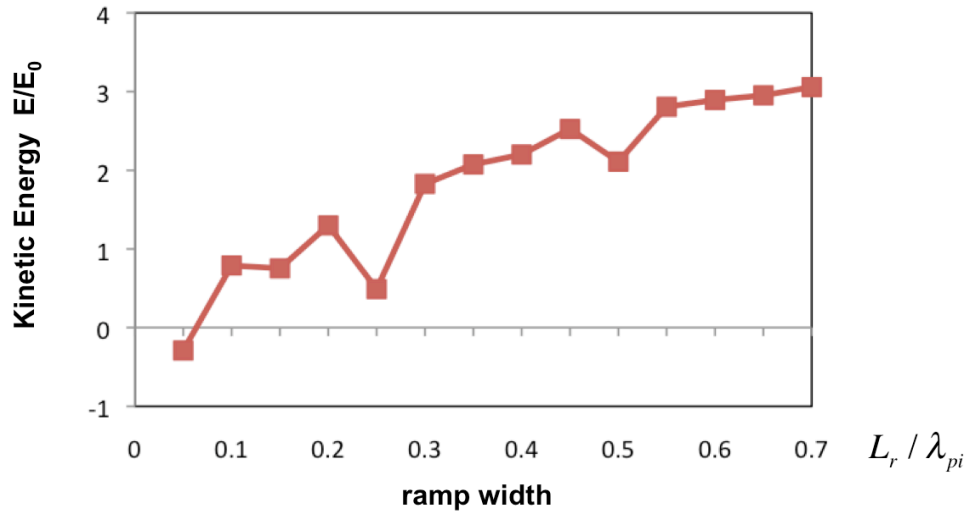


Figure 3-6 Deviation from theoretical limit vs. ramp width

The function shows $L_r^{1/2}$ dependence for protons at Mach 5

We present the surfing velocity ramp dependence for protons once again, but plotted with the aggregate shock drift velocity ceiling for comparison in Figure 3-7 below. While the ramp transit time argument holds, it is clear that the surfing ions at the widest ramp shown are reaching the shock drift ceiling asymptotically. Hence we infer that there does in fact exist a velocity ‘floor’ for all pickup ions.

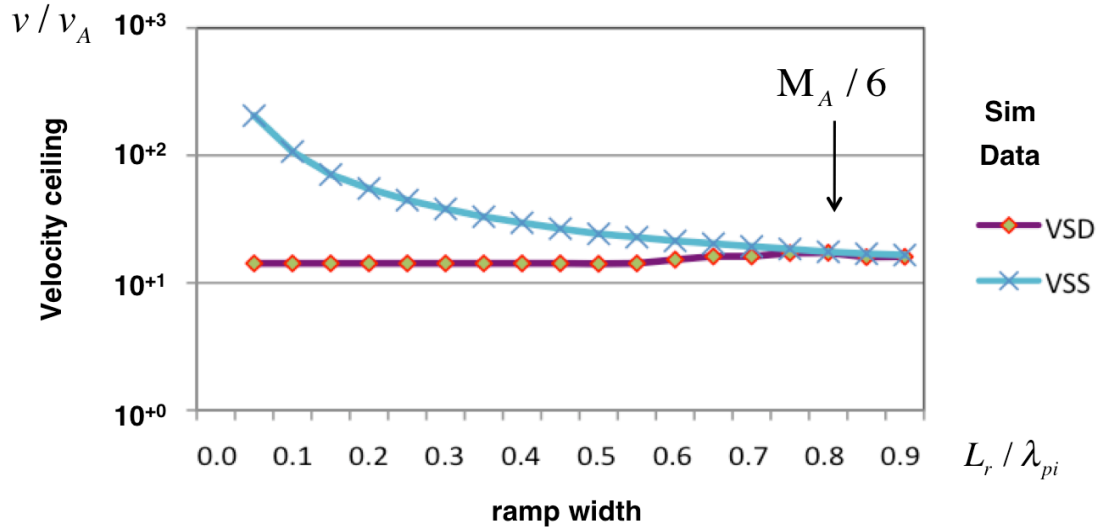


Figure 3-7 Comparison of SS and SD velocity ceilings

Velocity normalized to the Alfvén speed is shown as a function of ramp width for protons at Mach 5 shock. Ramp width is normalized to the ion inertial length. The surfing mechanism cannot promote ions above the SD ceiling as the ramp width approaches an ion inertial length.

We note in passing, that our force-balance equation in the ramp was a static approximation, giving only the condition when a surfing ion could no longer be restrained at the leading edge of the ramp. It was then assumed that its transit time through the shock would be rapid enough to be ignored. However, this can be converted into a dynamic problem, and the shock surfing ‘floor’ at wide ramps may be calculated precisely by solving the coupled differential equations $m\ddot{x} = q\dot{y}B(x(t)) - qE_{\perp}$ and

$$m\ddot{y} = qE_{\parallel} \text{ where } B(x(t)) = B_1 \left[1 + \frac{x(t)}{L_r} (R_c - 1) \right].$$

Mass dependence

We note that the maximum velocity function is independent of mass, and therefore expect our simulation results to validate this. We plot a family of curves representing four different masses in Figure 3-8 below. Aside from small discrepancies due to runtime limitations at the thinnest ramps, the mass independence is demonstrated by the superimposition of these curves.

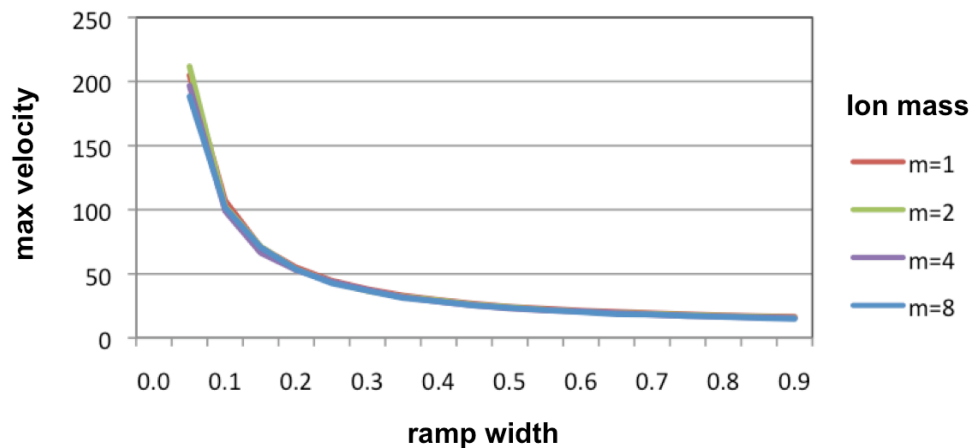


Figure 3-8 Surfing velocity ceiling as a function of ion mass

The maximum velocity attainable by four masses normalized to m_p are plotted against ramp width at a perpendicular Mach 5 shock. The curves are identical, but heavier ions take longer to reach their velocity maxima.

We also compare the velocity maxima for the shock drift mechanisms we have called single-bounce or ‘looping’ and the ‘crossing’ ions in Figure 3-9. Although the looping ions show mass-dependent ceiling attenuation, this dependence is negligible for crossing ions. One apparent feature is the small velocity increase for both species at the widest ramps. One possible explanation for this is the longer ramp transit time $\Delta\tau$

associated with wider ramps. Recall that no account of the shock-normal electric field was taken in the shock drift derivation. Clearly the retarding force must influence the trajectory and transit time, which in turn produce an additional drift along the convective electric field, independent of the grad-B drift. In this sense the two mechanisms SD and SS are not strictly disjoint. We did not explore this further, but note that the 'Crossing' velocity maximum is insensitive to variation in both mass and ramp width.

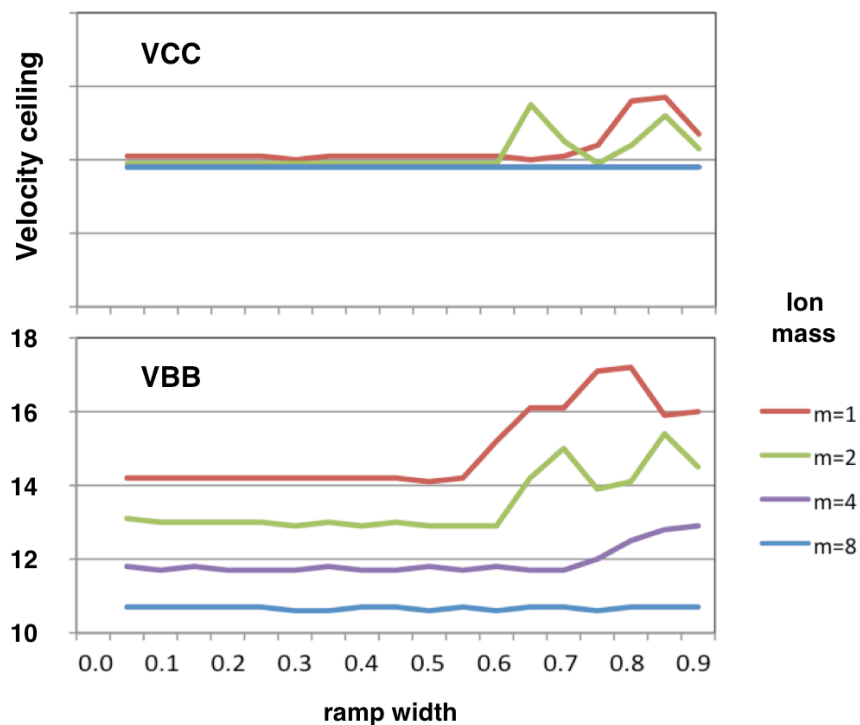


Figure 3-9 Shock Drift Velocity Maxima

The top panel represents 'crossing' ions and the bottom 'single bounce'.

The four curves in each panel represent four ion masses.

Accelerated Fluxes

We now examine the total fluxes of ions accelerated by each of the mechanisms, in terms of the fraction of the total pickup ion population they represent. Again, we are interested in mass and ramp width dependence as a smooth shock baseline study, not only to establish functionality in the best possible case, but also for comparison when we apply perturbations at the front.

We begin by comparing the fluxes of protons accelerated by the two different mechanisms; shock surfing (FSS), aggregate shock drift (FSD) and its individual components, 'single bounce' (FBB) and 'crossing' (FCC) as shown in the composite log-linear plot of Figure 3-10. The shock drift mechanism is only slightly affected by the ramp width, while the surfing ion flux diminishes with increasing width as predicted from the geometric construction shown in Chapter 2.

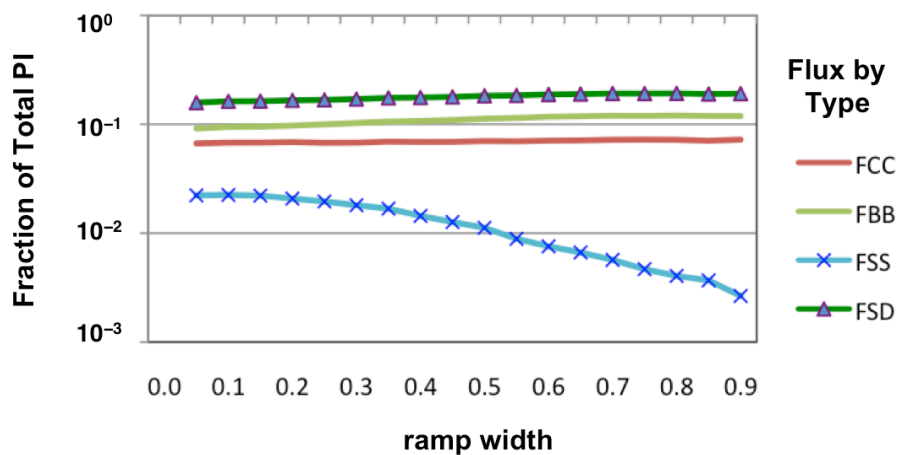


Figure 3-10 Total SS and SD Fluxes

The aggregate Shock Drift mechanism and its two constituents show little sensitivity to variation in ramp width, while surfing is severely suppressed.

Simulation is for protons at Mach 5 shock.

We next examine the mass dependence of the total shock surfing population in Figure 3-11. At the narrowest ramp width, the total flux mass dependence is as we predicted in Chapter 2. We note that the curves converge as the ramp widens as shown in Figure 3-12.

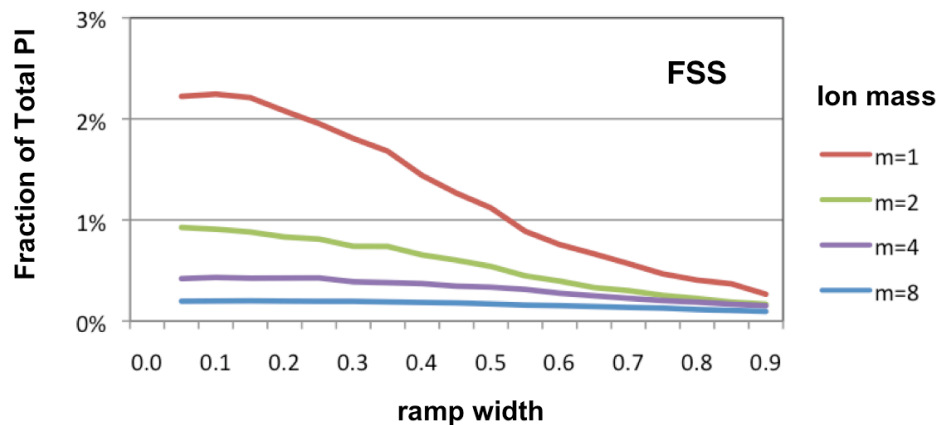


Figure 3-11 Shock Surfing Flux Mass Dependence

The flux is given as a percentage of the total incident PI flux.

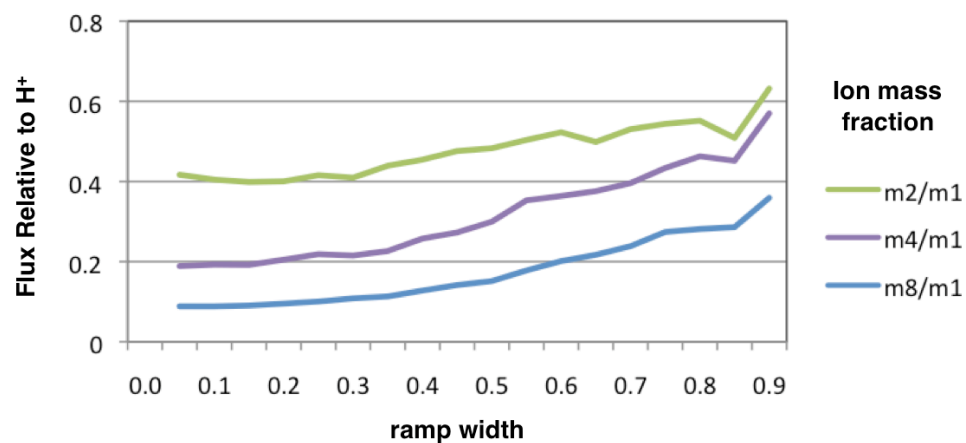


Figure 3-12 Convergence of fluxes at wide ramps

The fluxes of heavier ions are given as decimal fractions of the proton flux at the respective ramp width as shown in the previous figure.

We see in Figure 3-13 that the shock drift 'crossing' component (FCC) is relatively insensitive to ramp width but the flux of the 'single bounce' or 'looping' component (FBB) increases with ramp width, most noticeably for protons. This should not be surprising since the wider ramp offers more opportunity for ions to thread the ramp in this way. The mass dependence of the two sub-species is also easy to understand. The ramp threading of the 'looping' component is a function of reflection coefficient and, since this shrinks with increasing mass, so should the flux. The converse is true for the 'crossing' component as it is composed of transmitted ions, and that region grows with increasing mass.

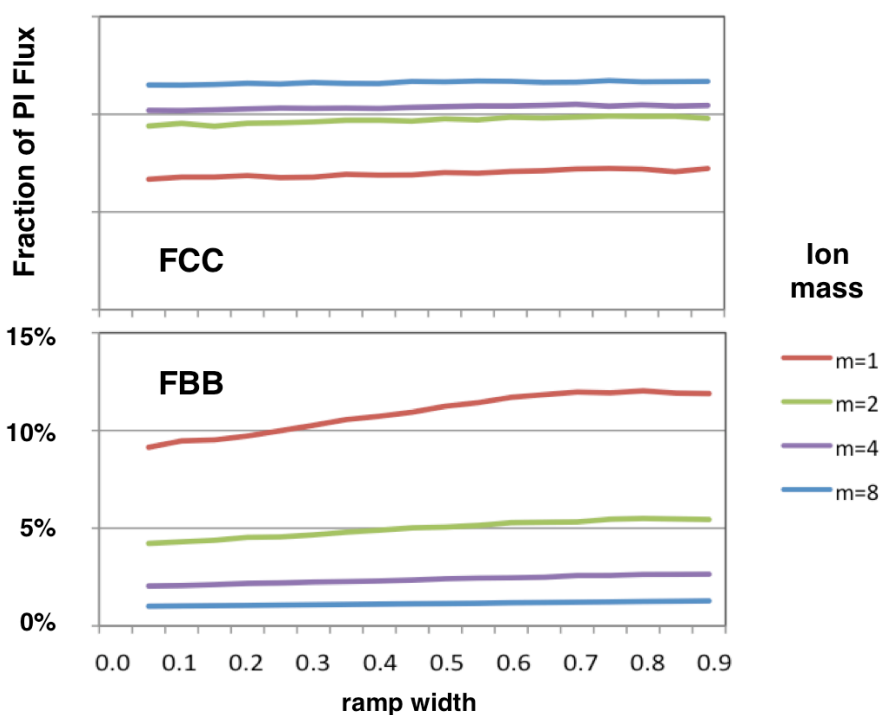


Figure 3-13 Shock Drift Component Fluxes

Curves representing four masses are plotted against ramp width. Single-bounce flux (FBB) is shown left, and crossing ion flux (FCC) on the right.

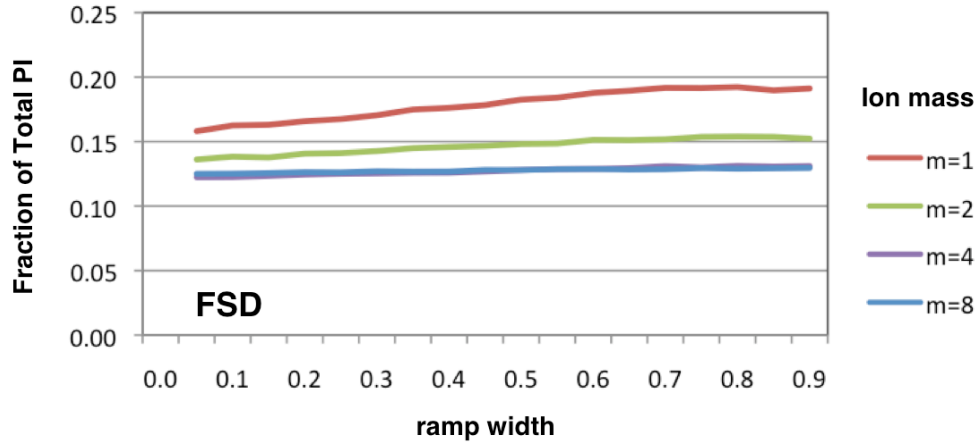


Figure 3-14 Aggregate Shock Drift Mechanism

Summed SD components ‘single-bounce’ and ‘crossing’ are plotted against ramp width. Weak mass sensitivity is evident from four curves.

The ordinate is given as a decimal fraction of the total incident PI flux.

As seen above in Figure 3-14, the aggregate shock drift species are relatively invariant compared to the rapid decline of the shock surfing flux. However, even at narrow ramps, neither the total surfing flux nor the velocity ceiling is of any consequence without some measure of the tail population above this base. We shall explore this requirement next.

Shock Surfing Efficiency

As a first attempt to quantify the effectiveness or ‘efficiency’ of the surfing mechanism, we compute the fraction of surfing ions which exceed the aggregate velocity ceiling of the other acceleration mechanisms, collectively considered as ‘shock drift’. For protons at encountering a Mach 5 shock, this is very nearly $\hat{v}_{SD} \sim 3u_1 = 15v_A$ but, as we shall see, it is not strictly invariant. We therefore exhibit first a plot of efficiency defined as the flux above the fixed velocity $3u_1$ in Figure 3-15 below that clearly shows a strong

dependence on ramp width. Indeed, when $L_r \sim 0.7$, the population above this marker has dropped by more than an order of magnitude and vanishes entirely at $L_r = 1$. In the absence of a specific and well-defined SDA threshold, there is no way to unequivocally establish the effectiveness of shock surfing as a pre-injection mechanism, but the criteria we have just established illustrate the worst possible case, namely, where no shock surfing ions can exceed the shock drift ceiling. Since the entire surfing population is scant compared with that of shock drift, we can say with absolute certainty that ramp widths above those shown must render the shock surfing mechanism completely untenable.

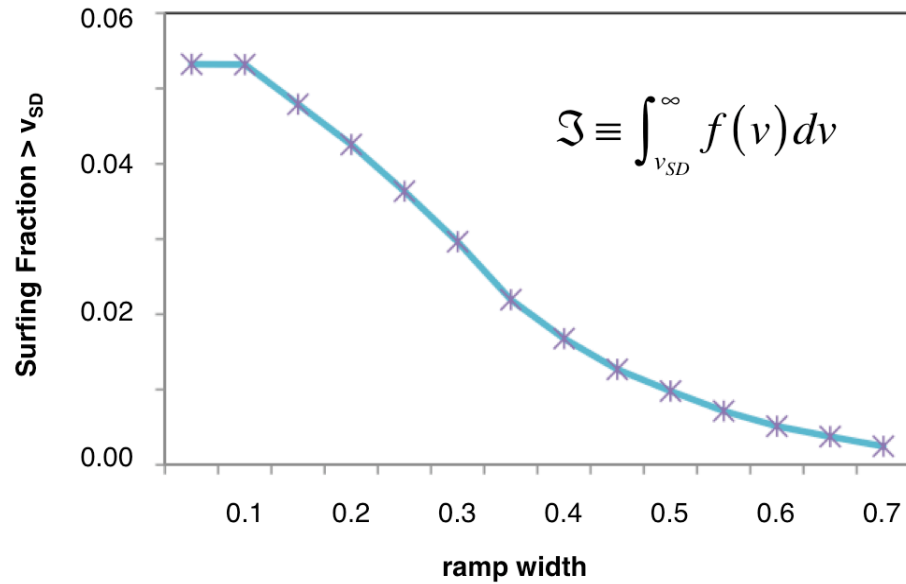


Figure 3-15 Surfing 'Efficiency'

The fraction of total PI flux promoted above 3 times the ram velocity is plotted against ramp width for protons at Mach 5 perpendicular shock.

Next, we illustrate the marked mass dependence on shock surfing efficiency in Figure 3-16 below, given as a log plot of the fraction of incident PI flux which shock surfing promotes ions above the minimum threshold marker, $v_{thr} = 3u_1$. The flux for ions $m = 8$ is more than an order of magnitude below that of protons. Again, the ramp dependence is even stronger, the flux dropping two orders of magnitude as the ramp width increases from $L_r = 0.05$, roughly two electron inertial lengths, to the maximum ramp width shown $L_r = 0.90$ which is just under an ion inertial length. Obviously, once the ramp width exceeds unity, the surfing velocity ceiling lies below that of shock drift, and the efficiency vanishes entirely.

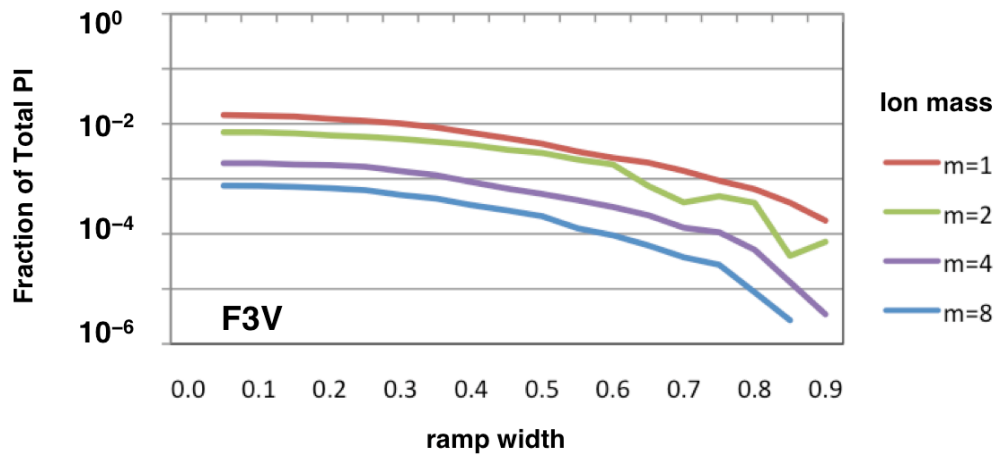


Figure 3-16 Mass comparison of SS Efficiency

Flux of Surfing Ions promoted above $3u_1$ plotted against ramp width at Mach 5. More than one order of magnitude separates species at the narrowest ramp $L_r = 0.05$. The flux for all ion species drops by two orders of magnitude as the ramp width approaches an ion inertial length.

Distribution Function

As shown below in Figure 3-17, the distribution function for each category of accelerated particle was our primary analytical tool. Bear in mind that the thrust of this investigation is to explore the viability of shock surfing as a pre-injection mechanism. As such, the velocity ceiling may rule out a particular mechanism, but the flux of surfing particles with velocity greater than that of other mechanisms becomes a valuable criterion in determining its efficiency.

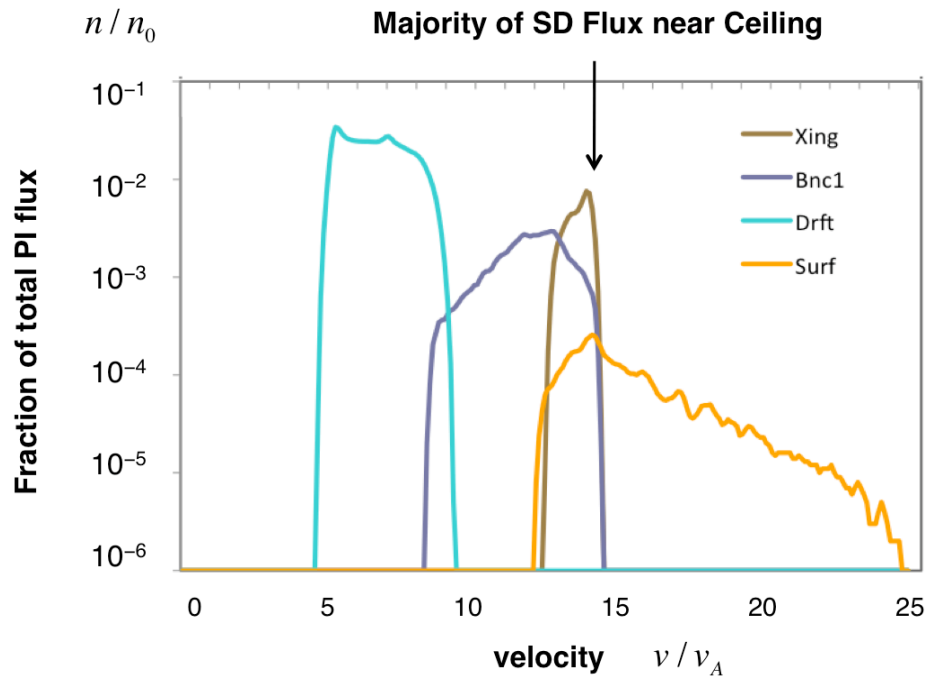


Figure 3-17 Transmitted PI Distribution Functions

Curves represent distinct subsets of ions accelerated by mechanisms as defined by 'tagging' particles. Unlike DSA, Shock Surfing generates an exponential tail rather than a power law, as seen in this log-linear plot.

Aside from the surfing ions, which exhibited the predicted ramp width dependence for maximum velocity, all other mechanisms had hard velocity ceilings for each given set of input parameters, and these were essentially insensitive to ramp width. It is important to note that, although the Shock Drift ceiling appears to indicate little acceleration – the PI incident velocity maximum in the shock frame being twice the Mach number – the distribution function for these ions is pushed upwards. In other words, acceleration is taking place, but not for the very fastest ions in the shell. It may be tempting to assume that the fastest shock-incident ions spend too little time in the ramp for effective displacement. Recall that $\Delta K = qE_y \Delta y$ and $\Delta y = v_{\nabla B} \Delta t$ where the ramp crossing time $\Delta t = L_r / v_x$ is inversely proportional to the normal velocity. A more likely explanation is that guiding-center displacement, and therefore energy gain, is due to threading the front. All trajectories become kinked there, but those with $v_y < 0$ return while those with $v_y > 0$ do not. Examination of energy gain within the initial pickup ion shell shows that this is so, and we prove this rigorously in the next Chapter.

Mechanism Source Regions

In Figure 3-18 below, we show the kinetic energy of all pickup ions after passing through the front. Each panel represents a different ramp width, broadening from left to right. Energy is plotted as intensity upon a velocity space projection of the pickup ion shell incident upon the shock ramp. As expected, the maximum energy decreases as ramp width increases. Moreover, there appear to be distinct regions within the shell, suggesting unique mechanisms. Note the following features: 1) rotation of the region terminator in the left hemisphere with increasing ramp width, 2) the small region of large energy gain near $v_{x0} = v_{y0} = 0$ which vanishes with increasing ramp width and 3) that,

with exception of that small ‘slice’, energy gain in the active regions increases with increasing normal incident velocity v_{x0} .

We may also compare some of the features with the predictive plots and graphical constructions from Chapter 2. Note that nearly vertical hard terminator in the left-hand panel, which represents the boundary between reflected and transmitted ions. By comparison, we observe that the boundary coincides with the inverted parabola from the finite-ramp equation, and that the parabola broadens with increasing ramp width, lowering and broadening the terminator.

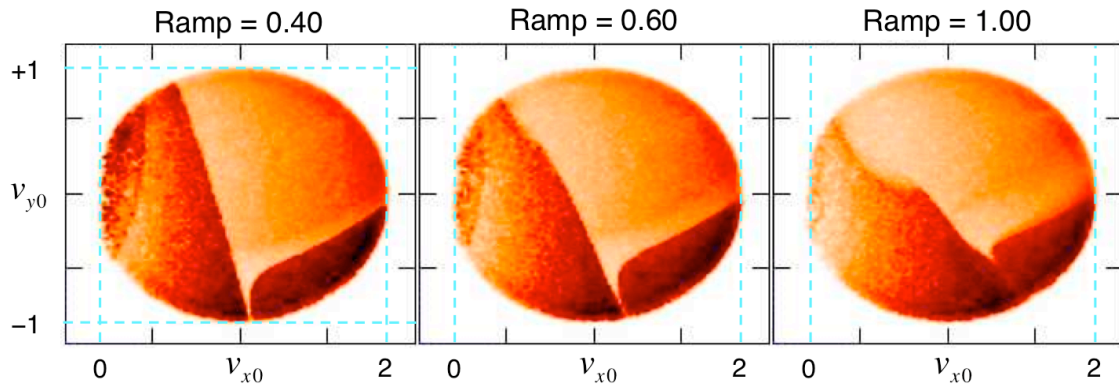


Figure 3-18 Ion Energy Gain

Transmitted ion energy is superimposed on incident velocity distribution. Velocity is normalized to upstream solar wind speed u_1 . Maximum energy is normalized to the piston energy $m_p u_1^2 / 2$. Energy maxima from left to right, $E_{\max} = \{16.0, 9.29, 8.68\}$. See text for further discussion.

The following set of plots shown in Figure 3-19 consists of superpositions assembled from distinct acceleration mechanisms, as discriminated by our tagging scheme. Here we clearly see that these mechanisms operate on disjoint subsets of the

counter-clockwise rotation of the highest energy region terminator of the shock drift component with increasing ramp width. The widest ramp shown is $L_r = 1.00$, and its angle of rotation $\Delta\theta_g \approx 10^\circ$ is easily explained by taking into account the ion gyrophase displacement during its transit time $\Delta\tau$ through the ramp. All ions shown are of the same species (in this case protons) and possess the same ion cyclotron frequency

$$\Omega_{ci} = \frac{eB}{mc}, \text{ equivalent to their gyrophase velocity. We can make a rough calculation by}$$

using the mean ion normal velocity and mean magnetic field strength through the ramp.

The fastest ions will maintain nearly their incident velocity $\langle v \rangle \sim 2u_1$ while the mean

magnetic field $\langle B \rangle = \frac{1}{2}(B_1 + B_2) = \frac{B_1}{2}(1 + R_C)$. The gyrophase displacement is just

$\Delta\theta_g = \Omega_{ci}\Delta\tau$ and the transit time is simply the ramp width divided by the velocity, so we

$$\text{have } \Delta\theta_g = \Omega_{ci}\Delta\tau = \frac{\langle B \rangle L_r}{m\langle v \rangle} = \frac{B_1(1 + R_C)L_r}{m(2u_1)}. \text{ Applying the normalized values for protons}$$

at a Mach 5 shock, we have $\Delta\theta_g = \frac{(1 + R_C)L_r}{(2M_A)} = 0.18L_r \text{ rad}$, which at the unity ramp

width gives $\Delta\theta_g = 10.4^\circ$. Note that since the shell rotates clockwise in time, the counter-

rotation is consistent with the original phase space shell being measured at the leading

edge of the ramp, in other words, retarded in time.

The Flux Problem

The contribution of Shock Surfing is clearly inconsequential when the surfing velocity ceiling lies at or below the drift maximum. However, simulations show that the total Pickup Ion flux accelerated by Shock Surfing is in any case at best two orders of

magnitude below that due to Shock Drift. So, even when the surfing ceiling is substantially above that maximum, two issues remain: (1) whether any part of the tail lies above the DSA injection threshold at all and, (2) if so, whether the integrated flux above that threshold is sufficient to account for the observed anomalous cosmic ray flux. One of the big problems that remain in making such an analysis possible is the absence of reliable Pickup Ion flux observations in the regions of interest.

On the other hand, we can make reasonable estimates on conditions that would render shock surfing completely untenable. Combining our surfing velocity ceiling equation with that for the potential step, and expressing the result in dimensionless (normalized) units, we have

$$v_y \leq \frac{M_A^2}{2L_r} (1 - R_C^{-2})$$

Since the velocity ceiling for the other species is always less than three times the Mach number and since the factor $(1 - R_C^{-2})$ is of order unity or less, we can safely assume that in order for shock surfing to function at all, we must in all cases satisfy

$$L_r < \frac{1}{6} M_A.$$

Note that this is an *absolute* ceiling and, as shown above, the surfing efficiency i.e., flux above the drift ceiling drops dramatically with increasing ramp width. Hence it is reasonable to demand that, to achieve any substantial surfing flux, $L_r \ll M_A$ must be satisfied by at least an order of magnitude, and preferably two.

Curvature Convection Limit

Our simulations are strictly one-dimensional, representing a plane front of infinite extent. But planetary bow shocks are concave, which places spatial and temporal constraints on acceleration processes due to convection along the front by the streaming solar wind. For curved magnetospheres, the region of viable shock surfing is confined to an arc of at most one radian, imposed by the geometry, which limits the span of quasi-perpendicularity. This in turn puts an upper bound on the path over which surfing acceleration can take place.

Assume the front ‘standoff’ radius and therefore the arc length are approximately $R_s \approx 16R_E \approx 10^8 m$. Taking generous upstream solar wind parameters $B \approx 5 \times 10^{-9} T$ and $v_{sw} = 5 \times 10^5 m \cdot s^{-1}$, we may compute a velocity limit for protons,

$$v_{\max} = \sqrt{2aR} = \sqrt{2 \left(\frac{q}{m} B v_{sw} \right) R}, \text{ or } v_{\max} \approx 6.9 \times 10^6 m \cdot s^{-1} \approx 14 v_{sw} \text{ which, in our}$$

normalized units, is simply 14 times the Alfvén Mach number. Neglecting the thermal component, we calculate the energy of the ambient solar wind protons to be approximately 1.3 keV, which gives an upper bound of about 250 keV.

Although surfing ions accelerate perpendicular to the solar wind flow, they are also driven by convection along the shock surface, which imposes yet another constraint on the time of flight of roughly the same magnitude. Even though these limits may still be consistent with “field-aligned beam” energies, the ceiling clearly disqualifies surfing as the sole mechanism for producing anomalous cosmic ray candidates. Incidentally, although heavier ions have many other hurdles to overcome, the geometric constraint is not among them. Note that geometric limit $v_{\max} \sim m^{-1/2}$ so that, since surfing ions all

begin essentially at rest, the kinetic energy $\varepsilon \sim mv^2 \sim m(m^{-1/2})^2 = 1$ that is to say, there is no mass dependence on the geometric energy limit.

Geometric Flank Limit

In addition to other geometric considerations, a further limiting factor exists along the flanks of a shock. Assuming that the fluid flow is very nearly normal at the shock nose, it cannot be along the flanks, i.e., there must be some component of the front-incident velocity parallel to the surface. *In situ* observational data at the Jovian bow shock confirm this fact, as we show and discuss in the chapter on the outer gas giant planets. The salient point here is that, in modeling a plane one-dimensional front – regardless of the Parker angle – we obtain results which are extremely conservative. In other words, the performance of shock surfing so obtained is the best possible. If the mechanism fails under the assumptions of this model, it must fail under all conditions.

Shock reformation

One last point to consider in light of the discussion of variability in the previous section is that of shock oscillation and reformation. Again, theoretical, numerical and observation sources all suggest this process takes place on the scale of an ion cyclotron period. Even if the shock does not overturn, but merely ‘breathes’ – its ramp narrowing and widening – the oscillation in ramp width must seriously impair or even terminate the surfing mechanism’s efficiency. A trivial calculation gives the maximum surfing velocity limited by such ramp modulation

$$v_{\max} = a\tau = \frac{qE_{\parallel}}{m} \tau = u_1 \left(\frac{qB_1}{mc} \right) \frac{2\pi}{\omega_{ci}} = 2\pi u_1$$

Applying a conservative estimate that surfing is viable during only half the cycle gives $v_{\max} \sim 3u_1$, which is roughly the same as the Shock Drift ceiling. Note that this constraint is particularly deleterious to heavy ions since the time required for acceleration scales inversely with mass.

Summary and Conclusions

We have presented an exhaustive exposition of baseline studies made at smooth quasi-perpendicular shocks, including models to explain the observed parametric dependencies of SS velocity maxima and fluxes. For the interest of investigators familiar with the formalism of Lee et al (1996) we refer the reader to Appendix A where we show the effect of more realistic ramp widths on the fundamental variables of their formalism, e.g., the adiabatic invariant, which deviate substantially from those theoretically required. Since the range of ramp widths we have explored is extremely narrow compared with those of shocks observed in the heliosphere, the fact that these variables deviate so rapidly from those predicted by theory indicate that such a formalism cannot be expected to reliably predict the behavior of the surfing mechanism in the more general context of the manifold shocks actually found in nature.

We have seen that at finite ramp shocks, a maximum velocity exists for surfing ions, that this ceiling is inversely proportional to the width $v_{\max} \sim L_r^{-1}$ and drops to the shock drift velocity ceiling $v_{\max} \rightarrow v_{SD}$ when the ramp widens to the ion inertial length $L_r \rightarrow \lambda_{pi}$. The total shock surfing flux, as a fraction of the entire incident pickup ion population, drops approximately an order of magnitude from the thinnest physically possible ramp to the widest $L_r : \lambda_{pe} \rightarrow \lambda_{pi}$. The surfing flux is also mass dependent and approximately inversely proportional. Conversely, both the velocity ceiling and total flux

are essentially invariant for the shock drift mechanism regardless of variation in ramp width or mass.

The efficiency as defined by the flux of surfing protons above the shock drift ceiling is approximately 1% of the total incident PI flux, and drops by two orders of magnitude when the ramp widens to an ion inertial length. The oxygen flux begins at 0.1% and drops by three orders of magnitude at the widest ramp. Hence, we may definitely conclude that shock surfing even at the ideal shock strongly favors light ions and very narrow shock ramps to be effective at all. We will now see what happens when the shock front is not ideal.

Chapter 4 The Perturbed Shock

The baseline studies presented in the previous chapter provide results for idealized plane-front shocks of infinite extent. These studies establish a set of data for smooth shocks that demonstrate the parametric dependence of both maximum velocity and ‘efficiency’ of Shock Surfing (SS) as functions of shock Mach number, ion mass, ramp width and magnetic field angle. In order to better evaluate surfing effectiveness, we compared the effects that the same parametric variation has on the Shock Drift (SD) mechanism as well. Finally, we made a cursory study of deviations from the idealized shock, with estimates on limitations imposed by curved fronts such as those found at planetary bow shocks, and temporal constraints imposed by periodic shock reformation.

We now inquire whether the SS mechanism can remain robust when the front is perturbed. These investigations are motivated by observations made over several decades which have shown that the Earth’s Bow Shock is quite active. Magnetosonic, Shear Alfvén, Lower Hybrid, Ion Cyclotron and Whistler waves have all been detected, as well as populations of gyrating reflected ions and field-aligned beams. While Small Large Amplitude Magnetic Structures (SLAMS) are commonly found ahead of quasi-parallel regions, their effects may also be seen along the quasi-perpendicular flanks of inner planets where the Parker angle is steeper, as some fraction of energetic ion populations observed there may in fact originate at SLAMS and be swept downstream.

Until recently, such detailed observations and analyses were necessarily restricted to the Earth’s magnetosphere. However, analyses of *in situ* data gathered by missions to the outer heliosphere have begun to appear in the literature, giving new

insight into the behavior of shocks at the gas giants as well as the heliospheric termination shock.

Plasma Waves

Since shock surfing is a coherent process, we first explore the effects of simple plasma waves on particle acceleration, and refer the reader to Appendix B for a more complete exposition and analysis of the results we summarize here. Shear Alfvén, Lower Hybrid and Magnetosonic waves in turn were imposed upon the background fields employed to drive proton test particle trajectories at a perfectly perpendicular Mach 5 shock. A battery of simulations was performed while varying the amplitude, wavenumber and polarization appropriate for each wave type. We chose an extremely thin ramp width at $L_r = 0.05 \lambda_{pi}$ or about two electron inertial lengths so as to observe the most profound effects since wider ramps, as we have shown, suppress both the velocity ceiling and fluxes of surfing ions. The parameters explored ranged from zero to the limits of physically reasonable field amplitudes.

Our simulations show that Lower Hybrid waves propagating parallel to the solar wind can suppress the velocity ceiling and surfing efficiency by up to 70%, but only for extremely large field amplitudes, with energy density more than an order of magnitude greater than that of the ambient solar wind magnetic field. Waves traveling along the shock front produced negligible effects unless the wavelength exceeded scales comparable to planetary bow shock structures. Alfvén waves produced almost no effect whatsoever on the velocity maxima of surfing ions. Some influence was seen on the flux efficiency, but this varied according to wavelength and was apparent only at the highest amplitudes. Magnetosonic waves propagating in the fluid direction normal to the front

suppressed the flux and the velocity by a factor of two – again only at high field amplitudes – while waves traveling along the shock surface had smaller effect.

Whether mechanisms exist to drive such waves at the amplitudes and wavelengths required for suppressing shock surfing is beyond the scope of this investigation. However, it should be noted that although strong turbulent energy spectra have been observed, such waves are incoherent, stochastic ensembles, and therefore more likely to generate velocity diffusion. Moreover, the attenuation of the surfing velocity ceiling and tail populations were observable only for the very steepest shock ramps, since velocity and flux suppression above $L_r = 0.1 \lambda_{pi}$ exceed that attributable to plasma waves. This is not to say that waves produce no effect at all on pickup ion trajectories. We investigate wave-induced velocity diffusion in Chapter 5, but show that time scales required for significant diffusion to evolve coupled with the long wavelengths required, are inconsistent with shocks inside the heliosphere.

The Corrugated Front

We shall now demonstrate that physical ripples in the structure of the shock front profoundly suppress shock surfing. We present quantitative results and discuss in detail the mechanisms responsible for terminating the acceleration process. Having already discussed the effects of spatial and temporal constraints at shocks of finite extent, we now assume that the deformations of interest are on length scales small compared with the curvature of the overall shock structure, that is, a corrugation or rippling along the surface of the front.

Such corrugations have been observed at the Bow Shock by the Cluster mission, at outer planet shocks by Cassini and Voyager I missions, and at the Termination Shock

by the Voyager II mission. Ripples in shock surfaces have also appeared in recent multi-dimensional numerical simulations, e.g., Whalen and Norman (2008). We anticipate that this type of perturbation will produce greater effects on SS than waves alone, since the shock-normal electric field – at least inside ramps thin enough for effective surfing at all – is orders of magnitude greater than other ambient fields, and changes in its direction and geometry must necessarily be felt more strongly by surfing ions.

We begin with a simplified model of fluid entering, flowing through, and exiting the ramp normal to the surface as shown in Figure 4-1 below. Since the magnetic field in the z -direction is unperturbed, it follows that the convective electric field must always be tangent to the front. Since the electrostatic potential step at the front is fixed, we further assume that the perpendicular fluid path through the ramp is also fixed, producing a uniform retarding electric field inside the ramp and normal to the front. We assume that the fluid flow is uniform far from the ramp itself.

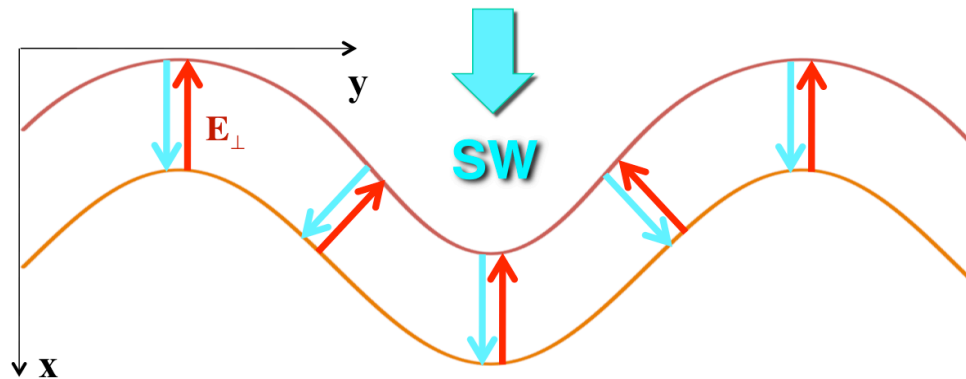


Figure 4-1 Velocity field inside the ramp.

The solar wind (blue) flows from top (upstream) to bottom (downstream) in the x -direction. The velocity between the two corrugated boundaries (inside the ramp) is normal to the surface and the fluid paths are everywhere equal. The retarding electric field inside the ramp is antiparallel to the flow throughout.

In order to realize this model, we begin with a single static sinusoidal perturbation for the potential flow and include exponential decay both upstream and down to effect a smooth transition from ambient to perturbed values

$$\psi = -ux - a \frac{u}{k} e^{-kx'} \cos ky \quad \text{where } x' = \begin{cases} x & x > 0 \\ -x - L_r & x < -L_r \\ a & -L_r < x < 0 \end{cases}$$

Here a is the perturbation amplitude, k is the wavenumber, and x' is the distance from the base of the ramp (upstream), or distance from the top of the ramp (downstream). Both a and x' are normalized to the ion inertial length $\lambda_{pi} = c / \omega_{pi}$, while

k is normalized to the inverse length $k_{pi} = 2\pi\lambda_{pi}^{-1}$. Since $V = -\nabla\psi$ and $\mathbf{E} = -\frac{\mathbf{u}}{c} \times \mathbf{B}$, we

obtain the convective electric field for the region

$$\mathbf{E}_C = \tilde{E} \left(a e^{-kx'} \sin ky \right) \mathbf{i} + \tilde{E} \left(1 + a e^{-kx'} \cos ky \right) \mathbf{j}$$

where $\tilde{E} = -\frac{u_1}{c} B_1$ is the unperturbed upstream convective electric field. The chosen

function satisfies the requirement of potential flow, namely, that the Laplacian vanishes

$\nabla^2\psi = 0$ ensuring that $\nabla \times \mathbf{E} = 0$.

Inside the ramp, the retarding electric field $E_\perp$ is superimposed over this, and is rotated so as to be normal to the front. The perturbation is localized at the front and decays exponentially with distance both upstream and downstream, producing the static “corrugated” structure shown schematically in Figure 4-2. The magnetic field in the $\mathbf{z}$ -

direction (out of the page) is unperturbed, but a test run using $B_z = B_1 \left(1 + ae^{-kx'} \cos ky\right)$ to account for possible compression produced no noticeable difference in our results.

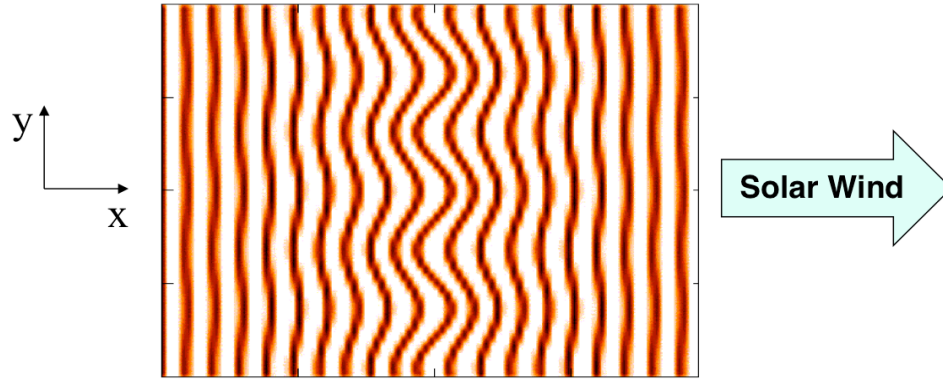


Figure 4-2 Convective Electric Field Lines

The magnetic field is out of the page and the perturbed shock front lies between the two corrugated streamlines straddling the vertical center line.

As in earlier wave studies, our damping scheme for the imposed field strength was intended to restrict the perturbation to the immediate region of the ramp and so eliminate any contributions to the ion distribution functions by velocity diffusion. A schematic diagram of the ambient field profiles is shown in Figure 4-3. The solar wind plasma density and magnetic field strength are set unity upstream and increase linearly through the ramp to their downstream values, equivalent to the shock compression ratio R_C . The magnitude of the convective electric field throughout is uniform since density and fluid velocity are inversely proportional, i.e., $u_1 n_1 = u_2 n_2$ and the magnetic field is frozen into the plasma, hence $E_{C1} = -u_1 B_1 = -u_2 B_2 = E_{C2}$.

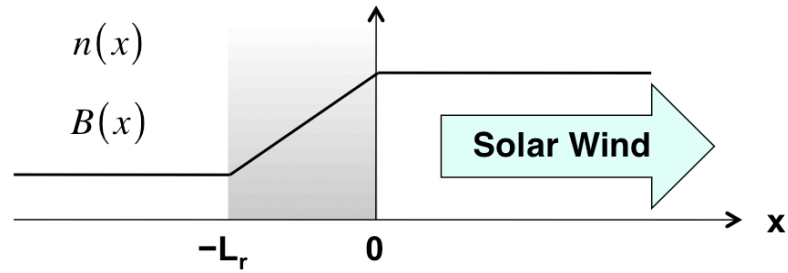


Figure 4-3 Density and Magnetic Field Profiles

Field strength increases from unity upstream (left) to the shock compression ratio downstream (right) through the shock ramp (shaded)

The damping function initially employed to rotate the convective electric field and magnetic vectors is illustrated in Figure 4-4 below.

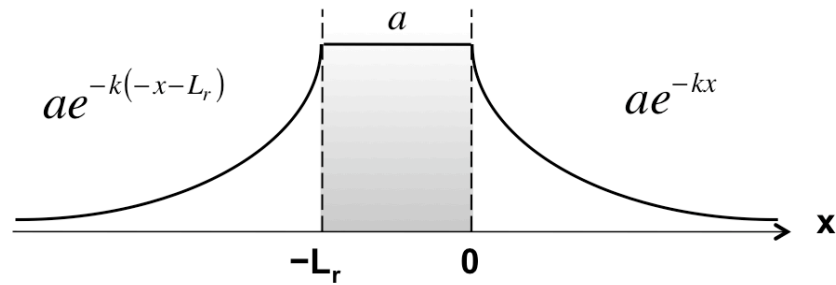


Figure 4-4 Perturbation Damping Function

Strength is uniform inside the ramp (indicated by shaded region) and decays exponentially with distance both upstream (right) and down (left)

The ambient magnetic field was varied linearly to its downstream value through the shock ramp as in our baseline studies. Several other schemes e.g., sine-squared profile to achieve continuous spatial derivatives were performed but produced negligible differences in the final data. After initial trials, an alternate scheme was tested in an attempt to provide more expedient computation. Since our goal was to extensively

explore parameter space which, given the sheer number of runs required, would have consumed more computational resources than were available, trials were made over several parameter variations wherein the unperturbed moduli for the electric fields $E_{\perp}$ and $E_{\parallel}$ were held constant, and the vectors simply rotated along the shock front according to the assumptions illustrated in Figure 4-1.

Further trials were performed where the downstream perturbation did not decay but remained constant at the front value. Finally, tests were made wherein the perturbation strength of the convective electric field was held fixed throughout the simulation box, and test particles injected within less than one ion inertial length ahead of the ramp to minimize velocity diffusion prior to injection.

Since it was found that the overall results differed between these various methods by less than a few percent, this last and simplest scheme was adopted. It should be noted that, in the final analyses, the effects of corrugation at the front produced attenuation in both maximum velocity and flux of surfing ions by orders of magnitude. Hence we consider the discrepancies in results between our initial idealized model and the simplified one to be of little consequence.

Insight and Motivation

Before proceeding to detailed analyses of the numerical data with respect to the performance of the surfing mechanism, we immediately noted from the exploratory simulations a definite suppression of velocity maxima with increasing wave amplitude and number, with substantial attenuation of the surfing flux above the Shock Drift velocity ceiling. We generated and examined trajectory plots at once in order to determine if

some obvious features might suggest a responsible mechanism, and illustrate two such plots below in Figure 4-5.

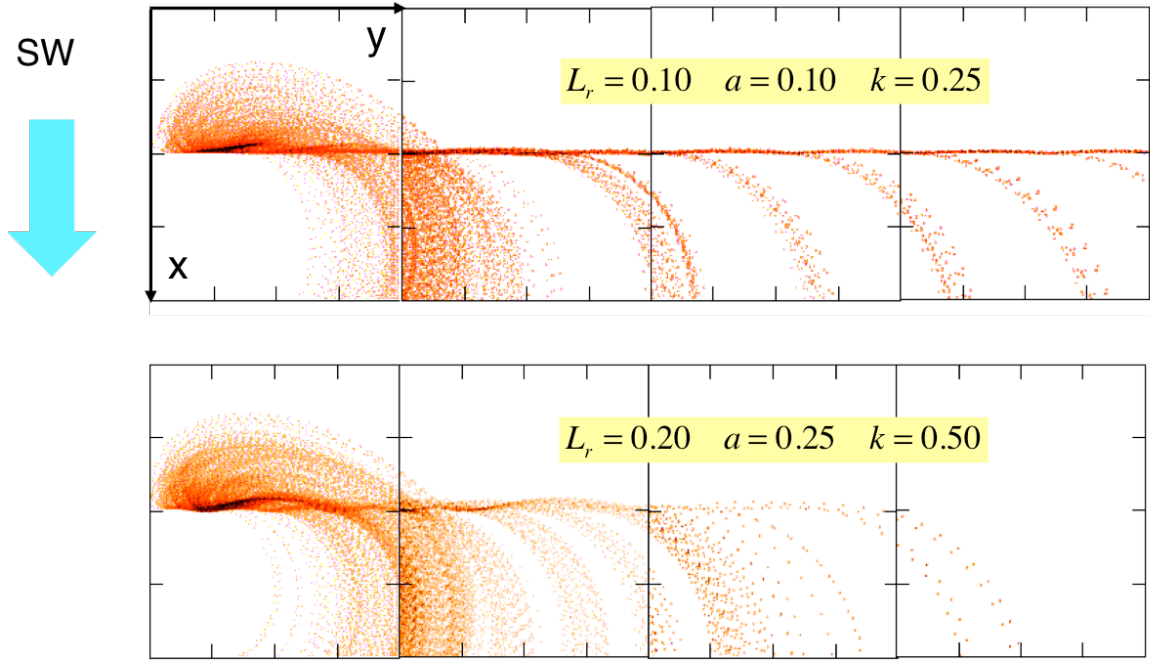


Figure 4-5 Proton trajectories at corrugated fronts

Two Mach 5 perpendicular shocks are shown, with ramp widths L_r and ripple amplitudes a given in terms of the ion inertial length λ_{pi} and the corrugation wavenumbers normalized via $2\pi / \lambda_{pi}$.

One identifiable feature is the periodicity in the flux of transmitted surfing ions that escape downstream. On closer inspection, we note that the length scale between the ‘escape channels’ is inversely proportional to the wavenumber k . We examined the source of this more closely by extending the resolution and expanding the scale of one such test as shown in Figure 4-6.

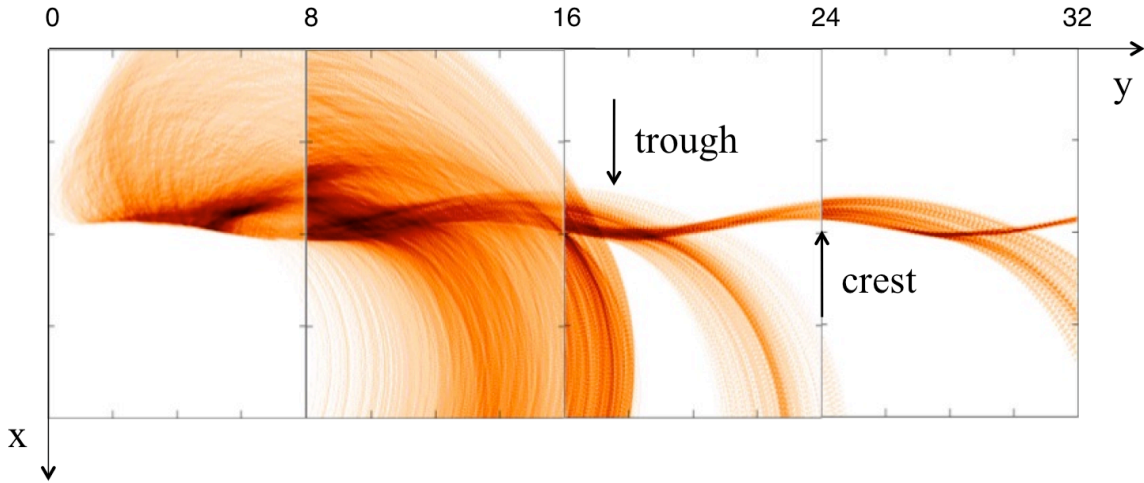


Figure 4-6 Ion trajectories at a corrugated front

The parameters for this shock are Mach number $M_A = 5$, field angle $\theta_B = 90^\circ$, ion mass $m_i = 1$, ramp width $L_r = 0.2$, perturbation amplitude $a = 0.5$ and perturbation wavenumber $k = 0.5$.

In these simulations, test particles were injected at a single point just ahead of the ramp base. This was done specifically to facilitate examination of the trajectory characteristics owing to the curvature of the front and, in addition, the depletion of the surfing ion population with distance along the curved path. This plot strongly suggests that the transmitted surfing ions escape the shock in streams with the same periodicity as the corrugation. Closer inspection shows that exit points correspond to locations of maximum front curvature.

Two possible factors immediately present themselves. At the 'crest' of the corrugation, the cyclotron radius of some ions exceeds the front curvature. These ions then temporarily escape the ramp and its normal electric field and, without that constraint, are free to assume upstream guiding center trajectories. In so doing, their

velocity vectors freely rotate until they re-enter the ramp – now with greater velocity normal to the surface.

In these cases of free-flight phase angle rotation, there is time for the ion to acquire sufficient normal velocity to breach the shock. This is exacerbated at the subsequent front encounter where the shock corrugation angle is pitched forward. Then the combined rotation of the ion velocity vector and the shock angle itself guarantee transmission through the front. Those that undergo the greatest excursion and associated velocity mixing acquire sufficient normal velocity to breach the shock. Indeed, the largest number of reflected particles by far does so only once before ‘puncturing’ the front. In addition, these excursions from the shock violate one of the fundamental tenets of shock surfing as defined in Lee et al (1996), namely, monotonically decreasing bounce normal amplitude.

While faster ions appear to continue accelerating without suffering this fate, it is the combination of parallel velocity and normal oscillation with its associated normal velocity component that ‘launches’ ions prematurely. However, even those that ideally encounter the shock with zero normal velocity must ultimately fall prey due to this force imbalance combined with normal oscillation induced by the corrugation itself.

Maximum curvature also occurs in the ‘troughs’ of the rippled surface. Force balance violation is particularly severe here since the curvature is concave rather than convex, opposite to the natural cyclotron orbits. The fastest ions closely hugging the surface can no longer follow the front, but are pushed through by the Lorentz force and their own momentum, and escape downstream.

Metaphorically, one might picture shock surfing as loosely analogous to skipping a stone on a pond. A smooth and polished stone must be chosen, of course, but a

shallow attack angle and a very smooth and calm surface are absolutely essential for multiple bounces to take place. If the pond surface is choppy, the stone must at some point encounter the leading edge of a wave and, with considerable normal velocity, breach the surface.

The trajectories show that ripples play an important role in disrupting the surfing process, with suggestions as to a possible source mechanism. Before exhibiting our numerical data, we first consider a physical model and derive a mathematical expression for the velocity ceiling at a corrugated front in order to both validate our simulation results and also produce a ‘litmus test’ for evaluating surfing viability from observation data we present in subsequent chapters.

We examined two physical models to explain the effect of corrugations on ion trajectories shown in the figures above. The derivations are somewhat lengthy, especially with regard to the justification of certain simplifying assumptions, and the complete detailed exposition is given in Appendix C.

Force Balance Model

The first model is based on force balance, and assumes that ions which ultimately reach the highest velocities have trajectories that remain just inside the leading edge of the ramp prior to escape downstream. The net force is comprised of the retarding electric field $F_E = -qE_{\perp}$, the Lorentz force $F_B = q\frac{v}{c}B$, and the centripetal force $F_{\perp} = \kappa m_i v^2$ required for an ion trajectory to follow a front where $\kappa = ak^2$ is the local geometric curvature at a convex wave crest and m_i the reduced ion mass. Force balance produces a quadratic equation whose solution yields the velocity ceiling as a

function of the shock potential Φ and ramp width L_r , as well as the ripple amplitude a and wavenumber k . This ripple-limited velocity may be expressed in terms of the

dimensionless variable $\tilde{v}_{pert} = \frac{B}{cE_{\perp}} v$, as $\tilde{v}_{pert} = \Gamma_1 \left[\sqrt{1 + \frac{2}{\Gamma_1}} - 1 \right]$ where $\Gamma_1 \equiv \frac{\xi_1^2 L_r}{2\Phi \tilde{m}_i a k^2}$

and the fitting parameter ξ_1 , which represents the magnetic field strength as a function of ramp penetration depth. In the limit of vanishing corrugation, the function correctly returns the ramp-limited velocity ceiling, that is, $\tilde{v}_{pert} \rightarrow 1$ as $a, k \rightarrow 0$.

As verified in Appendix C, this function fits the simulation data only for wide ramps or weak perturbations, more specifically, when the perturbation parameters violate the critical constraint $\tilde{m}_i a k^2 < 0.18$ in normalized units. We infer from this failure that the very fastest ions do not in fact traverse the shock as initially assumed under these conditions. For completeness, we also treat ion trajectories at troughs, and conclude that the mathematical model based on these assumptions is ill posed, since the solution is not defined outside an extremely limited range of parameters, evident from the critical constraint.

Curvature Model

As seen in the trajectory plots, some ion trajectories temporarily escape the ramp upstream where, unconstrained by the normal electric field, their gyro-phase is free to evolve, whence they breach the shock on next encounter. We assume that this occurs when the upstream ion trajectory radius exceeds that of the curved front, and prove in Appendix C that the cyclotron radius may serve as a sufficiently accurate approximation for the radius of these trajectories at crests. The escape criterion for such ions is then found by equating the curvatures of the trajectory and shock front, again yielding a

velocity ceiling function. We include several corrections for minor effects in order to establish confidence in our model and computational results over a wide range of parameters. This function is given in term of the same dimensionless velocity ceiling

$$\tilde{v}_{pert} = \max \left\{ \left[1 + \frac{2}{\Gamma_2} \right]^{-1} \sqrt{1 + m_i^{-1} v_{SD}^2 \left[1 + \frac{2}{\Gamma_2} \right]^2}, \hat{v}_{SD} \right\}$$

again with dimensionless parameter $\Gamma_j \equiv \frac{\xi_j^2 L_r}{2\Phi \tilde{m}_i a k^2}$ and shock drift ceiling v_{SD} as

derived in Chapter 3, roughly equivalent to the shock compression ratio.

As may be verified from the data section in Appendix C, this predictive function agrees with the simulation data over the entire range of parameters explored, including ramp widths to $L_r = 0.3 \lambda_{pi}$ and masses up to that of singly ionized helium. This is particularly remarkable since one and only one fitting constant was applied uniformly throughout the entire data series, and therefore validates not only the assumptions made in constructing our suppression model, but also the veracity of the simulation code itself.

Finally, we emphasize here that the simple ramp width velocity ceiling function derived in Chapter 3 is completely predominant when the perturbations are small. Conversely, in Chapters 6 and 7 we show from observational data that the perturbation strength in actual shocks overshadows the ramp width velocity limits, which conveniently reduces our predictive function to the simple form $\tilde{v}_{max} = (m_i a k^2)^{-1}$, normalized to the local Alfvén speed. In other words, physical conditions as observed *in situ* at heliospheric shocks represent the limiting case of our model, abrogating the need to employ the more comprehensive predictive function we have derived.

Simulation Results

We now present data from our corrugated front studies. A battery of simulations was performed over a broad range of parameters, including wavenumber $0.1 \leq k \leq 0.7$ in increments of $\Delta k = 0.1$, amplitude $0.05 \leq a \leq 0.7$ with $\Delta a = 0.05$, magnetic field angle $60^\circ \leq \theta \leq 120^\circ$ with $\Delta \theta = 5^\circ$, ramp width $0.025 \leq L_r \leq 0.4$ with $\Delta L_r = 0.025$, shock Mach number $M_A = 3, 5, 7, 9$ and ion mass $m = 1, 2, 4, 8$. We show results from analyses of the velocity ceiling, total accelerated fluxes, and surfing efficiency as defined by the flux of surfing ions promoted above the shock drift ceiling. For the sake of expediency and clarity, we offer abbreviated examples of data sets for exhibition and comparison, and exhibit a more comprehensive set in Appendix D. We also include a discussion of limitations on the length of waves supportable at planetary bow shocks, which we supplement with examples from observations made by the Cluster Mission. From these results and arguments, we justify our choices of range in the wavelength and amplitude of the corrugations imposed in our numerical studies.

Velocity Maxima

For completeness and easy reference, we first recapitulate the example given in the text. In Figure 4-7 below, we show raw data for protons of mass $m = 1$ at a Mach 5 perpendicular shock, in abbreviated parameter sets superimposed by wavenumber $k = 0.1, 0.3, 0.5, 0.7$ and grouped for comparison at four ramp widths

$L_r = 0.05, 0.1, 0.2, 0.3$. The velocities given are normalized $\tilde{v} = v_{\max} \frac{B}{E_\perp}$, so that the

maximum unperturbed velocity for each ramp width at zero amplitude is unity. These unity maxima correspond respectively to the absolute velocities $v_{\max} = 212, 106, 53, 38$

given in units normalized to the upstream Alfvén speed. We also display for comparison the normalized shock drift velocity ceiling at $k = 0.5$ as tabulated at each amplitude increment in the simulations, and shown by the dotted curve labeled ‘SDN’. This not only emphasizes the scale difference between the panels, but also highlights the degree of surfing suppression by comparison with shock drift. Note that the SD ceiling velocity is not strictly invariant, but increases slightly with perturbation strength.

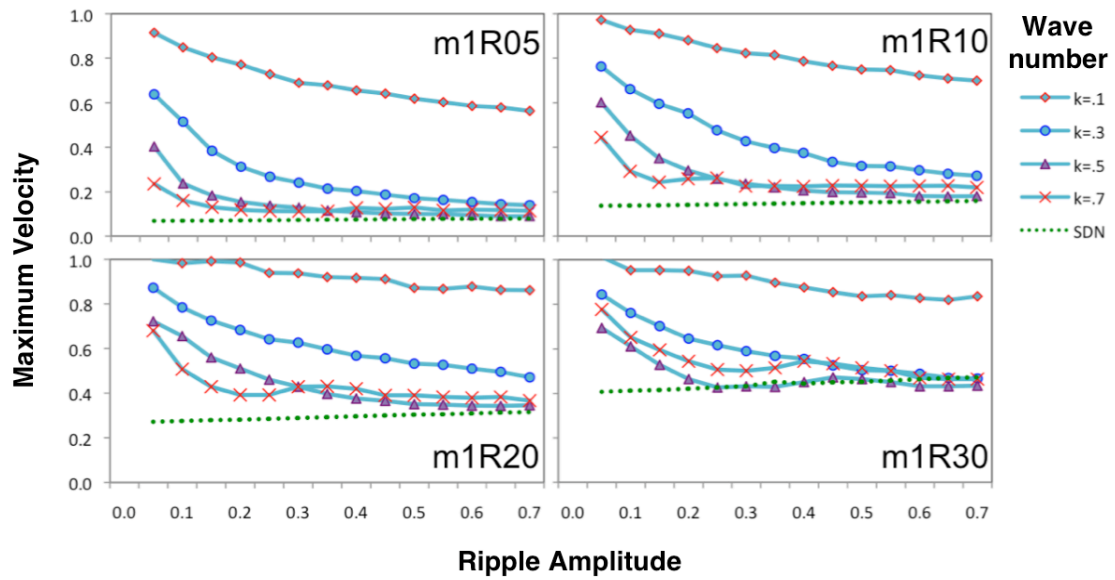


Figure 4-7 Velocity Suppression of Surfing Protons

Results for perpendicular Mach 5 shock are shown at four ramp widths (m1R10 means mass $m = 1$ and ramp width $L_r = 0.10$). Four

wavenumber curves are superimposed in each panel, with velocity normalized to the maximum attainable without ripples at each ramp width.

The scale expands since $v_{\max} \sim L_r^{-1}$, evident by the shock drift velocity SDN, which is nearly insensitive to ramp width.

Suppression of the SS velocity ceiling is immediately apparent. Even at the very thinnest ramp, the smallest wavenumber $k = 0.1$ (longest wavelength) attenuates the

velocity by over 40% at the amplitude $a = 0.7$, while the shortest wavelength corresponding to $k = 0.7$ reduces the maximum surfing velocity all the way down to the shock drift ceiling at half that amplitude. Bearing in mind that the displayed velocity scale contracts with increasing ramp width, at the widest ramp shown where $L_r = 0.3$, the $k = 0.7$ curve quickly reaches its velocity floor, while the $k = 0.5$ curve is not far above it in absolute velocity, and suppression is still substantial at the modest wavenumber $k = 0.3$. Since we expect the maximum velocity for strong perturbations to drop roughly as $v_{\max} \sim (mak^2)^{-1}$, it is not surprising that wavenumber dependence is more pronounced than that of amplitude. The difference between wavenumber curves at the widest ramp is less drastic, simply because the maximum unperturbed velocity is already substantially attenuated by the wider ramp.

We extended these studies to include deuterons and singly ionized helium, and include these results in Appendix D. In short, the effect of front ripples becomes worse with increasing ion mass, consistent with the abbreviated velocity suppression function. We say that the surfing mechanism fails completely when the velocity ceiling reaches the Shock Drift velocity. This floor is reached far more quickly for deuterons, that is, at lower amplitudes and longer wavelengths than those required for similar suppression of protons. For helium ions, the mechanism fails catastrophically at the modest wavenumber $k = 0.3$ even at the very weakest applied amplitude, $a = 0.05$. Extrapolating these results – which agree remarkably well with our predictive theoretical function – to heavier ions shows that singly ionized oxygen O^+ is incapable of surfing a front even with ripples of wavelength the order of the bow shock standoff, and turbulent amplitudes well below any actually measured *in situ*.

Velocity Summary

The plots shown in Appendix D along with analyses therein, give us a convincing and comprehensive overview of the disruptive effects of front corrugations upon the shock surfing mechanism. Even at very modest amplitudes and long wavelengths, velocity suppression can be significant and, for physical parameters actually observed or reasonably expected at heliospheric shocks, is sufficient to render the surfing mechanism completely inconsequential when compared to shock drift. Further, disruption of the surfing mechanism is more pronounced for heavier ions, especially boding ill for the application of shock surfing as a DSA pre-injection mechanism to explain the observed oxygen fluxes in the anomalous component of the cosmic ray spectrum.

Cluster was the first mission capable of resolving ripples in the Earth's bow shock and in fact detected them, and we discuss these and other mission observations more thoroughly in Chapter 6. Briefly, the ripple amplitude was measured at $a \simeq 76 \text{ km}$ and the wavelength $\lambda \simeq 700 \text{ km}$. Based on the ambient shock parameters as reported by Narita (2005), the values normalized to the local ion inertial length become $\tilde{a} \simeq 1.5$ and $\tilde{k} \simeq 0.47$. Since these are both large, we ignore the ramp width limiting term, and approximate the velocity ceiling as $v_{\text{max}} \simeq (mak^2)^{-1}$. For protons, this gives the value $v_{\text{max}} \simeq 1.6 u_1$, well below the shock drift ceiling. Hence, we conclude that shock surfing was not capable of functioning at the bow shock, at least not during the period of these observations.

We have demonstrated that surfing velocity is constrained at the Earth's Bow Shock by spatio-temporal factors, the Parker angle and shock reformation time scales.

However, if the ripples observed by Cluster are a ubiquitous feature of planetary bow shocks throughout the system, this poses serious doubts as to whether shock surfing can be supported anywhere within the heliosphere.

Accelerated Fluxes

Besides perturbation effects on surfing velocity maxima, we are also interested in the fluxes of pickup ions accelerated by Shock Surfing (SS) and Shock Drift (SD) as defined by the fraction of the total incident Pickup Ion (PI) flux which is acted upon by each mechanism. Clearly, if the velocity ceiling for any mechanism falls below the minimum DSA injection threshold, pre-injection cannot be achieved at all. However, the flux of surfing ions that exceed the SD ceiling is more significant still since, even when $v_{SS} \gg v_{SD}$, the flux above the SD ceiling may be of little consequence if the distribution tail is tenuous. Hence, we must also consider as a measure of efficiency the flux of pickup ions promoted above that threshold even when the velocity ceiling is not completely suppressed. Since we have not yet established a DSA threshold, we shall first investigate the most conservative value possible, setting our threshold marker at the shock drift ceiling, or roughly three times the piston velocity.

We begin by examining the total fluxes of protons acceleration by each mechanism at a purely perpendicular Mach 5 shock. As we did for velocity maxima, we investigate the fluxes of more massive ions in Appendix D and summarize the results here. We show in Figure 4-8 the total fluxes of protons accelerated by each of the three ‘tagged’ mechanisms at selected ramp widths in groups of four perturbation wavelengths, each plotted against the full range of perturbation amplitudes. For comparison, we superimpose the flux curves of shock surfing ions (FSS), crossing ions

(FCC), and single-bounce (FBB) mechanisms. We also plot the sum of the latter two, which we have shown are essentially the same mechanism and refer to as aggregate shock drift (FSD). The flux function values plotted on the ordinate represent the decimal fraction of the total incident PI population.

The acronym ‘m1R10’ represents mass $m = 1$ and ramp width $L_r = 0.10$, a form we use throughout this paper. We note that the aggregate shock drift flux (FSD) remains nearly invariant with perturbation strength both in amplitude and wavenumber, while the ‘crossing’ (FCC) and ‘single-bounce’ (FBB) fluxes diverge linearly with increasing amplitude. Although the surfing flux (FSS) increases slightly at the strongest perturbation strength shown, the difference is small in absolute terms. Finally, the shock drift flux is roughly an order of magnitude below that of the total PI population, most significant, the shock surfing flux is an order of magnitude below that. This illustrates the importance of an extended SS tail since, if the shock drift and surfing velocity ceilings were comparable, the SD flux would dwarf that of SS, rendering the mechanism inconsequential.

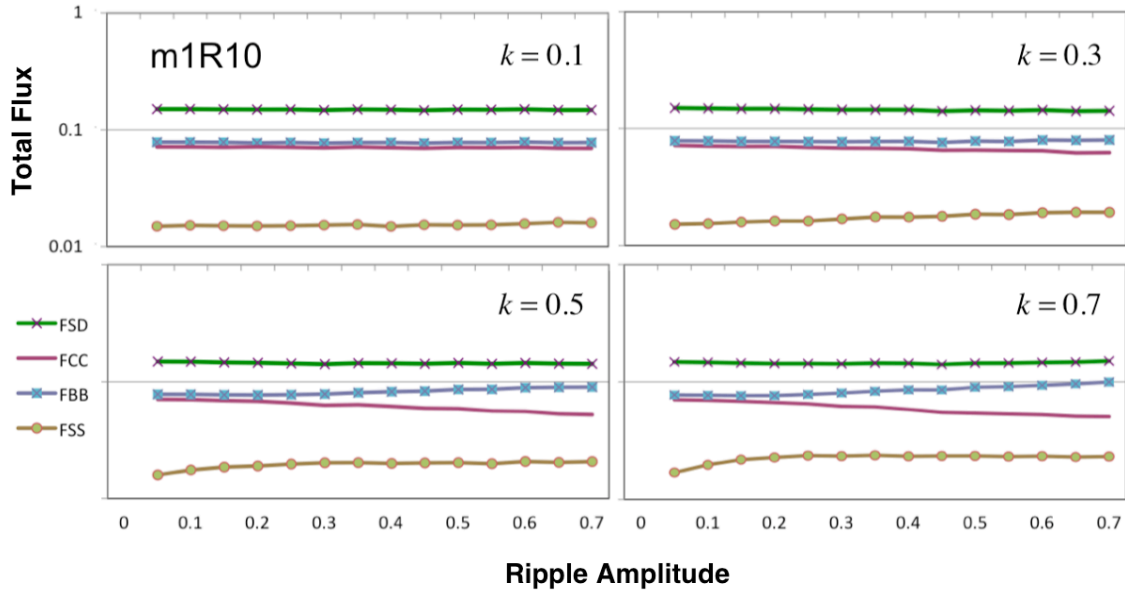


Figure 4-8 Total Accelerated Proton Fluxes

The panels in the quad represent four wavenumbers as shown. The aggregate Shock Drift flux (FSD) composed of the ‘Crossing’ (FCC) and ‘single-bounce’ (FBB) components is insensitive to perturbation strength. Note that for the ramp shown, $L_r = 0.10$, the surfing flux FSS is an order of magnitude below FSD.

One feature of interest is the mass dependence of the surfing fluxes, which may be seen at a glance from the following set of four plots shown in Figure 4-9 below. As in the figure above, the abscissa indicates the perturbation amplitude in reduced units $\lambda_{pi} = c / \omega_{pi}$ and the ordinate the decimal fraction of the total incident pickup ion population. The four panels have identical scales for better comparison, and represent two wavenumbers (top and bottom) and two ramp widths (left and right) as indicated in the legends. These plots show that the total surfing flux is inversely proportional to ion mass, but fairly insensitive to ripple amplitude and wavelength. At the same time, the

aggregate SD flux is almost completely independent of ion mass, which once again bodes ill for the fluxes of heavier surfing ions when compared with those of Shock Drift.

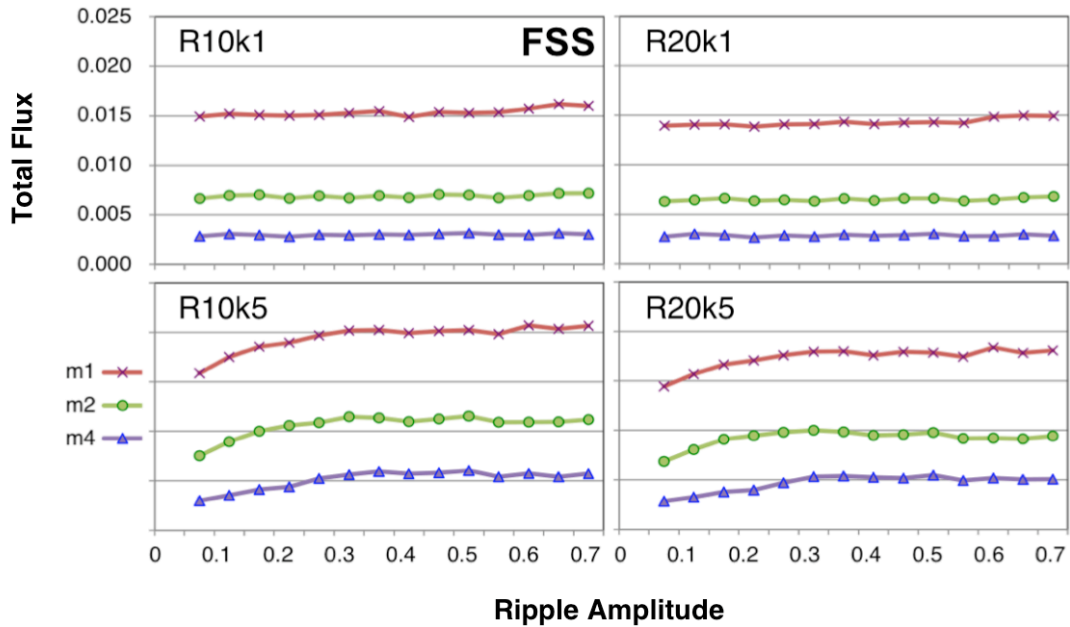


Figure 4-9 Shock Surfing Mass Comparison

Ramp width $L_r = 0.1$ is shown in left panels and $L_r = 0.2$ on the right.

Perturbation wavenumbers $k = 0.1$ and $k = 0.5$ are shown top and bottom respectively. The surfing fluxes are inversely proportional to mass, consistent with the derived dependence in the shock front reflection coefficient. All panels have the same vertical scale.

These plots are entirely representative of the ensembles explored in parameter space, as can be verified in the Flux section of Appendix D, and the features of interest may be summed up thus:

- The aggregate Shock Drift flux, composed of the ‘crossing’ and ‘single-bounce’ components, is almost completely insensitive to variation in perturbation strength parameters, ramp width and ion mass.

- The ‘crossing’ component flux FCC decreases with perturbation strength while that of the ‘single-bounce’ species increases roughly linearly in both wavenumber and amplitude, but independent of shock ramp width.
- The crossing flux FCC increases with mass, but only marginally past $m = 2$, while the single bounce flux FBB is inversely proportional to mass.
- The flux of the aggregate SD ensemble is nearly an order of magnitude below the total incident pickup ion flux.
- The Shock Surfing flux is roughly an order of magnitude lower than SD, and thus two orders of magnitude below the total incident PI population.
- The total shock surfing flux is ramp-invariant and relatively insensitive to small perturbations, but increases at higher ripple strength for heavier ions. However, the total surfing flux should not be taken as a measure of efficiency, which we investigate next.

Surfing Efficiency

The most revealing and perhaps the most important observation we have made is the depopulation of the shock surfing distribution tail due to corrugations in the shock front. We have already seen that the surfing mechanism fails utterly when its velocity ceiling closely approaches that of the shock drift mechanism, but we have yet to quantify the mechanism efficiency otherwise. Since we as yet have no specific DSA threshold, we begin by examining the surfing flux of ions above velocities near the shock drift ceiling. In the figures below, we show plots of the surfing ion fluxes (given as decimal fractions of the total PI flux) that exceed specific critical velocity markers: ‘F3V’ for three

times the piston velocity $3\tilde{u}_{sw} = 3M_A$, ‘S>C’ the velocity ceiling of the ‘crossing’ ions, and ‘S>B’ for the single-bounce ceiling.

As we have seen, the aggregate shock drift flux is nearly invariant while the velocity ceiling increases with perturbation strength at large amplitudes and wavenumbers. However, in the spirit of presenting the most conservative estimates possible, and the possibility that our simplified model may have introduced artifacts, we will concentrate upon F3V, but show data for all three criteria for the sake of completeness. While the velocity $3u_{sw}$ is slightly above the unperturbed Shock Drift ceiling, F3V conservatively indicates when the surfing mechanism has failed catastrophically, especially in the absence of a definite DSA threshold. It also has the virtue of being completely unambiguous and, in light of some estimates in the recent literature asserting the DSA injection threshold velocity $v_{DSA} \approx 3u_{sw}$, directly addresses these proponents as well.

As we did for velocity maxima and fluxes, we now offer an abbreviated set of representative data ensembles, and refer the reader to the Efficiency section of Appendix D for a complete exposition of the simulation data and their analyses. These figures illustrate at a glance the parametric dependence of the flux tail attenuation, and show the profound extent to which front perturbations deplete the population of surfing ions in the distribution tail. For precision, we recapitulate our definition of efficiency by the fraction of the total incident PI flux promoted above a specific threshold velocity, for example, three times the ram speed $F3V(L_r, m_i, a, k) = \int_{3u_1}^{\infty} f(v_{ss}) dv$ where $f(v_{ss})$ is the surfing distribution density function. We specifically define this measure as a function

of ion mass, ramp width, ripple amplitude and wavenumber in anticipation of this dependence.

We begin by comparing the effects of two wavenumbers (top and bottom) at two ramp widths (left and right) on the function F_{3V} as shown in Figure 4-10 below. We compound the comparison with superimposed curves for three different masses $m = 1, 2, 4$, respectively the masses of protons H^+ , deuterons D^+ and singly ionized helium He^+ . The scales are log-linear, and the coordinates normalized as usual.

It is immediately apparent that the unperturbed efficiency decline is more rapid than inversely proportional with mass. Also evident is the negligible effect on all mass species at the long wavelength corresponding to $k = 0.1$, stronger suppression at $k = 0.5$ and the much stronger effect on heavy ions at this wavelength, more so at the wider ramp. As shown in the appendix, the other measures $F_{>C}$ and $F_{>B}$ show almost identical dependence.

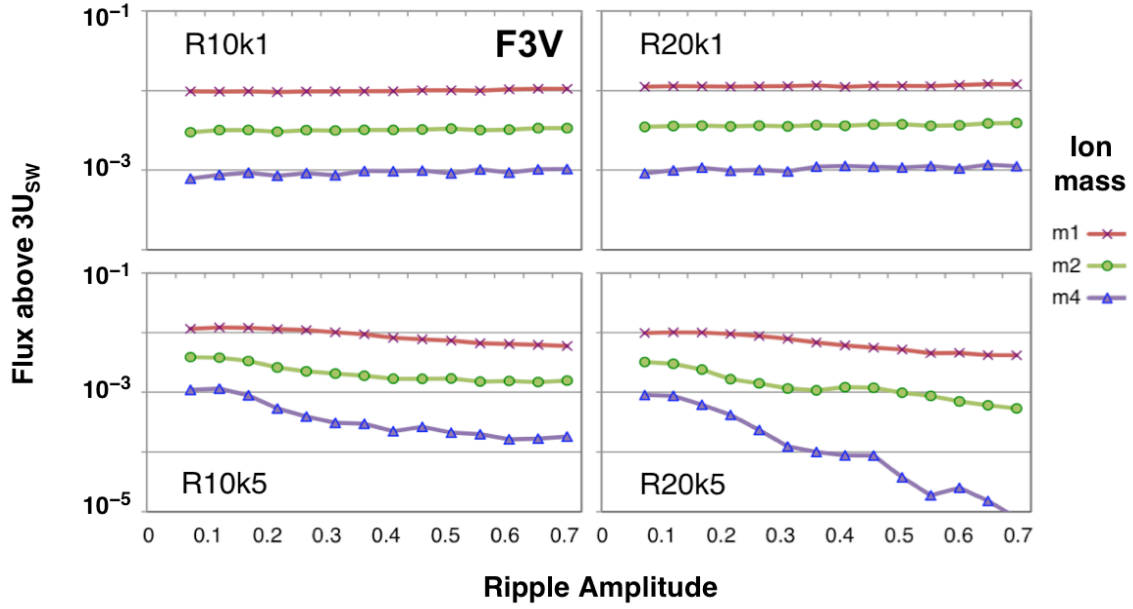


Figure 4-10 Surfing Efficiency F3V

Fractions of the total PI population promoted by SS above the velocity $3u_{SW}$. The unperturbed values are inversely proportional to mass, and show greater suppression at the wider ramp $L_r = 0.2$ on the right compared with $L_r = 0.1$ on the left. Effect is more pronounced for the heaviest He^+ ions shown as 'm4'.

At the Earth's bow shock, the wavenumber $k = 0.1$ is equivalent to a wavelength $63\lambda_{pi} \sim 6300\text{km}$, which is extremely gentle, especially when compared to ripples observed by Cluster at the bow shock. Indeed, wavenumbers exceeding $k = 0.5$ are not uncommon, as we show in Chapter 6. As shown in the appendix, the surfing efficiency by all three measures exhibits the expected dependence on every parameter – mass, ramp width and perturbation strength.

For the sake of expediency, we offer a more comprehensive visualization of data contained in the appendix in order to make the effects of front corrugation on surfing

efficiency abundantly clear. We demonstrate the ramp dependence with protons shown in Figure 4-11 below. Three efficiency measures F3V, S>C and S>B as plotted as separate curves. The abscissa is the ripple amplitude, the ordinate in powers of ten represent the fraction of the total incident PI population, and each panel represents a different ramp width as shown in the legends, e.g., the acronym ‘m1R20k5’ means mass $m = 1$, ramp width $L_r = 0.20$ and wavenumber $k = 0.5$. Here we hold mass and wavenumber constant for ramp width comparison. At the moderate wavenumber $k = 0.5$ shown, protons are only mildly affected at narrow ramps, and the flux is reduced by less than an order of magnitude at the largest amplitude shown, while suppression is more severe at wider ramps, especially for the measure function S>B.

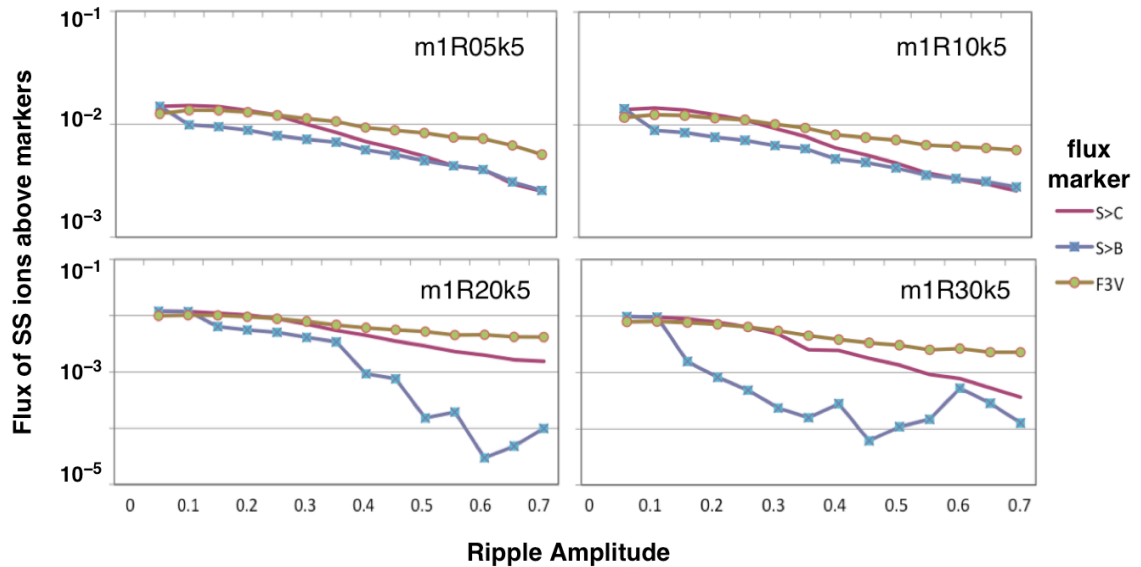


Figure 4-11 Efficiency Functions for protons

Curves are shown for all three measures at four ramp widths from narrow $L_r = 0.05$ to broader $L_r = 0.30$ all at moderate wavenumber $k = 0.5$

In Figure 4-12, we show four panels with different wavenumbers, ramp widths and ion mass all plotted against amplitude. The selected parameters and the vertical log scale highlight the profound effect of ripples on shock surfing. The attenuation in fluxes is orders of magnitude in scale, and especially severe for the efficiency measure $S > B$. Although the single-bounce velocity ceiling is variable and sensitive to front perturbations, the dramatic appearance of the measure $S > B$ should not be entirely dismissed, since it strongly indicates a rapid drop in the surfing distribution function as the velocity approaches $v \sim 4u_1$, only marginally greater than the F3V threshold and ‘crossing’ velocity ceiling VCC.

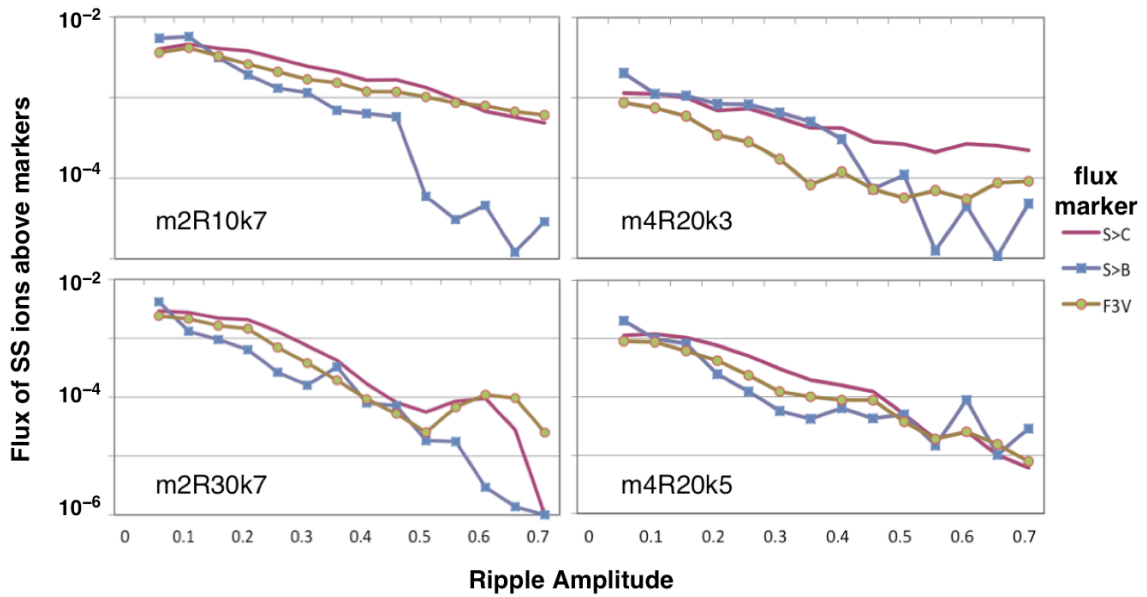


Figure 4-12 Representative Efficiency Ensembles

Three efficiency measures are plotted in each of the four panels of various mass, ramp width and wavenumber. The acronym m2R10k7 means mass $m = 1$, ramp $L_r = 0.10$ and wavenumber $k = 0.7$. Erratic appearance of $S > B$ may be due to paucity of data points, or real features such as resonant velocity enhancement.

As seen from the data in Appendix D, the maximum single-bounce velocity VBB increases with increasing perturbation strength. However, this increase is by no means extreme, as demonstrated by the example shown below in Figure 4-13. In this case, we plot the velocity VBB in units of the upstream solar wind speed, making clear that the deviation above $3u_1$ is small – approximately $0.5u_1$ or about 16% – which validates our assertion that the surfing efficiency measure ‘S>B’ is not unreliable, but rather a strong indication of the rapid decline in surfing flux above the shock drift ceiling.

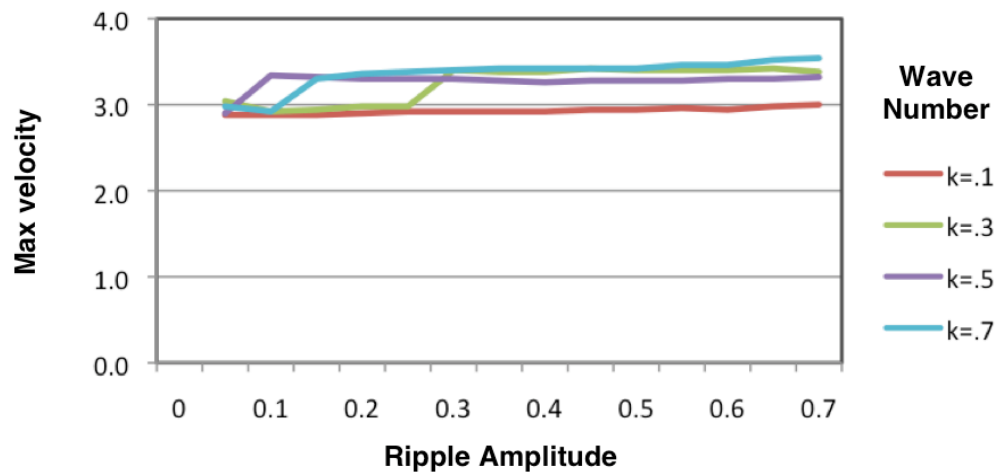


Figure 4-13 Single Bounce Velocity Ceiling

Figure represents protons at Mach 5 shock of ramp width $L_r = 0.05 \lambda_{pi}$ at four corrugation wavenumbers. Velocity is given in units of the upstream wind speed.

Efficiency Function

We now seek to quantify the effect of perturbations on shock surfing efficiency as we did when producing a predictive mathematical function for the suppression of the surfing velocity ceiling. Recall that we define surfing 'efficiency' as the flux of surfing ions promoted above some collective ensemble of Shock Drift velocity ceiling markers

$$\mathfrak{S}(m_i, L_r, a, k) = \int_{v_{SD}}^{\infty} f_{SS}(v) dv, \text{ where the distribution function } f_{SS}(v) \text{ is normalized to the}$$

total incident Pickup Ion flux. Noting that the plots were nearly linear, visual fits to the data were made on each of the function ensembles across the entire data set, and an analytic predictive function found with remarkably accurate global agreement. Taking the equally weighted mean of the three measures above and discarding extrema (especially in $S > B$) we can express the efficiency globally by the simple exponential function

$$\mathfrak{S}(m_i, L_r, a, k) = m_i^{-1} 10^{-p(m_i, L_r, a, k)} \text{ where } p(m_i, L_r, a, k) = 2 + 10 m_i^{1/2} a k L_r$$

The 29 points of data tabulated by mass, ramp width, amplitude and wavenumber were sorted in ascending order by the observed exponent and plotted against the above function on those parametric variables as shown below in Figure 4-14.

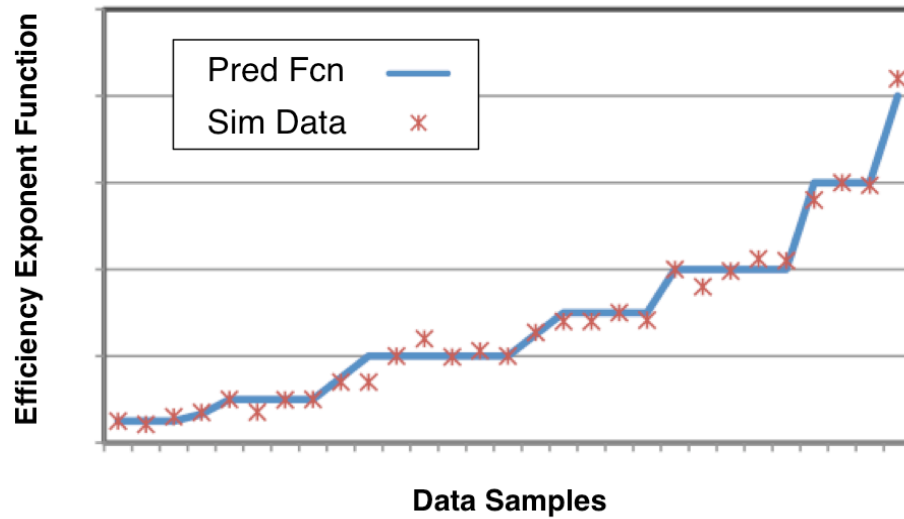


Figure 4-14 Plot of Efficiency Scaling Function vs. Data.

Flux suppression tabulated from data over a range of values: mass, ramp width, ripple amplitude and wavenumber. Attenuation data were sorted in ascending order along with predictive function values.

See text for detailed discussion.

Naturally, it would be better to analyze all the surfing distribution tails and discover a similar function that would universally prescribe the tail profile, in terms of a power-law or exponential function, the form we have seen in the baseline studies. This would be an extremely painstaking exercise unless automated, and we encourage the reader to do so if so inclined. However, we believe that this would be of little consequence, unless it could be proved that the DSA threshold were to lie extremely close to the shock drift ceiling. Further, its application within the heliosphere would be entirely academic since, as we have shown, other mitigating factors such as spatio-temporal constraints and shock reformation are in force. We have pursued the matter this far only to lend additional weight to these other factors, and offer this flux

suppression analysis for future application, possibly to shocks such as those generated by supernovae and other violent cosmic events.

Summary of Flux Analysis

We have seen that the flux of surfing ions FSS is inversely proportional to mass. As with the aggregate shock drift flux FSD, the unperturbed shock surfing flux is relatively insensitive to ramp width, and the effect of perturbations over a wide range of strength parameters quite modest.

However, ripples along the shock front produced a substantial suppression in the population of the SS distribution tail. The ‘efficiency’ defined as the fraction of incident PI accelerated above the aggregate shock drift velocity ceiling – an extremely conservative measure – shows that ripples not only suppress the surfing velocity ceiling, but reduce the efficiency of the mechanism as well. Even for reasonable amplitudes and wavenumbers found at heliospheric shocks, the degree of this attenuation is sufficient to render SS entirely ineffective as a pre-injection mechanism, as we verify with observational data in Chapter 6. This is particularly applicable to the observed oxygen flux in the anomalous cosmic ray spectrum, since we have proven that acceleration of heavy ions is especially affected as shown by the efficiency function – once in the factor m^{-1} and once again as $-m^{1/2}$ in the exponent.

Magnetic field angle study

In the interest of computational economy, a cursory study with protons was performed for a thin ramp $L_r = 0.05$ at Mach 5 over a discrete set of corrugation amplitudes and wavenumbers, representative of the range we explored. We chose a set of identical values, beginning with $a = k = 0.1$ and increasing to $a = k = 0.7$ as shown in

the figure legend. We plot these superimposed curves in Figure 4-15 as for comparison. Even for small perturbations, the effect is moderate where SS is effective at all. Outside the range $80^\circ \leq \theta \leq 100^\circ$, ripples are altogether negligible when compared to field angle suppression. However, for $a = k = 0.3 \Rightarrow ak^2 \sim 10^{-2}$, the velocity attenuation becomes severe, abrogating a more detailed study. As evident from the graphic, the effect of ripples is most pronounced for field angles near perpendicular.

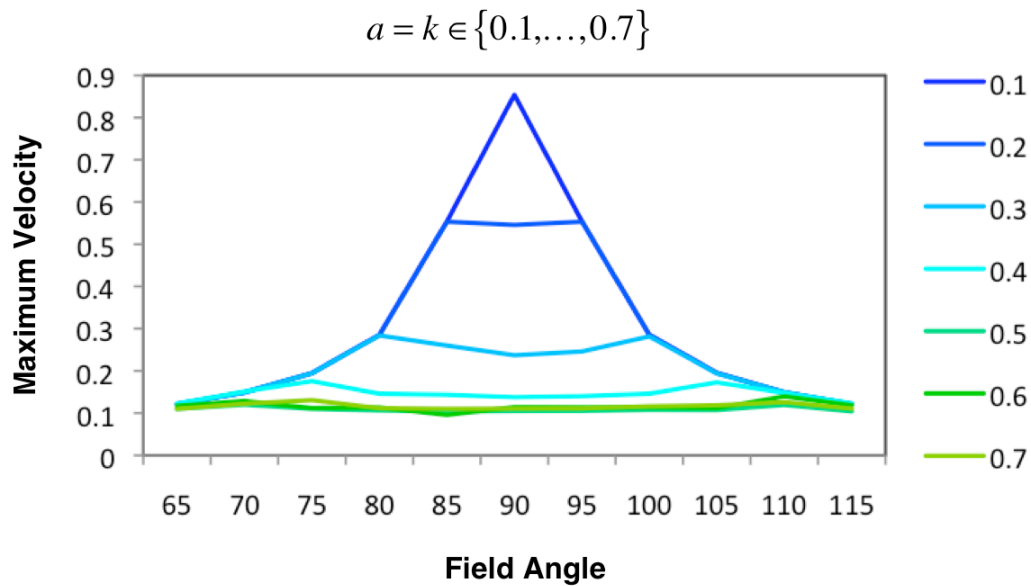


Figure 4-15 Perturbed Velocity Ceiling versus magnetic field angle
Velocity is normalized to the unperturbed perpendicular ceiling.

In addition to the velocity suppression caused by ripples in the shock front, we also note that the fluxes of accelerated ions are also attenuated beyond those due simply to the magnetic field angle. As shown in Figure 4-16 below, these effects are comparable in magnitude to those seen in the velocity maxima above, and further demonstrate the deleterious effect on shock surfing imposed by corrugations in the front. One feature of

note is the strong flux suppression produced by moderately weak perturbations when the magnetic field angle is very nearly perpendicular, and effect which diminishes with obliquity within the range $85^\circ < \theta < 95^\circ$. Another, is the profound effect at $a = k = 0.3$ which reduces the flux within the entire prime surfing range to that when $\Delta\theta = \pm 10^\circ$.

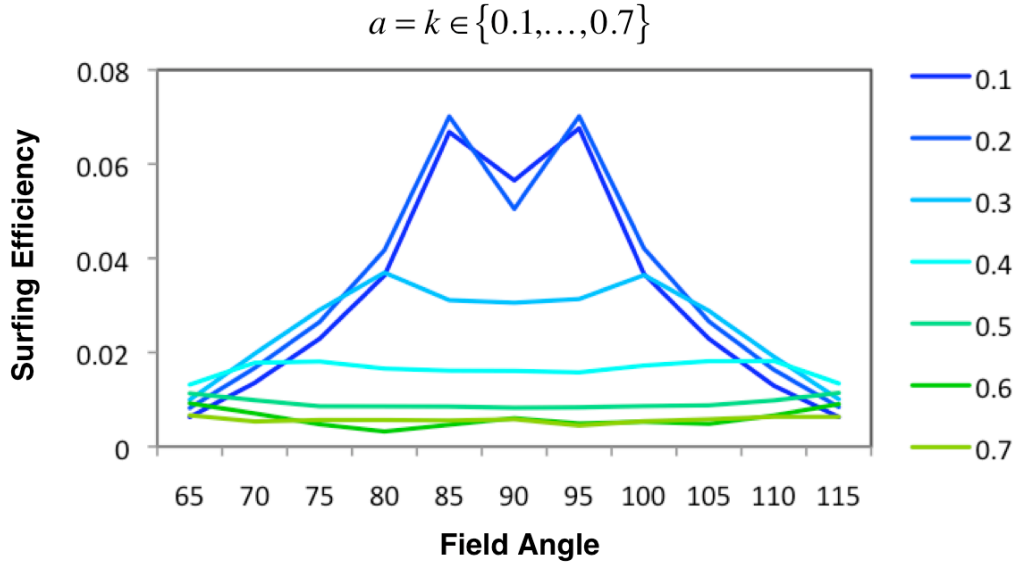


Figure 4-16 Surfing Efficiency versus magnetic field angle

The family of curves represents fluxes of ions promoted above $3u_{sw}$ plotted against magnetic field angle, each at one of several perturbation values from $a = k \in \{0.1, \dots, 0.7\}$.

Summary and Conclusions

We have shown that ripples in the shock front degrade the shock surfing mechanism by suppressing the maximum attainable velocity. We have advanced a physical model to explain this suppression and derived a tractable expression for calculating the velocity ceiling as a function of the perturbation parameters. We have

shown that the flux of ions promoted above the shock drift ceiling becomes tenuous with increasing perturbation strength. Within physically reasonable parameter ranges, ripples can attenuate this population by several orders of magnitude, and we have produced an analytic function from the data sufficiently accurate for predictive purposes.

Although we feel confident in our physical analyses of the mechanisms responsible for shock surfing disruption, for completeness we also address the effects of corrugations upon the theoretical constructs of shock surfing as put forth by Lee et al (1996). We present results and discussion of this investigation in Appendix E, and conclude that the imposition of ripples in the shock front rapidly cause surfing ion trajectories to become erratic, thus undermining the mathematical foundations of Lee et al. We also address in some detail, the ranges of shock Mach numbers and perturbation wavenumbers we have chosen for this study, which may be found in the discussion and comments section at the end of Appendix C.

As a final comment on perturbations we note that front corrugations may form even in absence of externally driven turbulence, due to the manifold plasma instabilities inherently associated with astrophysical shocks. Indeed, one entire subject of spirited discourse is the question of whether such perturbations induce turbulence or vice versa. We have seen that effective shock surfing requires the thinnest possible sustained shocks but, while steeper shocks are certainly better for surfing, they may also be more favorable for the growth of instabilities.

We now turn to case studies in the heliosphere, and examine the data from several missions sent to planetary shocks and the heliospheric termination. We analyze and interpret this data applying the predictive functions just derived, and investigate whether the shock surfing mechanism is in fact viable.

Chapter 5 Velocity Diffusion

We have shown that Shock Surfing (SS) can, under idealized conditions, accelerate Pickup Ions to high velocities at sufficiently steep and smooth quasi-perpendicular shock fronts. Even in the presence of plasma waves near the front, SS may still remain capable of promoting ions to substantial energies. On the other hand, we have also demonstrated several ways in which deviations from the ideal shock can suppress not only the velocity ceiling, but also the flux of surfing ions promoted above the velocities achieved by Shock Drift (SD).

Recall that SS was originally conceived in conjunction with Anomalous Cosmic Ray (ACR) production as a ‘pre-injection’ mechanism for Diffusive Shock Acceleration (DSA) at quasi-perpendicular shocks. Since DSA is believed to require a minimum injection velocity, it is of paramount importance to precisely establish that threshold in order to determine whether SS can remain robust as a pre-injection mechanism in the presence of turbulence.

In this chapter we investigate velocity diffusion due to magnetic turbulence and various resonant and non-resonant wave-particle interactions and present the results of theoretical and computational studies in order to establish the dependence of the diffusion rate on ion mass and velocity, turbulent energy density and wavelength. There is general agreement that DSA requires spatial and temporal scales inconsistent with those at planetary bow shocks. We confirm this by showing that velocity diffusion is limited by the turbulent wavelength ceiling, that is, constraints on waves supportable downstream of such shock systems which in turn place a ceiling on the velocity attainable through diffusion.

To discover whether SS may be of any consequence in the absence of DSA at planetary shocks, we begin with a simple criterion necessary for downstream ions to return to the front, and show that none of the diffusion rates seen in simulations satisfy this requirement, even for the most potent turbulent parameters likely to exist at shocks within the heliosphere.

We then apply the results of these diffusion studies to establish a phenomenological form for the DSA scattering parameter in preparation for analysis of the Termination Shock (TS). We argue that the Webb form of this parameter is inconsistent with our conclusion that scattering is more effective on lighter, slower moving ions. Since velocity diffusion processes disfavor heavy ions, we conclude that the DSA mechanism may therefore be unlikely to account for the observed ACR fluxes heavy ions especially those of Oxygen. We also show that, since the scattering strength and therefore the diffusion rate are inversely proportional to ion velocity, there should be no DSA injection threshold velocity at all – a result validated by simulations of both thermal and pickup ions. Finally, based on a variety of diffusion studies, we question the ability of DSA or any other scattering-based mechanism to promote ions to the extremely high energies observed in the CR spectrum.

We shall restrict our investigation and analysis to the nearly perpendicular case for several reasons. First, we have seen that SS ceases to be effective for magnetic angles deviating from the perpendicular by more $\Delta\theta \sim 10^\circ$. It is, in fact, primarily in the region $80^\circ < \theta_{Bn} < 100^\circ$ where SS was conceived to facilitate multiple shock transits across magnetic field lines. Further, such transits produce less effective acceleration with increasing obliquity since the convective electric field $E_C = \frac{u}{c} \times B$ diminishes as

$\sin \theta_{Bn} \rightarrow 0$, which in turn lowers the energy gain per transit due to the magnetic gradient drift. We begin with a brief exposition of the DSA mechanism and its dependence upon turbulent scattering.

Diffusive Shock Acceleration

In a seminal paper, Webb et al (1995) presented exhaustive analyses of the DSA mechanism, including a Green's Function for calculating various moments of interest. In short, the diffusive shock acceleration process relies on multiple front crossings by turbulent scattering and associated energy gains through drifts along the shock electric field to promote ions to very high energies. The salient point of interest to our investigation is the assertion that a minimum injection threshold velocity is required in order to begin the process. Simply put, an ion must possess sufficient velocity to return upstream after the first turbulent downstream 'collision'. Thermal ions are believed not to achieve this, but Pickup Ions may, by means of a pre-injection mechanism, reach this threshold.

There is widespread agreement that perpendicular shocks offer the most rapid and efficient means of acceleration. The Webb study concludes that, at a perfectly perpendicular shock, DSA can begin provided that a minimum injection velocity

threshold v_{inj} is attained which satisfies $\frac{v_{inj}}{u_1} >> \frac{3}{4} \sqrt{1 + \eta_1^2}$ upstream, and

$\frac{v_{inj}}{u_1} >> \frac{4}{R_C} \sqrt{1 + \eta_2^2}$ downstream, where R_C is the shock compression ratio and u_1 is the

upstream fluid velocity. The local parameters $\eta_{1,2}$ are dimensionless scalars inversely proportional to the local scattering strength, upstream and down respectively. Since the

authors admit that “the kinetic theory of energetic charged particle transport in turbulent electromagnetic fields, especially the theory of perpendicular diffusion, is complex and not well understood”, this scattering parameter takes on paramount importance in defining the transport coefficients of the diffusion tensor, after the phenomenological approach of Jokipii (1987),

$$\kappa = \begin{pmatrix} \kappa_{\perp} & \kappa_A & 0 \\ -\kappa_A & \kappa_{\perp} & 0 \\ 0 & 0 & \kappa_{\parallel} \end{pmatrix} \text{ where } \kappa_{\parallel} = \eta \frac{vR_C}{3}, \quad \kappa_{\perp} = \frac{\kappa_{\parallel}}{1 + \eta^2} \text{ and } \kappa_A = \frac{vR_C}{3} \frac{\eta^2}{1 + \eta^2}.$$

In the Webb model, the scattering parameter η is defined simply by $\eta = \lambda_{\parallel} / r_g$ where $\lambda_{\parallel}$ is the mean free path parallel to the ambient field and r_g the ion gyro-radius, both dependent on the local plasma parameters. Note that, in the weak scattering limit, $\eta \rightarrow \infty$ as $\lambda_{\parallel} \rightarrow \infty$, hence it is implicit that the minimum threshold velocity scales inversely with scattering strength. Note especially that, this form of the scattering parameter implies that the scattering strength itself scales with the particle momentum $\eta^{-1} \sim r_g \sim mv_{\perp}$, so that particles with higher momentum are more readily scattered, which is counterintuitive.

While these relations have been often cited and incorporated in subsequent treatments, there is by no means general agreement on how to calculate the scattering parameter η . Zank et al (2008) calls this into question since, “for a particle experiencing diffusion through some form of complex magnetic field random motions, there is little reason to assume that the (mean free path) scales with gyroradius.”

In addition, computing the mean free path itself appears to have been no easy matter. One popular estimate of the quiescent solar wind parallel mean free path is

based on a hard-sphere collision cross section $\lambda_{\parallel} \sim 1AU = 1.5 \times 10^8 km$, (e.g., see

Biebert, *et al* 1980) which leads at once to amusing consequences. Interpreting the

gyroradius to be that of the ion in question $\eta = \frac{\lambda_{\parallel}}{r_{gi}(v_{\perp})}$, we may simplify the estimate for

the injection criterion $v_{inj} \gg \frac{3}{4}u_1\sqrt{1+\eta^2}$ by making the approximation $\eta = \frac{\Omega_{ci}\lambda_{\parallel}}{v_{\perp}} \simeq \frac{\Omega_{ci}\lambda_{\parallel}}{v_i}$

whence the equation may be rewritten as a quadratic in the reduced variable $\tilde{w} \equiv \frac{v_i^2}{u_1^2}$

thus, $\tilde{w}^2 - \tilde{w} - K^2 = 0$ where $K = \frac{\lambda_{\parallel}\Omega_c}{u_1}$. Applying local solar wind parameters at the

Earth's Bow Shock, $\Omega_c \simeq 0.5s^{-1}$ and $u_1 \simeq 500 km s^{-1}$, we obtain $K \sim 1.5 \times 10^5$. The

solution $\tilde{w} = \frac{1}{2}\left(1 \pm \sqrt{1+4K^2}\right)$ for $K \gg 1$ can then be approximated by $\tilde{w} = K$ or

$v_{inj} \gg u_1\sqrt{K} \sim 2 \times 10^8 m/s$, which is absurd. For this formalism to be meaningful, the

scattering in question must be due to some wave-particle interaction and a function dependent upon the magnetic turbulence parameters in the region of interest.

To date, no compelling estimate has been established for the DSA injection threshold velocity. Meanwhile, articles continue to appear in the literature covering a broad variety of estimates. In a conference presentation, Giacalone & Ellison (Aspen, 2005) assert simply that $v_{inj} \simeq 3u_1$ i.e., the same as at parallel shock. Why should this be so? Where is the turbulent strength or other parametric dependence? Note that, even at a perpendicular shock with a ramp too broad for shock surfing, some Pickup Ions will undergo a front crossing. As we have shown, some of these ions proceed downstream

with velocity nearly $3u_1$ and, if this is indeed the injection threshold, there is no need for shock surfing at all since the Shock Drift flux is more than an order of magnitude greater than that of the entire surfing population.

Simply put, the scattering parameter must be a function of the local turbulent energy density and spectra, as well as ion mass and velocity. To address this problem, we now turn our attention to velocity diffusion studies, which will provide the basis for constructing a phenomenological form of the scattering parameter, and also give us insight into the possible role of shock surfing where DSA itself may not be viable.

Front Return Criteria

We begin under the assumption that velocity diffusion by turbulent scattering alone is insufficient to produce the high energies observed in the CR spectrum, and that shocks crossing are an essential component for processes, including DSA, to accomplish this. In order to return to the front, the rate of velocity diffusion must be sufficient to increase the cyclotron radius faster than the guiding center is swept downstream in the bulk fluid. In each orbital period $\tau_{g(2)}$ the guiding center moves a distance $s = u_2 \tau_{g(2)}$ where u_2 is the downstream fluid velocity. To achieve retrograde normal motion, the cyclotron radius must increase faster than this. At a perpendicular shock, this amounts to satisfying the condition $\dot{r}_g > u_2$ or, equivalently, $\dot{v}_\perp > \Omega_{c2} u_2 = \Omega_{c1} u_1$ where Ω_c is the local ion-cyclotron frequency, and the second equality follows from the fact that $u_1 B_1 = u_2 B_2$. At an oblique shock this condition is, of course, relaxed since velocity diffusion along the magnetic field may be more easily

achieved than cross-field diffusion but, as we have shown, oblique shocks do not promote efficient surfing acceleration.

We begin by examining the results of velocity diffusion studies to see whether the condition $\langle \Delta v / \Delta t \rangle \geq \Omega_{c1} u_1$ can be met. Since it is common in numerical simulations to normalize time to the ion cyclotron period Ω_c^{-1} , velocity to the Alfvén velocity v_A and kinetic energy to the piston energy $\varepsilon_0 = \frac{1}{2} m_p u_1^2$, this is equivalent to satisfying the normalized condition $\langle \Delta \tilde{v} / \Delta \tilde{t} \rangle \geq M_A$ or, when results are presented as energy spectra, $\langle \sqrt{\Delta \tilde{\varepsilon}} / \Delta \tilde{t} \rangle \geq M_A$ where M_A is the Alfvén Mach number, $\Delta \tilde{t} = \Delta t \Omega_c$, $\Delta \tilde{v} = \Delta v / v_A$ and $\Delta \tilde{\varepsilon} = \Delta \varepsilon / \varepsilon_0$.

Pitch Angle Scattering

As in many areas of scientific endeavor where problems may be mathematically intractable, those in plasma physics where the degrees of freedom and types of interactions are manifold can be particularly daunting. The question of velocity diffusion has been approached through several formalisms including Vlasov (Boltzmann with collision operators), Fokker-Planck (by way of the Master Equation), Hamiltonian analysis via canonical transformations and generating functions, scaling arguments and Quasi-linear (perturbation) theory. Each requires some form of approximation, which may render useful results within a limited regime. However, these methods must be based upon a kernel containing the physics responsible for driving the interactions in question.

Scattering of particles by waves and Coulomb collisions are known to be different from that produced by classical hard-sphere collisions. However, there are certain features that all scattering should at least intuitively have in common. First, one expects that a particle's deflection should decrease with greater momentum. In the field of accelerator physics, this concept is applied to magnetic deflection by the term 'rigidity' which is defined as $\mathfrak{R} = B\rho = p/q$ where B is the magnetic field, ρ the particle gyroradius, p its momentum and q its charge. Hence it is reasonable to expect that faster and more massive particles are less susceptible to deflection. Similarly, it follows from the Lorentz force law that the deviation of a charged particle from its trajectory in an otherwise uniform electric or magnetic field should scale not only with electric charge, but also with the field strength or energy density of magnetic fluctuations.

The need for a more accurate kernel is non-trivial, as pointed out by B. R. Ragot (1999) who notes that there has been "no satisfying agreement between the quasi-linear predictions and observations. The initial quasi-linear result (Jokipii 1966), for a slab magnetostatic turbulence of Kolmogorov spectrum, gives a rather strong decrease of the mean free path with decreasing particle rigidities..." The failure of quasi-linear theory may be due to misapplication, since its derivation was performed in the limit of weak turbulence. However the cited observations reaffirm our notion that scattering must be less pronounced on particles with higher momenta.

This is not to say that the diffusion rate itself should scale accordingly, but certainly the particle velocity should be incorporated into any random walk model where the step size in velocity space is based on a particle's susceptibility to deflection by the magnetic scattering strength.

Magnetic Turbulence Simulations

Giacalone and Ellison (2000) performed both three-dimensional hybrid and test-particle simulations to study velocity diffusion in a turbulent field. They demonstrate, as we will also show in our own standing wave studies below, that velocity diffusion takes time to evolve. A kinetic hybrid code was employed to create the turbulent thermal background plasma with test particles into which the shock propagates, and the downstream distribution functions for Pickup Ions recorded as shown in Figure 5-1. Different simulation box sizes were employed, each limiting the length of the longest waves along the coordinate axes to the dimensions of the box.

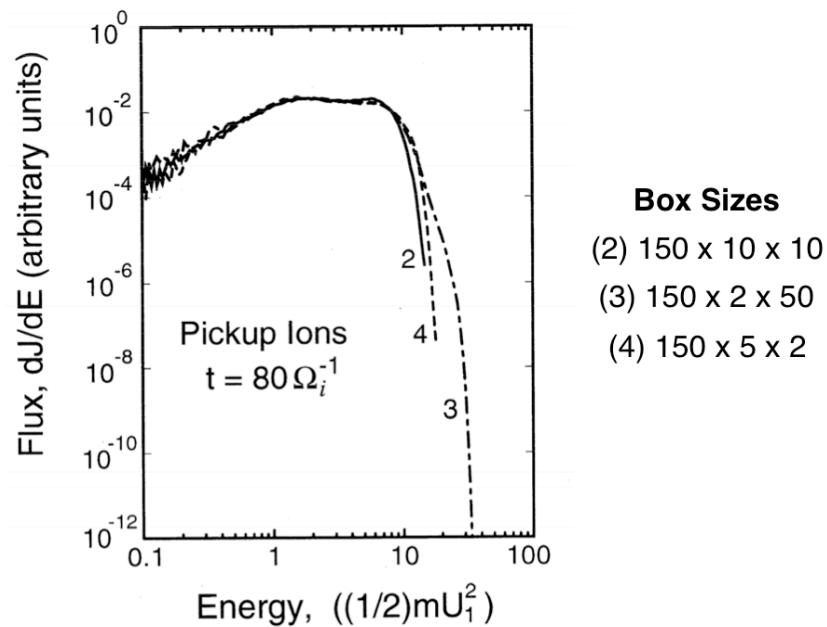


Figure 5-1 Pickup Ion Velocity Diffusion by Stochastic Waves
Maximum velocity grows with time but reaches a wavelength-limited ceiling constrained by the simulation box. The three curves show distribution functions for different box sizes $x_{\max} \times y_{\max} \times z_{\max}$ as shown in units of λ_{pi} . After Giacalone & Ellison (2000)

Parameters were set to model the Heliospheric termination shock, with a broad ramp to ensure the absence of acceleration due to specular reflection. In their simulation scheme, the authors used a closed box with fluid flowing in continuously from the left, resulting in a left-moving shock of Mach number $M_A = 5.4$. A Kolmogorov spectrum was applied to the inflowing plasma, assuming fully developed turbulence, and the total integrated power was set to one-third the ambient solar wind magnetic energy density, i.e., $\frac{\delta B^2}{B^2} = \frac{1}{3}$. While this spectrum may be typical at large astrophysical shocks, data from the Cluster mission indicate only a slightly different spectrum at the Earth's Bow Shock, with exponent $p = -1.9$ rather than the Kolmogorov $p = -5/3$.

For the most pronounced acceleration represented by curve #3, the maximum energy at time $t = 80 \Omega_{ci}^{-1}$ is approximately thirty times the piston energy. Note that, at this energy, the flux is ten orders of magnitude down from the incident PI flux – is this enough to account for the observed CR flux? We now show that, in any case, the velocity diffusion *rate* is insufficient to return ions to the front for any convex shock in the Heliosphere. The maximum normalized downstream kinetic energy after time $\Delta \tilde{t} = 80$ is $\Delta \tilde{\epsilon} \approx 30$. Assuming this is the minimum time required to attain this energy, and linear growth in velocity with time, we have $\sqrt{\Delta \tilde{\epsilon}} / \Delta \tilde{t} \approx 7 \times 10^{-2} \ll 5.4$, nearly two orders of magnitude too small to return ions to the shock front. To study longer wavelength and time domains, turbulent background plasma was synthesized according to the kinetic hybrid simulation data and test particles injected with results shown in Figure 5-2. Although this was not a time study, the shock-normal box dimension gives sufficient information for a rough estimate. The downstream box length $\lambda = 100 \lambda_{pi}$ while the

downstream velocity is roughly twice the upstream Alfvén speed, giving a time $\Delta t \sim 50\Omega_c^{-1}$. This gives a best-case estimate $\sqrt{\Delta\tilde{\epsilon}} / \Delta\tilde{t} \approx 0.28$, again well below the minimum requirement for return to the shock. One feature of interest observable from these studies, however, is that both the diffusion rate and velocity ceiling increase with increasing wavelength.

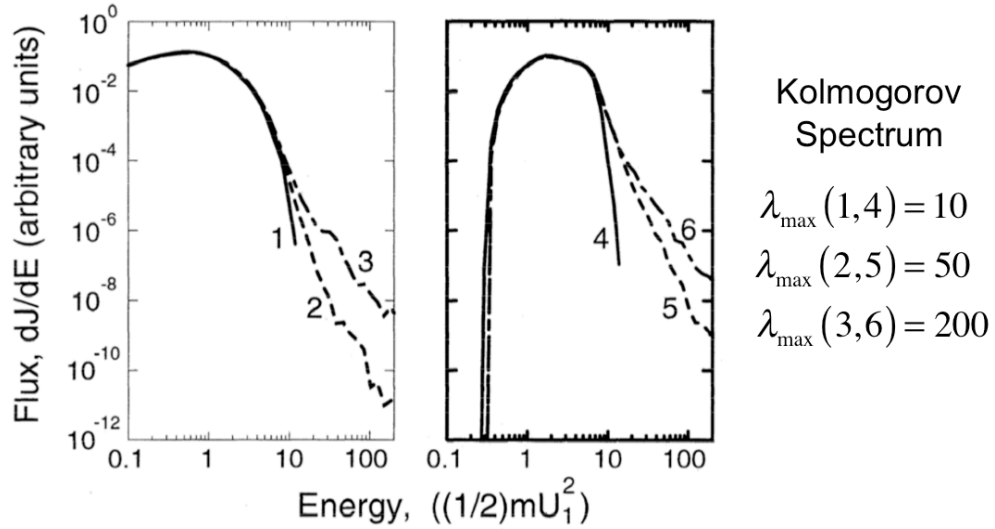


Figure 5-2 Velocity Diffusion Simulations

The left panel represents thermal ions, the right Pickup Ions. The shock propagates through a Kolmogorov spectrum of turbulent magnetic waves with minimum wavelength $\lambda_{\min} = 0.5$ for all simulations, given in ion inertial length. The numbered curves represent maximum wavelengths as shown and described in the text. After Giacalone & Ellison (2000)

Giacalone and Jokipii (2005) extended these simulations, as exhibited in Figure 5-3. In the steady state study shown in the left panel, the amplitudes of these distributions appear at first glance to be linearly dependent (in their notation) on the perturbation strength $\frac{dJ}{dE}(E) \propto (\Delta B/B)^2$ for the values $(\Delta B/B)^2 = 0.1$ and $(\Delta B/B)^2 = 1$.

However, the tail for strength $(\Delta B/B)^2 = 0.3$ suggests that, at higher energies, a ceiling exists which may be due to saturation of the mechanism. Either that, or the velocity diffusion rate slows with increasing energy to the point where the maximum velocity is not attainable within the simulation runtime. The former appears to be more likely, as the following figure shows. The authors explain that, aside from being limited by finite geometry at bow shocks, “acceleration takes time... the ideal power law is not created instantly” which is evident from the blue curve $\theta_{Bn} = 90^\circ$ as shown by one example of their time-dependent studies in the right hand panel in the figure.

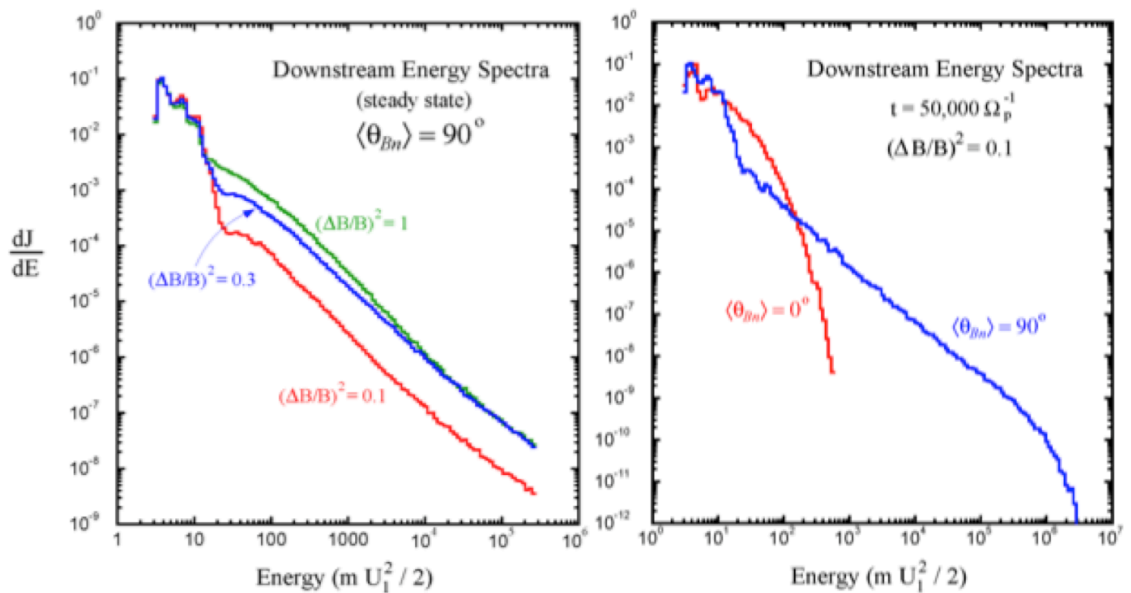


Figure 5-3 Turbulence Driven PI Velocity Diffusion

Left panel shows steady-state linear dependence on wave energy density and right panel shows the time dependence. The blue curve represents the perpendicular shock with modest scattering strength.

After Giacalone & Jokipii (2005)

Again applying our shock-return criteria, the fastest ions attain normalized energy $\Delta\tilde{\varepsilon} \approx 2 \times 10^6$ in time $\Delta\tilde{t} = 5 \times 10^4$ whence $\sqrt{\Delta\tilde{\varepsilon}} / \Delta\tilde{t} \approx 3 \times 10^{-2}$, nearly two orders of magnitude too small. Note that, although the perturbation strength $(\Delta B/B)^2 = 0.1$ is only moderate, the previous figure indicates that we should expect no more than an order of magnitude increase in the flux, and perhaps less in the maximum velocity when perturbation strength is increased to unity.

Resonant Pitch Angle Scattering

Karimabadi et al (1992) addressed this problem beginning with a Hamiltonian approach in the presence of an imposed spectrum of waves and demonstrated mathematically and by resonance diagrams how diffusion may occur when resonances overlap. In such cases particles may hop from one H-surface to another, which they refer to as “second-order Fermi acceleration” with effective scattering dependent upon the particle resonance as well as the polarization of the waves with respect to the background magnetic field. This is illustrated in Figure 5-1 below where the panel (a) represents a single wave with two non-resonant and one resonant particle, while panels (b) and (c) depict a single particle with two waves at different polarization angles. The integration time for the first panel is $t = 2000 \Omega_c^{-1}$ while the other two are at $t = 6000 \Omega_c^{-1}$ and the total wave amplitude $\delta B = 0.4 B_0$ equivalent to an energy density $\delta\varepsilon = 0.16$ normalized to the background. We note that, although the latter panels show some diffusion, the velocity gain $\Delta v_{\perp} \sim v_A$ is minimal even after such a substantial elapsed time.

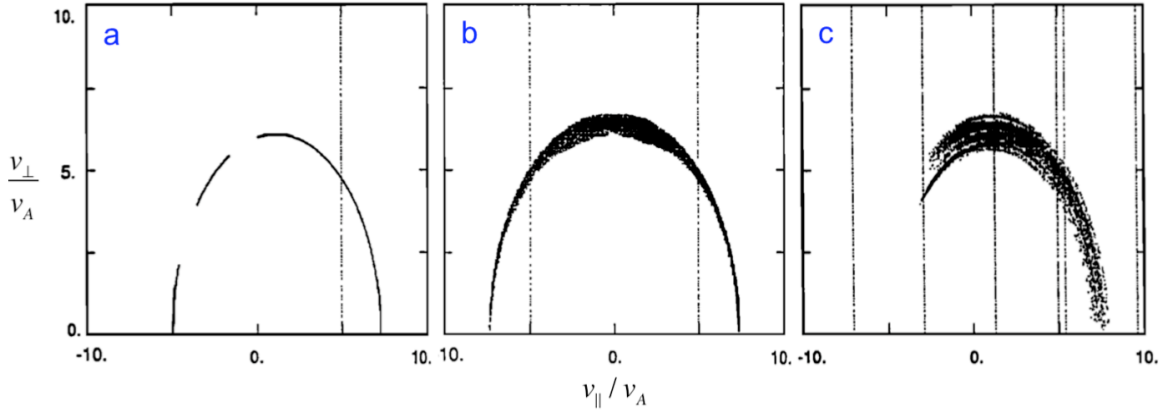


Figure 5-4 Velocity Diffusion in the presence of waves

Frequency $\omega = 0.3\Omega_c$ and amplitude $\delta B = 0.4B_0$ are the same for all waves. Panel (a) shows three particles of different initial velocities, and a single wave at angle $\alpha = 0^\circ$. Panels (b) and (c) are for two waves, the first with $\alpha_1 = 0^\circ$ and $\alpha_2 = 180^\circ$ and the latter $\alpha_1 = 0^\circ$ and $\alpha_2 = 30^\circ$.

After Karimabadi et al (1992)

Figure 5-5 below represents the same wave parameters as panel (b) above, except the wave angles $\alpha_1 = 30^\circ$ and $\alpha_2 = 210^\circ$. Here the acceleration is evident since the maximum attained perpendicular velocity $v_{\perp\max} \simeq 30v_A$. Although these studies were not performed in conjunction with a standing shock, we may assume as a rough benchmark that the downstream fluid speed of such shocks should be in the neighborhood of the local Alfvén velocity. For these cases, we recast our criteria more generally in open space with $u_2 \lesssim v_A$ for these citations. Hence, using $\tilde{v} \equiv v/v_A$, the minimum criteria for retrograde motion simply becomes $\dot{\tilde{v}} > 1$. Applying this criteria as before, we see that $\frac{\Delta v_\perp}{\Delta t} = \frac{30v_A}{6000\Omega_c^{-1}}$, or $\dot{\tilde{v}} = 5 \times 10^{-3}$, far too slowly to achieve retrograde motion. Extrapolating the wave energy to unity as observed *in situ* and modeled in other

studies cited herein, we may increase the diffusion rate linearly by a factor six which still gives $\dot{\tilde{v}} = 3 \times 10^{-2}$ still more than an order of magnitude too small.

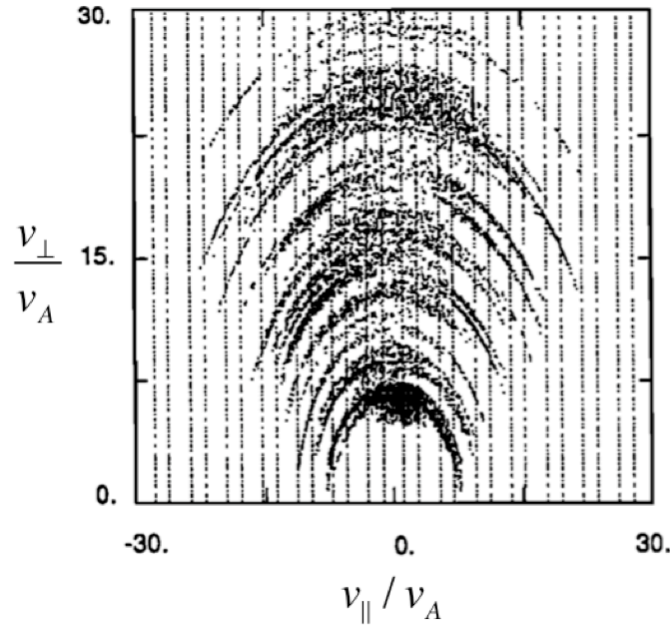


Figure 5-5 Enhanced diffusion by choice of wave angle
Wave amplitude and frequency are the same as the figure above, but the angles are now $\alpha_1 = 30^\circ$ and $\alpha_2 = 210^\circ$. After Karimabadi et al (1992)

Of particular interest to us in formulating an accurate scattering parameter is the authors' derivation of velocity diffusion coefficients and perpendicular diffusion rate, both of which include the factor $\Omega_C / v_\perp$ and the wave energy by way of the magnetic vector potential and wavenumber. We have already seen the latter dependence from the work of Giacalone & Jokipii, but here we explicitly have the ion mass and velocity dependence as well since $\frac{\Omega_C}{v_\perp} = \frac{eB}{mcv_\perp}$, so that the diffusion rate and therefore the scattering strength scale inversely with both ion mass and velocity.

Ion Cyclotron Velocity Diffusion

S. R. Cranmer (2001) has studied velocity diffusion in the solar corona, but with sufficient generalization in the underlying mathematical formalism to be more widely applicable. The main focus of the paper is to explain and rectify certain shortcomings of quasi-linear theory in predicting pitch-angle scattering. By way of the Vlasov equation, he derives a time dependent distribution function with parametric dependence on ion mass, charge and velocity, as well as the fluctuation spectral energy density. Numerical simulations were conducted with parameters appropriate to the solar corona, but we may adapt the results for our purposes. Figure 5-6 below shows the evolution of the velocity distribution of O^{5+} at four discrete times.

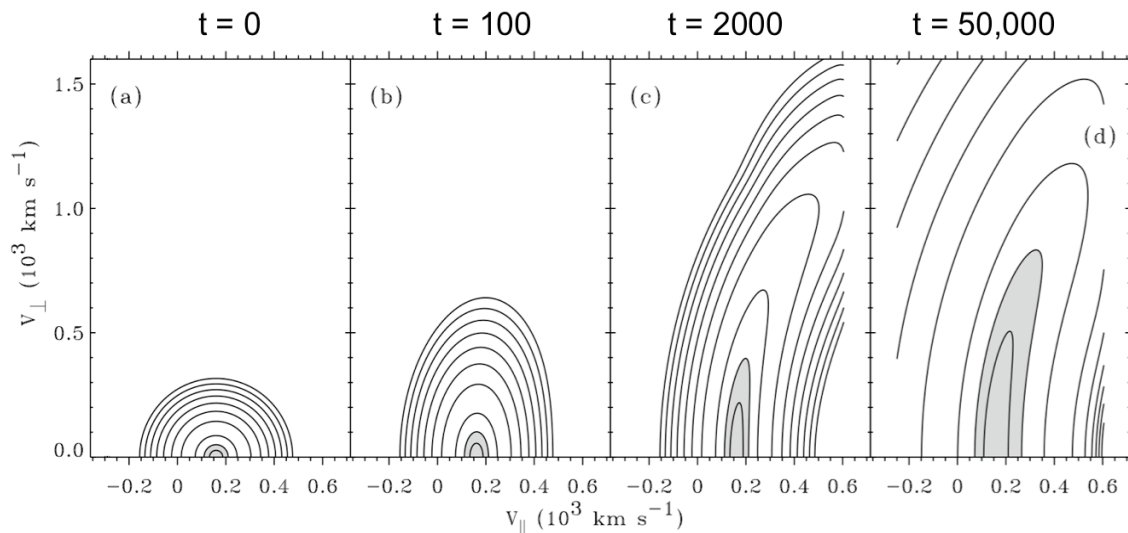


Figure 5-6 Velocity Distribution of O^{5+}

Contour plots at four elapsed times $t = 0, 100, 2000$, and $50,000$ seconds.

The charge/mass ratio is roughly equivalent to that of tritons H_3^+ and singly-ionized helium He_3^+ . After Cranmer (2001)

The Alfvén velocity in the region of the corona has been modeled at approximately $v_A \approx 2550 \text{ km s}^{-1}$, and the cyclotron frequency $\Omega_c = 7660 \text{ rad s}^{-1}$.

Normalizing the elapsed time and velocity maximum from the snapshots, we find that $\dot{v}(10^2 \text{ s}) \approx 1.8 \times 10^{-7}$, $\dot{v}(2 \times 10^3 \text{ s}) \approx 2.7 \times 10^{-8}$, and $\dot{v}(5 \times 10^4 \text{ s}) \approx 1.7 \times 10^{-9}$ and the maximum velocity asymptotes to $\dot{v}_{\text{max}}(t) \propto t^{-1}$. This also implies that $v_{\text{max}}(t) \propto \ln(t)$ or $\dot{v}_{\text{max}}(v) \propto e^{-Kv}$. The author further notes that, “at late times, the diffusion seems to stop before the extended core reaches the highest resonant values of $v_{\parallel}$ near the upper boundary of 610 km s^{-1} .” For the ions modeled, it is clear that the diffusion is insufficient to achieve retrograde motion. Extrapolating the results for O^{5+} to protons introduces a multiplicative factor $\frac{m}{q} = \frac{16}{5} \approx 3$ but the diffusion remains trivially small. Further, since the diffusion rate scales with wave power, in order to effect retrograde motion, the spectral energy density would have to be increased several orders of magnitude, well beyond any reasonable value.

Standing Waves

There are compelling arguments against standing waves of sufficient length scale and power enduring for more than brief periods at shocks inside the heliosphere. Bale, et al (2005) deduced that large amplitude waves downstream of the Bow Shock were not stationary in the shock frame, based on the lack of coherence in the downstream field profiles measured by the four Cluster spacecraft. Indeed, the variability of the shock itself would suggest this, and waves along the shock front itself were not observed at wavelengths greater than the spacecraft separation of around 600 km.

Nevertheless, for the sake of completeness, we performed a battery of simulations to study the effects of magnetic waves standing in the shock frame, and our results are given in the Appendix. We stress that the mechanism responsible for velocity distribution in this case is not strictly scattering, but simply the result of guiding center displacement due to a magnetic gradient drift. Even so, the results validate our intuitive predictions regarding the parametric dependence on mass, velocity, etc., insofar as they appear to agree at least in character with those observed in the turbulent diffusion studies cited.

Pitch Angle Scattering Model

We now seek a ‘kernel’ for either the Master Equation in the Fokker-Plank formalism or simply an estimate of the step size in velocity space, which will yield a reasonable estimate on the diffusion rate. We model pitch angle scattering by magnetic turbulence and show analytically that long wavelengths are required for effective velocity diffusion and derive a velocity ceiling due to wavelength constraints. We do not address the connection between turbulence and $\lambda_{\parallel}$, nor do we discuss mixing between parallel and perpendicular velocities since those diffusion coefficients have already been dealt with in great depth and, once $v_{\perp}$ is large enough, any gain through scattering $v_{\parallel} \rightarrow v_{\perp}$ should yield only small contributions.

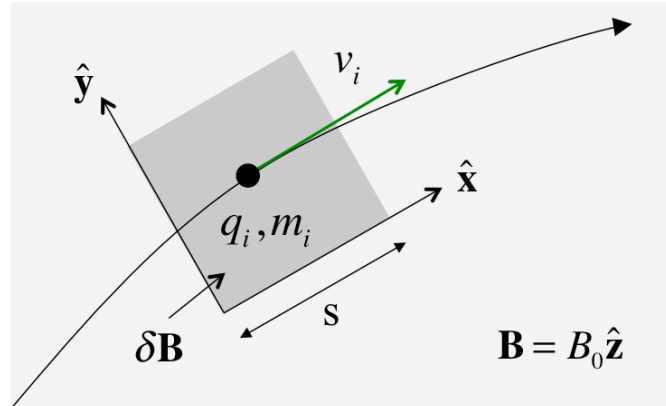


Figure 5-7 Scattering by a 'cell' of magnetic turbulence

Consider an ion of mass m_i and charge q_i traveling with velocity v_i on a guiding center trajectory in a region with ambient magnetic field $\mathbf{B}$ as shown in Figure 5-7. In the presence of magnetic turbulence, we may model the region as a patchwork of disjoint compact cells each of finite extent s through which the particle must pass. Granted that such turbulence is composed of an ensemble of waves over a broad spectrum of frequencies and wavelengths, we nevertheless claim the following. Within the cell, the effect from turbulent spectrum due to higher wavenumber components should stochastically 'wash out' and leave only the longest wavelengths to do the scattering. This can be seen by noting that, over a compact region, any waveform can be Fourier decomposed and whether the magnetic field components are coherent waves or randomly oriented, averaging the contributions over the path through the cell should leave only the 'DC' component as a constant 'bias' field. We can also offer justification for treating the turbulent waves as time-stationary or frozen within the cell if the particle's transit time is shorter than the shortest relevant wave period – say for Lower Hybrid

wave. Certainly we can ignore the electrostatic waves, ω_{pe} and ω_{pi} and simply require that $v_i > \Omega_{LH}s$.

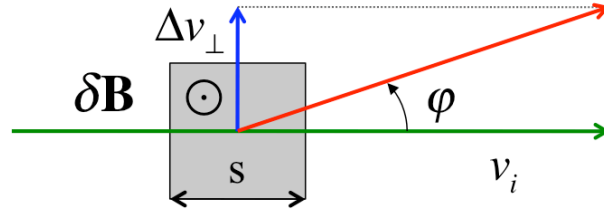


Figure 5-8 Impulse through a turbulent cell

In Figure 5-8 we show an ion passing through a region in space of extent s where the magnetic field perturbation $\delta\mathbf{B}$ is nearly uniform, and we have rotated the coordinate system so that $\delta\mathbf{B}$ points out of the page. The unperturbed path will be nearly straight provided $r_g \gg s$ which implies either a small cell size or high velocity v_i . We can then compute the perpendicular velocity component resulting from the impulse through this region by ‘integrating’ over the unperturbed path

$$\Delta v_{\perp} = a_{\perp} \tau = \frac{F_{\mathbf{B}}}{m} \tau = \left(\frac{q}{m} \frac{v}{c} B \right) \frac{s}{v} = \Omega_{ci} (\delta\mathbf{B}) s$$

which is independent of the incident ion velocity, and approximate the scattering angle at high velocity for a given field strength $\delta\mathbf{B}$, by $\tan \varphi \approx \varphi = \Delta v_{\perp} / v = \Omega_{ci} (\delta\mathbf{B}) s / v$. Note that the scattering angle φ is proportional to the cell size s which, in keeping with our assumption that δB be nearly uniform within the cell, can be no larger than the wavelength of the perturbation $s \sim \lambda_{\delta B} \sim k_{\delta B}^{-1}$. This should explain the wavelength dependence of the velocity ceiling as observed by Giacalone, Ellison and Jokipii.

Further, the scattering angle is inversely proportional to the ion velocity, again consistent with the observed due to an ensemble of waves of fixed strength and spectral dependence, but limited below a maximum wavelength. Finally, we assume that the integrated power in magnetic turbulence, both observed and assumed for simulations, is of order unity with respect to the background. Hence we are justified, within an order of magnitude approximation, in further taking $\Omega_{ci}(\delta\mathbf{B}) \sim \Omega_{ci}(\mathbf{B}_0)$ where $\mathbf{B}_0$ is the background magnetic field strength.

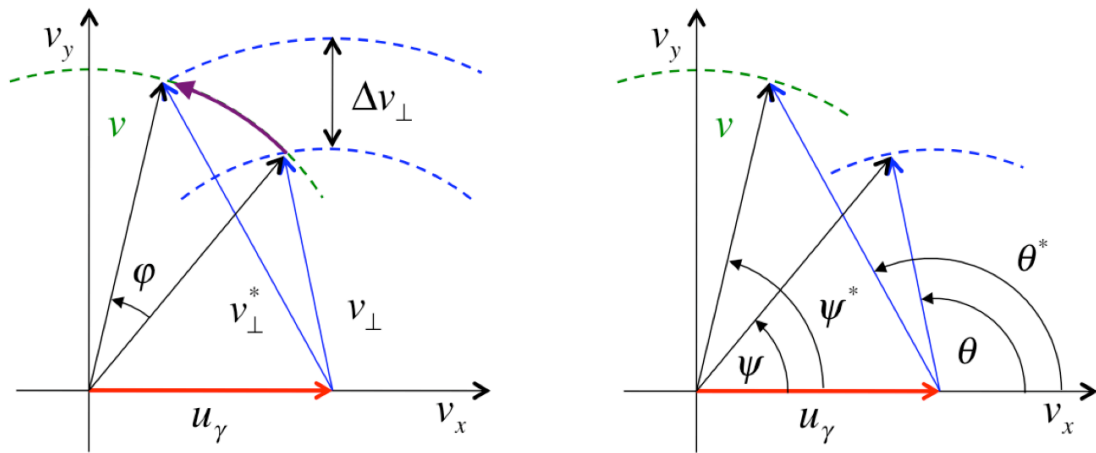


Figure 5-9 Pitch Angle Scattering

The blue rays represent pre- and post scattering velocities in the guiding center frame while the black rays represent the same velocities in the magnetic wave frame, offset by the wave speed u_γ .

To complete the model, we now consider a construction of pitch angle scattering in the plane of an ion gyrating about its guiding center as shown in Figure 5-9. Note that strictly magnetic waves standing in the fluid frame cannot effect a change in energy, but merely rotate the gyro-phase angle leaving the ion cyclotron velocity $v_\perp$ invariant. We

therefore assume that the scattering is caused by waves traveling at velocity u_γ with respect to the guiding center or fluid velocity. Since the magnitude of the ion cyclotron velocity in the wave frame is invariant, we equate the velocities in the two frames both pre-and post-scattering to obtain $ve^{i\psi} = v_\perp e^{i\theta} + u_\gamma$ and $ve^{i\psi^*} = v_\perp^* e^{i\theta^*} + u_\gamma$. Here $v_\perp$ and $v_\perp^*$ are the perpendicular cyclotron velocities in the guiding center frame, before and after scattering, θ and ψ are the gyro-phase angles in the guiding center and waves frames respectively. The ion velocity magnitude v is invariant in the wave frame, which moves with respect the guiding center at speed u_γ . The left panel shows a scattering event in the wave frame that causes a change in the guiding center cyclotron velocity, which may be viewed as a ‘kick’ into a higher orbit. Dividing through by the relative wave speed, we have $\tilde{v}e^{i\psi} = \tilde{v}_\perp e^{i\theta} + 1$ and $\tilde{v}e^{i\psi^*} = \tilde{v}_\perp^* e^{i\theta^*} + 1$ which, upon dividing the two equations gives $e^{i\varphi} = \frac{\tilde{v}_\perp^* e^{i\theta^*} + 1}{\tilde{v}_\perp e^{i\theta} + 1}$. Letting the ion perpendicular velocity grow without bound implies that $e^{i\varphi} \rightarrow \frac{\tilde{v}_\perp^*}{\tilde{v}_\perp}$ as $v_\perp \rightarrow \infty$ or equivalently, $\Delta v_\perp \rightarrow 0$ unless the scattering angle φ in the wave frame becomes extremely large. Indeed, the incremental velocity change per ‘collision’ is limited to the wave velocity u_γ and can only be achieved at high ion velocities by scattering through an angle $\varphi = \pi$. However, we have already shown above that this angle necessarily diminishes with increasing ion velocity, all other factors being equal, and thus conclude that scattering by magnetic waves alone is fundamentally limited in promoting ions to extremely high velocities.

The DSA Scattering Parameter

Having established that scattering strength must be proportional to the wave

energy density $\left| \frac{\delta B}{B} \right|^2$ and inversely proportional to the wavenumber k , mass to charge

ratio m_i / q_i and ion velocity v_i , we now seek a functional form for the scattering

strength and the parameter η used in the Webb model and applied in many subsequent

treatments. We also show that the scattering parameter in the Webb model is inverted in

its dependence on the ion gyroradius which contains independent variables that we have

shown are germane to the problem. We naturally require that, for all parametric

dependence, the function properly converge in all their limits, which can only be true for

a product of factors. Since velocity diffusion is more pronounced at longer wavelengths,

the spectrum – not just the integrated power – must be taken into account. Consider the

differential element $\frac{\partial \delta B(k)}{\partial k} dk$. We may include the spectral dependence by integrating

the elements over the entire spectrum in order to obtain the weighted mean

$\langle k \rangle = \frac{1}{\delta B} \int k \frac{\partial \delta B(k)}{\partial k} dk$ where $\delta B = \int \frac{\partial \delta B(k)}{\partial k} dk$. But we can simplify this by noting

bow shock observations indicate a power spectrum exponent $p = -1.9$, while fully

developed Kolmogorov turbulence has $p = -5/3$. Assuming that the spectrum near a

shock is sufficiently hard (unless the turbulent bandwidth is exceedingly narrow) we may

safely take $\langle k \rangle = k_{\min}$ and apply the total turbulent energy at that wavenumber. As long

as the exponent $p > 1$ and the bandwidth is greater than a decade $\lambda_{\max} / \lambda_{\min} > 10$,

$$\int_{k_{\min}}^{k_{\max}} k^{-p} dk = \frac{1}{1-p} \left| k_{\max}^{1-p} - k_{\min}^{1-p} \right| = \frac{1}{p-1} k_{\min}^{1-p} \left[1 - \left(\frac{k_{\min}}{k_{\max}} \right)^{p-1} \right] \approx \frac{1}{p-1} k_{\min}^{1-p}.$$

We have also shown that scattering is dependent inversely upon an ion's momentum and, via the Lorentz force, directly proportional to charge so that, in the non-relativistic regime, we must include the factor $m_i v_i q_i^{-1}$ where m_i , q_i and v_i are the ion mass, charge and velocity respectively. The ion's instantaneous gyroradius r_{gi} will do nicely, since its product with the wavenumber is dimensionless, providing the ion velocity is sufficiently greater than the ambient fluid speed so that $v_i \approx v_{\perp}$.

Recall that we are not seeking a statistical ensemble solution, but rather a specific velocity required for *any* individual particle to be injected. Since η by definition of use is inversely proportional to the scattering strength, we construct a new downstream dimensionless parameter with explicit velocity dependence

$$\eta_2 = \langle k \rangle r_{ci} (v_i) \left(\frac{B}{\delta B} \right)^2 = \langle k \rangle v_i \Omega_2^{-1} \left(\frac{B}{\delta B} \right)^2 = k_{\min} v_i Z_i \frac{m_i}{m_p} \Omega_p^{-1} R_C^{-1} \left(\frac{B}{\delta B} \right)^2$$

where $\Omega_{1,2}$ are the upstream and downstream ion cyclotron frequencies and Z_i the atomic number. Our function contains the proper pitch angle scattering dependence on perturbation strength, wavelength weighting, particle velocity, charge and mass. The implicit Webb downstream injection criterion

$$\frac{v_{inj}}{u_1} \gg \frac{4}{R_C} \sqrt{1 + \eta_2^2 (v_i, m_i, q_i, \delta B / B)}$$

may then be rewritten explicitly in velocity as

$$v_{inj} \gg \frac{K_1}{1 - K_1^2 K_2^2} \text{ where } K_1 = \frac{4u_1}{R_C} \text{ and } K_2 = \frac{k_{\min}}{\Omega_1 R_C} \left(\frac{B_0}{\delta B} \right)^2.$$

Although this still contains the ambiguous ‘much much greater than’, we may approximate some limits of this function, at least within the Heliosphere. Clearly, a solution exists only when $K_1 K_2 < 1$ and a reasonably attainable injection velocity would require $K_1 K_2 \ll 1$. However, restricting ourselves for the moment to protons at the Earth’s bow shock, roughly Mach 5 and compression ratio $R_C \sim 2.6$ with reasonably strong turbulence of order unity, we find that $k \simeq 9.8 \times 10^{-3} \text{ km}^{-1}$ simply to obtain a solution. This gives an absolute floor on the turbulent wavelength of $\lambda_{\delta B}(\text{min}) \simeq 3700 \text{ km} \simeq 37 \lambda_{pi}$ simply to engender a finite velocity threshold at all. Even disregarding the ‘much much greater’ asserted in the Webb model would require wavelengths far greater than the scale of any planetary bow shock to bring such a threshold within realistic range of ion velocities observed.

Summary & Conclusions

We have shown that velocity diffusion by turbulent wave scattering is a function of the particle velocity, charge to mass ratio, turbulent energy density and wavelength. We have noted that a velocity ceiling exists whenever the maximum wavelength is constrained by spatial scales, especially those at planetary bow shocks. Outside the Termination Shock, these constraints may be relaxed, but we have demonstrated that, since diffusion slows with increasing particle velocity, even the strongest turbulence may be unable to return ions to the shock front against the downstream fluid flow. Further, the fact that such diffusion is more effective at lower velocities casts doubt on the need for a minimum injection velocity for any process including DSA. Indeed, the slowing growth rate in coupled with wavelength constraints calls into question the ability of any diffusion-

based acceleration mechanism to promote ions to the relativistic energies observed in the hard cosmic ray spectrum.

On the other hand, observations indicate that turbulence in the vicinity of a shock possesses spectra sufficiently hard to concentrate the majority of the energy at the longest wavelengths, and we have shown that these wavelengths are minimally destructive to the surfing mechanism. However, the inverse velocity dependence of the diffusion rate calls into question any advantage which may be gained through promotion by shock surfing, either as a pre-injection mechanism or simply an aid in facilitating shock front return.

We shall now examine case-by-case studies from observations of shocks in the heliosphere, and apply our derived shock surfing criteria along with results of diffusion studies to determine whether this mechanism warrants further consideration as a viable contribution anywhere within the heliosphere, and pay particular attention to recent observations of the Termination Shock.

Chapter 6 The Heliosphere

We have shown that velocity diffusion alone cannot return ions to quasi-perpendicular shock fronts unless the turbulent scattering strength significantly exceeds the background energy density, or the scale of the shock structure is large enough to support the extremely long wavelengths required for effective pitch-angle scattering. While the magnetospheres of outer planets, especially Jupiter, have scales two orders of magnitude larger than Earth's, the scales of interest are comparable, since the quasi-perpendicular region of the Earth's bow shock lies along the flank, while that of the outer planets resides at the nose. Hence the temporal, geometric and scale constraints at planetary bow shocks rule out Diffusive Shock Acceleration there.

The simulations of Giacalone, Ellison, Jokipii, Karimabadi and others has also established that the diffusion rate slows with increasing ion velocity and therefore in time, abrogating any advantage gained by a velocity boost via Shock Surfing (SS). Further, the similarity between the velocity distributions of boosted thermal ions and pickup ions strongly suggests that any mechanism that relies on velocity diffusion alone as the *primum mobile* may require no minimum threshold at all. Hence, assuming realistic values of turbulent energy that we justify below, we shall dismiss SS as an injection mechanism for any other possible form of diffusive acceleration at planetary bow shocks.

Nevertheless, SS may possibly play some other role by promoting ions to high velocities at perpendicular shocks. We now investigate as case studies, the data from several missions in the heliosphere in order to determine whether the surfing mechanism can function well enough to explain any of the energetic ion phenomena observed *in situ* near shocks there.

Review

We have shown that several factors can suppress the SS mechanism by limiting both the maximum attainable velocity and the ‘efficiency’ as defined by the flux of surfing ions promoted above the Shock Drift (SD) velocity ceiling. In this chapter, we examine observational data as case studies to determine whether SS may in fact be viable at shocks in the heliosphere and beyond. Specifically, we have demonstrated that:

- A velocity ceiling exists, which is limited by the strength and ramp width of the electrostatic potential. When the Lorentz force exceeds the retarding force of the shock-normal electric field, surfing ions escape downstream. The critical ramp width at which surfing ions can no longer exceed the SD velocity ceiling is roughly $L_{crit} \sim 0.1 M_A \lambda_{pi}$, applicable over a fairly broad range of shock parameters.
- Front corrugations prematurely terminate the surfing process when the ion cyclotron radius (as a function of velocity) exceeds the local radius of curvature. We have shown in test particle simulations that the velocity ceiling scales inversely with ion mass m_i , perturbation amplitude a and inversely as the square of the wavenumber k . Surfing ‘efficiency’ as previously defined is also suppressed, and we have derived analytic functions to describe both effects. Rippled shock fronts have not only appeared in fully self-consistent simulations, but also observed *in situ* at planetary shocks. We shall apply our suppression functions to that data to see whether SS can remain viable in the presence of these perturbations.

- The surfing velocity ceiling is severely limited by shock reformation time scales. Numerical simulations and *in situ* observations indicate that perpendicular shocks overturn and reform on the scale of an ion cyclotron period, which limits the acceleration time and therefore the maximum velocity attainable to roughly that of the shock drift mechanism. We will show that this periodic reformation is characteristic of all heliospheric shocks, and that the time scale is sufficient to render SS ineffective.
- The velocity ceiling is constrained by the obliquity of the magnetic field, which limits the regions of planetary bow shocks that may otherwise support shock surfing, and by temporal constraints due to solar wind advection along the flanks. Besides the shock geometry and Parker angle, the cross-shock potential and ramp width are function of the local SW angle of attack i.e., the normal velocity component, which becomes oblique along the shock flanks weakening the surfing mechanism.

We now turn our attention to case studies in the heliosphere and compare the observational data with our previously derived criteria in order to determine whether any of these constraints are sufficient to render surfing untenable. We will show, case by case, that shock surfing is not a viable mechanism anywhere in the heliosphere and, where it can possibly function at all, is incapable of promoting ions to velocities substantially greater than those achieved by SD.

The Inner Planets

One feature that distinguishes the inner from the outer planets is the Parker angle, which steepens with radial distance from the sun, becoming nearly perpendicular $\theta_p \rightarrow 90^\circ$ at the Termination Shock. As shown in Figure 6-1 below, the steep Parker angle at inner planets places the quasi-perpendicular region down the flanks of the shock. This is unfavorable for surfing since upstream fluid advection along the flank quickly sweeps surfing ions out of range. At an even more fundamental level, the off-normal fluid flow reduces the cross shock potential (the compression ratio is smaller along the flanks) and the convective electric field $\mathbf{E}_c = -\mathbf{u} \times \mathbf{B} / c$ (responsible for acceleration) that diminishes as the sine of the angle between the magnetic field and the local fluid velocity vector.

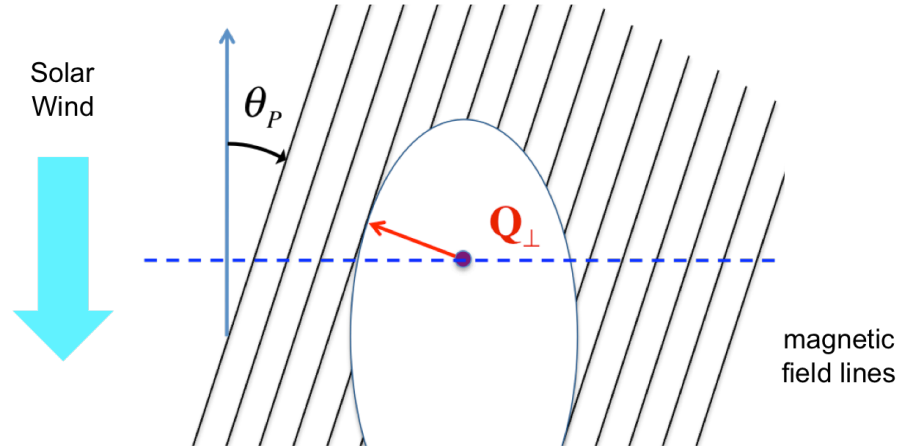


Figure 6-1 Parker Angle at an Inner Planet

The shallow angle places the quasi-perpendicular region (red arrow) along the flank. The dashed line represents the orbital trajectory.

Earth is the only planet in the inner Heliosphere with a substantial magnetic field. The others are not magnetized, but still possess 'induced' magnetospheres. However,

these are too small and weak to support shock surfing. In addition, the Parker angle at the innermost planets, Mercury and Venus, puts the perpendicular region too far down their flanks to even be considered. Cometary bow shocks are also dismissible for the same reason.

Coronal Mass Ejections

CMEs originate from active solar regions with broad latitudinal distribution during solar maxima, but primarily in the coronal streamer belt near the magnetic equator during solar minima. They initially range in velocity from a fraction to several times the SW velocity. However, through interaction with the solar wind magnetic field, they eventually either slow or accelerate to v_{SW} . Initially supersonic CMEs obviously form shocks, which expand and weaken with time, as they lack a source to sustain them indefinitely. They are believed to form as near as 2 solar radii and are closely linked with the acceleration of Solar Energetic Particles, so there is no question that some acceleration mechanism is responsible. But, at such close distances, the Parker angle must still be nearly parallel at the nose and any ions boosted by surfing along the flanks must necessarily be left behind as the finite shock structure propagates outward. By the time the CME itself reaches 1 AU (one to five days) where the Parker angle is more favorable, it has generally slowed to the SW velocity (Gonzalez-Esparza 2005) and no longer viable for shock surfing.

Co-rotating Interactive Regions

CIRs are regions of enhanced magnetic field bounded by jumps in the solar wind speed that form at distances from the sun greater than 2.5 AU. As with other shocks, both propagating and standing, they have been shown to give rise to waves and

accelerated particles. However, no associated electrostatic potential gradient has been proposed or observed and, as this a key requirement for shock surfing, these regions are unlikely candidates for this mechanism.

The Earth's Bow Shock

Earth's Magnetosphere has, for the obvious reason of proximity, been studied the longest and in the greatest depth. Beginning with Explorer 1 in 1958, numerous probes – including dozens of 'Explorer' missions alone – have been launched carrying a variety of instrumentation that continued to provide valuable data. In 1964 IMP-1 provided the first evidence of the bow shock, trapped O^+ from the ionosphere was found in 1977 and direct measurement of its origin at the poles and upward transport observed by the USAF probe S3-3 in 1977.

The Cluster Mission was launched in 2000, and consists of four identical and roughly equidistant spacecraft in eccentric polar Earth orbits whose separation was altered as required for different mission objectives. Unlike previous single or even double probes, this multi-craft mission made possible for the first time high resolution spatial and temporal measurements of the Earth's bow shock as the probes executed repeated front crossings.

Large Scale Morphology

Many general features of the magnetosphere and the bow shock in particular have been known for some time prior to the Cluster Mission. Indeed, the high-resolution data recovered merely served in some cases to reaffirm these attributes. The overall shape and scale with locations of the quasi-parallel and quasi-perpendicular regions as

shown in Figure 6-2 were consistent with previous observations and theoretical models based upon them.

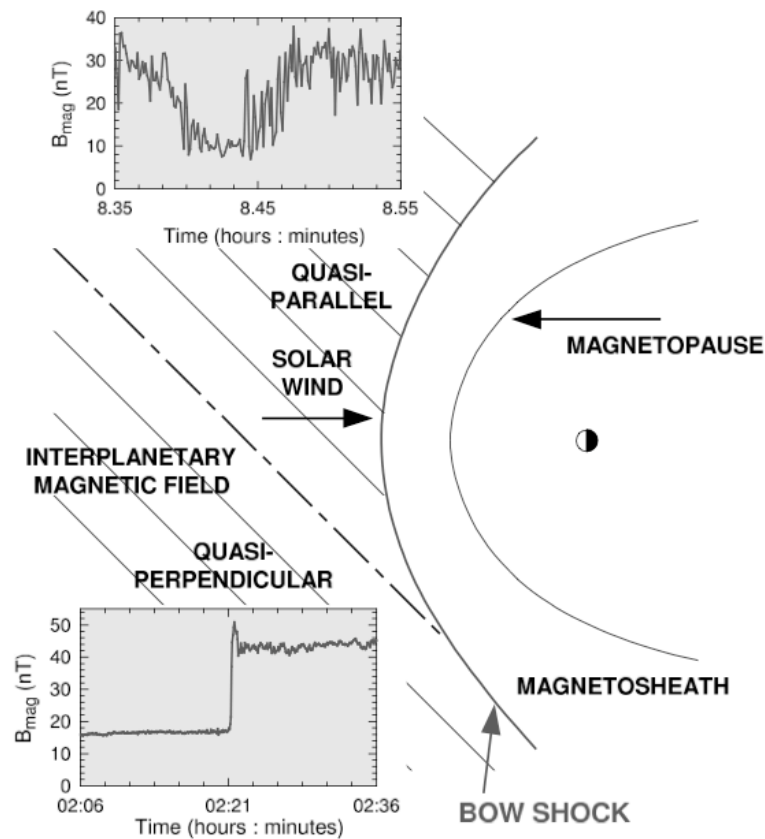


Figure 6-2 Sketch of the Earth's Bow Shock

The Solar Wind flows from the sun to the right. Typical shock-crossing magnetic field profiles are shown for perpendicular and parallel regions. The large-scale structure is remarkably stable while the standoff distance to the nose responds to temporal changes in the solar wind pressure.

After Balogh et al (2005)

The solar wind is not constant in time and, in addition to abrupt changes in turbulent wave energy due to violent solar events, the mean speed and pressure change in response to solar variation. Nevertheless, Horbury et al (2002) discuss the orientation and motion of the quasi-perpendicular bow shock in particular from data gathered during

the first two years of the Cluster mission when the spacecraft were separated by roughly 600 km. They found “the bow shock normal to be remarkably stable, under a wide range of steady upstream conditions” and essentially in keeping with previous observations and theoretical models. Under steady solar wind conditions in particular, they found the bow shock remains close to a parabolic shape, in agreement with the model of Peredo et al (1995).

However, it is the fine features that are of interest to our investigation of shock surfing. Cluster analysis was based upon selection of 48 “clean and steady” quasi-perpendicular front crossings, which represent a mere 1% of the data. Since the authors considered only data from reasonably clean crossings and rejected those with significantly different magnetic field profiles at different spacecraft (representing shocks in motion or reformation), this suggests that such “clean and steady” moments may be rare. Indeed, while these samples may be sufficient to establish the general morphology of the shock on large scales, much of the tractable crossing data indicate a substantial level of activity and variability at the smaller scales of interest to us. We show in Figure 6-3 below, three examples drawn by Horbury et al (2002) from the 48 “clean” crossings that highlight this variability.

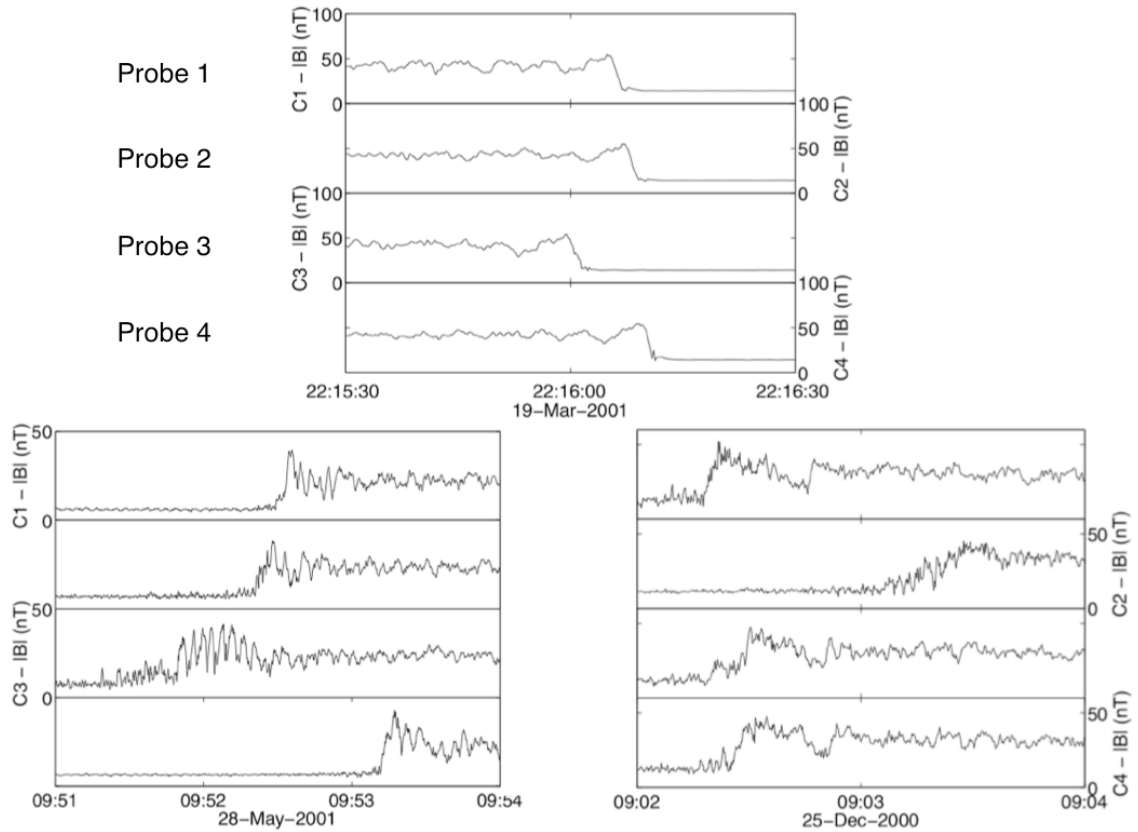


Figure 6-3 Observed Magnetic Field profiles

Top panel shows a sharp shock, bottom left a disturbed shock and bottom right one which was likely reforming. The field magnitude at each of the four spacecraft is shown in each panel from top to bottom, Cluster 1 to 4
After Horbury et al (2002)

Shock Magnetic Field Features

Before presenting the Cluster data and analyses of its shock crossings, we wish to emphasize one absolutely essential point. There are two unique features at the perpendicular shock front – the magnetic field gradient on the one hand, and the electrostatic potential gradient on the other. Each, aside from its detailed structure, may be characterized by magnitude and length scales.

Aside from fine structural features, two spatial gradients are of primary interest for surfing analysis – the shock ramp possessing the steepest field gradient, and the ‘foot’ or shallow precursor region ahead of the ramp. Although the scale of the magnetic ramp gradient has generally been estimated to be on the order of the upstream ion inertial length $\lambda_{pi} \equiv c / \omega_{pi}$, many alternatives have been proposed. From observations, Balikhin et al (1995) and references therein assert the ramp scale to be less the ion inertial length, with reports of one or two shocks of order $0.1c / \omega_{pi}$, after Newbury and Russell (1996) and Walker et al (1999).

From early theory and simulations, Leroy et al (1982) suggested that the magnetic ramp occurs on either an ion inertial scale λ_{pi} or the gyro-radius of an ion moving at the upstream flow velocity in the downstream magnetic field. The latter represents the downstream transit of the fastest incident pickup ions. It is most revealing by expressing the upstream gyroradius in units normalized to the local plasma parameters

$$\tilde{r}_g = \frac{r_g}{\lambda_{pi}} = \frac{v_{\perp}}{V_A} \Omega_c^{-1} \lambda_{pi}^{-1} V_A = M_A \frac{mc}{eB} \sqrt{\frac{4\pi ne^2}{mc^2}} \sqrt{\frac{B^2}{4\pi nm}} = M_A$$

whence the downstream radius for these same ions may be obtained simply by dividing r_g by the shock compression ratio $R_C = B_2 / B_1$ i.e., $r_{g2} = M_A R_C^{-1}$. At the Earth’s bow shock, this may be anywhere from slightly greater than unity to about three times the ion inertial length. The relationship between scaling of ion inertial length with Mach number was confirmed from Cluster observations by Bale et al (2002), as demonstrated in the lower panel of Figure 6-4 below. The upper panel suggests that the cross-shock scale (shock width) when normalized to the local ion inertial length remains roughly invariant.

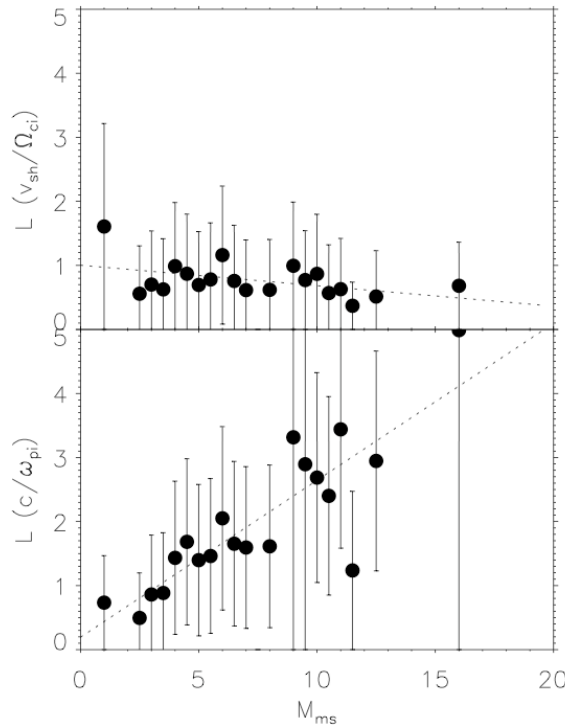


Figure 6-4 Inertial Scale versus Mach number

Top panel shows $L / (u_{sw} \Omega_{ci}^{-1})$ is approximately constant, while the bottom panel shows ion inertial scale increases with Mach number.

After Bale et al (2003)

The width of the foot was estimated by Sckopke et al (1983) and Livesey et al (1984) to be of order $L_{foot} \sim 0.68 u_{sw} \Omega_{ci}^{-1}$ or equivalently $\tilde{L}_{foot} \sim 0.68 M_A$ in units normalized to the local ion inertial length, which is not substantially different from other estimates. On the other hand, there appears to be some question as to whether the scales discussed in the literature actually refer to the foot region or to the magnetic shock ramp itself.

Cross Shock Potential

Although the plasma density tends to follow the magnetic field (Scudder et al., 1986), the electric field has shown finer features. While the scale and structure of the magnetic component of the shock are of intrinsic interest to other investigators – although the foot region may, in fact, play a role in limiting surfing efficiency through pre-expansion of the pickup ion shell – it is the electric field in the ramp that is key for our purposes. Since shock surfing relies entirely upon the shock-normal electric field to sequester pickup ions at the front, determination of the strength, scale and features of the cross-shock potential is of paramount importance in determining the feasibility of this mechanism in the light of observational data.

The difference in energy gain between a few electrostatic reflections and that of ‘Multiply Reflected Ions’ is vast. Indeed, it was the high energy attainable through multiple reflections that made the concept of shock surfing so attractive in the first place, and this difference intimately depends upon the ramp electric field scale being on either that of the ion or electron inertial length.

Many treatments & discussions have been approached under the assumption that the magnetic and electric field profiles at perpendicular shocks are, at least in general character, similar to those proposed in the most elementary pedagogical models. In particular, both the physical coincidence and structural similarity of the field in the transition region – the foot, ramp and overshoot – has sometimes been presumed. We show that detailed observations of the fine field structure do not support these assumptions.

One ubiquitous feature of all the exploratory missions is the relative ease and accuracy with which the magnetic field has been measured when compared with the

electric field. Due perhaps to the lack of highly resolved measurements from early missions, observations of the electric field were very sparse. Some early theoretical treatments predicted fundamentally different scales for the electrostatic versus magnetic field components of the quasi-perpendicular shock ramp, and the apparent difficulty in measurement has prolonged the debate.

For example, Formisano (1982) reported from ISEE data that the electric field ramp observed at quasi-perpendicular shocks lasted longer than the magnetic ramp crossing itself, while numerical simulations by Lembège et al (1999) showed that the scale lengths were of the same order. Other authors have also asserted that the spatial scales of the potential and magnetic field in the ramp region are similar while Scudder (1995) proposed the potential scale length should be larger than that of the magnetic scale length. More recently, Scholer et al. (2003) showed that the main potential drop can occur over several ion scales in the foot region, but also estimated that the ramp region also significantly contributes to the total potential over much smaller scales, typically $5 - 10 \lambda_{pe}$, roughly corresponding to the range $0.12 < \tilde{L}_r < 0.23$ in ion inertial length units.

Among the early estimates were based on initial results from ISEE by Heppner et al. (1978) who reported observation of short duration spike-like electric field features. In light of the wide range of views upon this subject, a possible point of confusion in the past may have been the misinterpretation of the 'spike' scale with that of the entire cross-shock potential gradient itself.

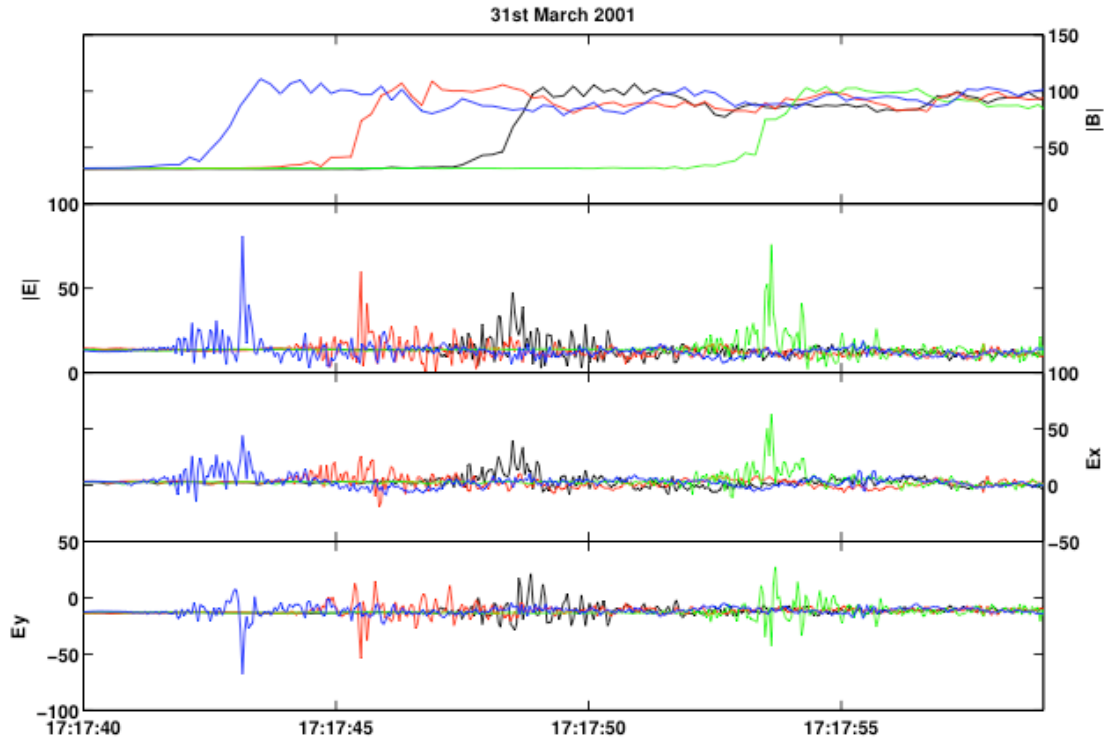


Figure 6-5 Shock Crossing Field profiles measured by Cluster Data from four spacecraft are superimposed. The time delay is due to the spacing of the Cluster probes. Magnetic field (top) Electric field (second) and its components E_x and E_y in the satellite spin-plane are shown.

After Walker et al (2004) as cited in Bale et al (2007)

With the higher resolution data provided by the Cluster mission, we are now in a better position to resolve this issue. As shown in Figure 6-5 above, large amplitude spikes were observed in the electric field through the shock, but at very small scales on the order $1 - 5\lambda_{pe}$. Roughly half of the entire cross-shock potential occurred within these small regions of the shock. However, two features are clear from the figure. The first is that these electric field spikes are distributed over regions extending beyond and in some cases well beyond the scale of the magnetic ramp. The trace from Cluster Probe 4

(green) extends upstream beyond the foot region. The second feature of note is that, even in the short time elapsed between individual probe crossings the shape of the magnetic structure changes noticeably, acquiring a second step in the magnetic field profile. These features are accompanied by substantial changes in the electric field profiles, as can be seen from the heights and lengths of the spike trains. Note particularly the apparent concentration in the Cluster 4 trace (blue) as compared with the spread through same shock passage for Cluster 2 (black). The elapsed time between them is less than five seconds, or about two ion cyclotron periods. Finally, it is clear from the changing juxtaposition between the magnetic and electric field profiles that the largest potential step does not always coincide with the magnetic ramp spatial position.

During the early stages of our own investigation, a brief suite of Test Particle simulations was performed with multiple potential step profiles rather than a single uniform ramp field. We concluded that such a disjoint set of potential steps, especially when spread over scales approaching an ion inertial length, substantially reduced the ability of the shock front to restrain ions upstream. In other words, we deduce that in the absence of a nearly uniform and sustained cross-shock electric field, the surfing mechanism ceases to be effective at all. The particular set of Cluster measurements we have just described was taken during a shock crossing associated with the passage of a coronal mass ejection (CME) where the local upstream plasma parameters were $\mathbf{B} = 30nT$, $\theta = 87^\circ$, $n = 19cm^{-3}$ and $M_A = 3.6$, and a histogram of the scale size distribution is shown below in Figure 6-6. However, the authors also note that similar large amplitude, short duration features in the electric field have been observed at a number of shock crossings. A scatter plot of those distributions as a function of shock Mach number is shown in Figure 6-7.

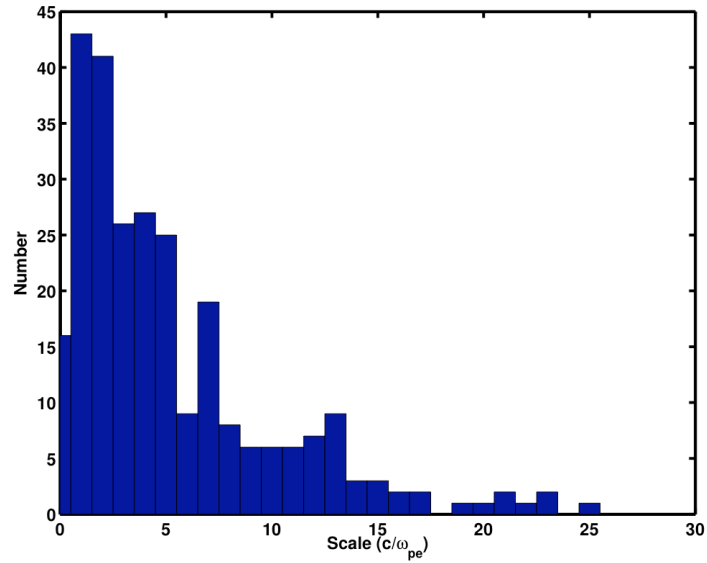


Figure 6-6 Distribution of spike-like electric field features
The histogram shows counts versus electron inertial length
After Walker et al (2004) as cited in Bale et al (2007)

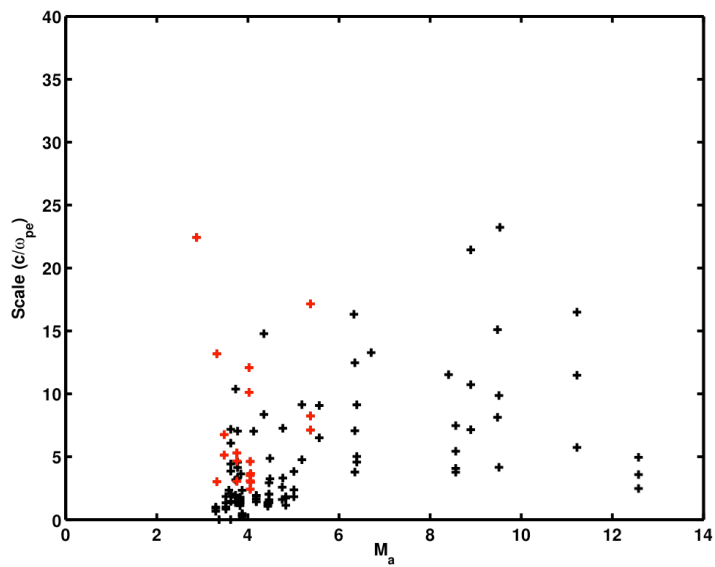


Figure 6-7 Electric field scale versus Mach number
The data in red correspond to the CME event previously shown.
Balikhin et al (2002) cited in Walker et al (2004)

Finally, we show in Figure 6-8 the distribution of the electric field ‘spike’ scales as a function of shock obliquity, and note that the distribution broadens considerably with deviation from the perpendicular. Our test particle simulations showed considerable velocity suppression as the magnetic field angle deviated from the perpendicular, but with the ramp width held constant. Here we see that the electric field scales are coupled to that obliquity as well, which places an additional constraint on regions of the bow shock otherwise theoretically capable of supporting shock surfing.

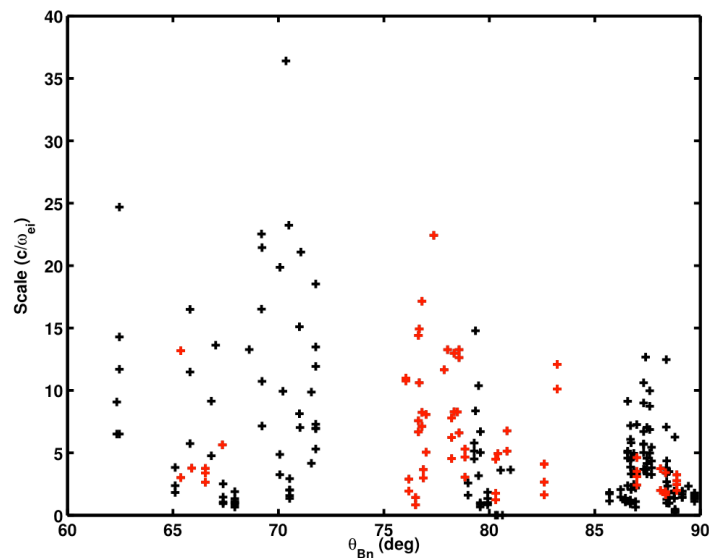


Figure 6-8 Electric Field scale vs. magnetic field angle
After Walker et al (2004) as cited in Bale et al (2007)

In summary, we conclude that although observations show very small scale spike-like features in the electric field – on the order of a few electron inertial lengths – a substantial portion of the cross-shock potential gain occurs over scales on the order of an ion inertial length, and sometimes much larger. Indeed, in those cases where the gradient extends into the magnetic foot region, we may characterize the field as

extending over several ion inertial lengths. This is sufficient to severely disable the surfing mechanism entirely, leaving only singly reflected ions to contribute to any observed upstream gyrating population.

Shock Variability

It is now widely believed that cyclic reformation is a ubiquitous feature of quasi-perpendicular collisionless shocks. Citations include evidence from numerous observations as well as from numerical simulations. Prior to Cluster, oscillations on the order of an ion cyclotron period have been observed in simulations by Leroy et al (1981), Lembege and Dawson (1987) and Lembege and Savoini (1992), while *in situ* observational evidence has confirmed this behavior at the bow shocks of Earth (Vaisberg et al, 1984) and Uranus (Bagenal et al, 1987). This variability has been seen even in the early laboratory experiments of Morse et al (1972), although some details of the processes involved have remained open questions.

Evidence of shock variability has appeared even in the absence of external driving forces. Bale et al (2005) clearly demonstrates from Cluster data that “deviations are apparent in the shock foot, but they are not present in the ramp... This implies that the fluctuations in the foot do not propagate into the ramp, and therefore cannot be a source for downstream waves”. Horbury et al (2001) note that, “these may be signatures of unsteady reformation” citing Lembege and Savoini (1992), while Balogh et al (2005) assert that “Cluster electric and magnetic field observations have also highlighted considerable variability in the shock structure and profile even over relatively small scales associated with particle kinetic behaviour.”

All these comments taken together strongly imply that the shock either overturns and reforms on the one hand, or simply “breathes” on the other. But even in the latter case, the fact that the overall shock potential scale is nominally well above the electron inertial scale indicates that this periodic scale expansion alone is sufficient to limit the acceleration time scale for shock surfing. In other words, at the scales actually observed, periodic expansion of the potential scale essentially kills the mechanism. We recall from our Chapter 3 baseline studies the maximum velocity attainable by surfing during an ion cyclotron period

$$v_i(\tau) = \frac{q_i}{m_i} E_{\parallel} \tau = \frac{q_i}{m_i} \left(\frac{u_1 B_1}{c} \right) \frac{2\pi}{\Omega_{ci}} = 2\pi \frac{q_i}{m_i} \frac{u_1}{c} B_1 \left(\frac{m_p c}{q_i B_1} \right) = 2\pi u_1 \frac{m_p}{m_i}$$

Assuming a 50% “duty cycle” during shock reformation, this time scale limits the maximum attainable velocity of protons to $v_{\max} \approx \pi u_1$ or roughly the same as the shock drift velocity ceiling. Note that the ceiling also drops linearly with increasing ion mass.

Pickup Ion Disruption by Waves

One small point we have not previously covered in detail is that, for shock surfing to function optimally, a steady influx of pickup ions (PI) with the ‘standard’ velocity shell distribution is absolutely essential. Besides any suppression of the velocity ceiling due to other factors, the total flux of accelerated ions intimately depends upon the density of ions within a small region of velocity space as we demonstrated in our baseline studies. Although we have shown the inability of turbulent waves to return downstream ions to the front at planetary bow shocks, we nevertheless note that some velocity diffusion does occur. Even in the case where the cross-shock potential gradient may be sharp and confined to the ramp, such upstream diffusion, especially when coupled with the

expansion of the velocity shell through the magnetic gradient in the foot, will evacuate the PI population within the surfing candidate region.

Indeed, Kucharek et al (2007) have shown from both numerical simulations and Cluster observations that “the foreshock is characterized by extensive wave activity, which is caused by the interaction of the back-streaming particles with the inflowing solar wind”. The authors illustrate the deleterious effect that these waves produce in Figure 6-9 below, which clearly shows that upstream waves suppress particle reflection. We deduce by comparison with our own test particle simulations, that such waves deplete (through diffusion) the region in the velocity space where incident pickup ions are likely to be reflected. Although the simulation shown was configured to represent a quasi-parallel shock, the same conclusion may be drawn for the perpendicular case in light of the Cluster data presented by Bale et al (2005) and Narita et al (2007).

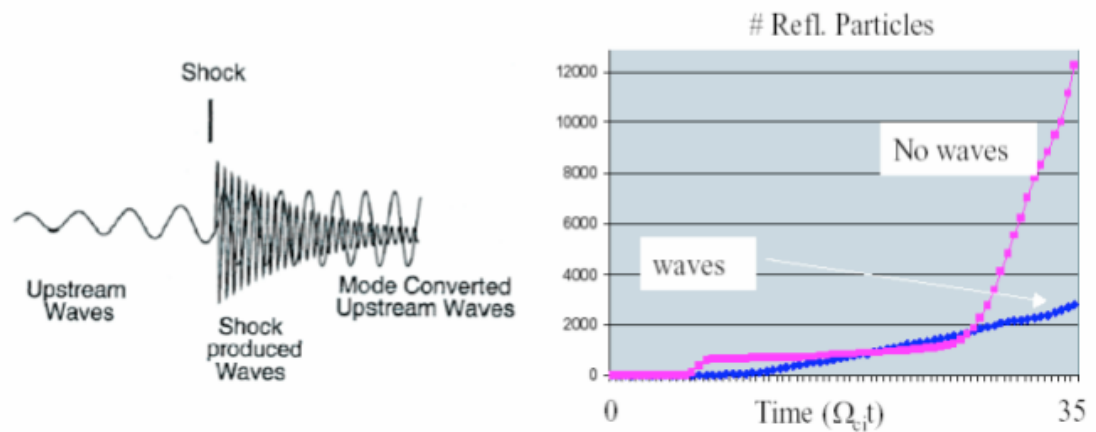


Figure 6-9 Effect of waves on reflected particles

Wave schematic is shown on the left and results of numerical simulations with and without upstream generated waves on the right.

After Kucharek et al (2007)

Turbulence Measurements

We now consider the effects upon shock surfing that may be associated with turbulence at or near the front, and justify our previous assumption that the scattering strength based upon the turbulent energy density is insufficient to return downstream ions to the front through velocity diffusion. The source of this turbulence, especially downstream, is still not well understood. Narita et al (2007) addressed this question in light of Cluster observations, and concluded that the increase in turbulent energy through the shock is inconsistent with the Rankine-Hugoniot relations, and that the intense spectra of compressible fluctuations observed in the magnetosheath requires some additional energy source at the front itself.

The plots from their study shown in Figure 6-10 below, exhibit the measured spectral energy density of three types of magnetic fluctuations observed by the Cluster spacecraft ensemble; two-dimensional slab, Alfvénic and compressible turbulence. The authors point out that, while Slab turbulence is primarily magnetic, Alfvén waves are correlated (or anti-correlated) with flow velocity fluctuations. Although this should favor the latter for driving shock ripples, there is no reason to eliminate the former as a coupling candidate.

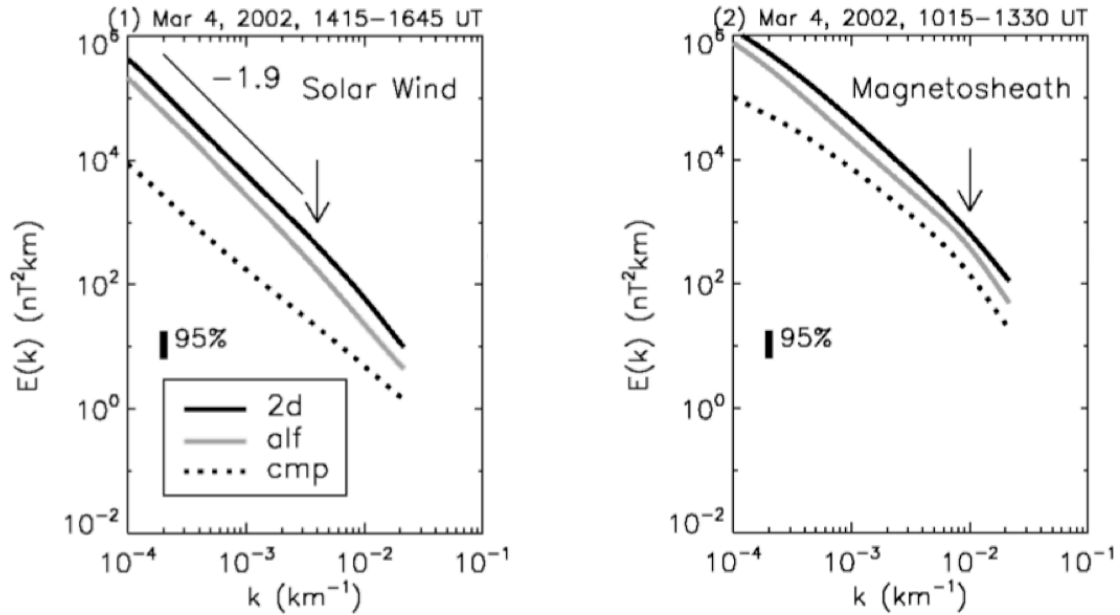


Figure 6-10 Turbulent Spectra at the perpendicular bow shock

The three curves represent two-dimensional, Alfvénic and compressible fluctuations as shown in the legend of the left panel (Narita et al 2007)

Note that the power of Slab and Alfvénic fluctuations share nearly the same spectral dependence, with the latter down in energy density by roughly a factor of 5. Whether this is a common feature at astrophysical shock fronts in general or peculiar to the Earth's bow shock has not yet been established. However, it should not be unlikely in the heliosphere since turbulent fluctuations are both transported in by the solar wind itself and generated by processes at the front. As already mentioned, this can be seen by noting the order of magnitude increase in turbulent power through the shock front, as seen by comparing the two plots in the figure. It is reasonable to assume that the turbulent energy density should scale through the shock as the square of the shock compression ratio, which in this case is approximately a factor of six, less than the factor

of ten actually measured. In other words, the spectral energy density coupled to the compression through the shock alone is insufficient to account for the increase.

In addition to these turbulence measurements, Horbury et al. (2001) were able to identify discriminate between the standing structure of the shock transition and large amplitude wave non-stationary in the shock frame. They found that the greatest contribution to the total downstream wave energy was posited in non-stationary ion-cyclotron waves, and no evidence at all of the long standing waves we have previously shown are required to promote rapid velocity diffusion.

It is of little consequence for the purposes of this investigation whether the source and nature of the turbulence is self-generated at the front by 'gyrating reflected ions', waves propagating or convected along the front from the quasi-parallel region upstream at the nose, or various plasma instabilities. The main fact is that the observed spectral energy density, when compared with our simulation results, is insufficient to promote the velocity diffusion necessary for ion shock return, but may perhaps explain ripples along the shock front.

Ripple Observations

From observations of the shock normal angle utilizing the Cluster Ion Spectrometry instrument (fluxgate magnetometer), Kucharek et al (2007) have demonstrated the existence of shock front ripples with amplitude $a = 76 \text{ km}$ and wavelength $\lambda = 700 \text{ km}$ as shown in Figure 6-11 below.

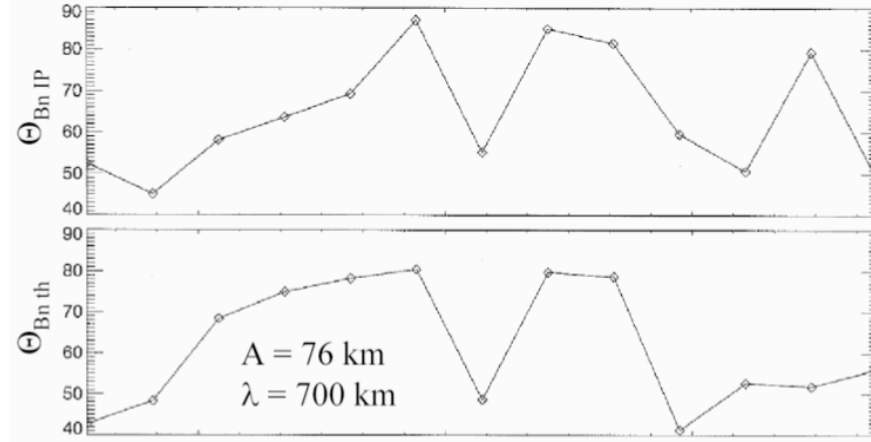


Figure 6-11 Evidence of Ripples at the Shock Front

The top panel represents observational data from Cluster and the bottom data synthesized from numerical simulation. The fitted parameters are for

amplitude $a = 76 \text{ km}$ and wavelength $\lambda = 700 \text{ km}$

After Kucharek et al (2007)

To estimate the effects of these ripples on the surfing velocity and efficiency, we first normalize the data to the local ion inertial length. Given the number density at the time of these measurements, $n \approx 19.0 \text{ cm}^{-3} \rightarrow \lambda_{pi} \approx 52.3 \text{ km}$ which gives us the normalized ripples parameters $\tilde{a} \approx 1.46$ and $\tilde{k} \approx 0.469$. These values are large enough to apply the abbreviated velocity ceiling function, giving $\tilde{v}_{\max} = \Omega_{cl} [\tilde{a} \tilde{k}^2]^{-1} \approx 1.6$, well below the Shock Drift ceiling, and only slightly greater than the upstream solar wind velocity. Calculating directly from the physical values gives roughly the same value $v_{\max} \approx 1.5 u_1$ given the nominal solar wind speed at the time of these observations at approximately 300 km/s.

Since the ripple curvature (and surfing velocity ceiling) depends more strongly on wavelength λ than amplitude (through the wavenumber squared) analysis of other data

would be useful in placing an upper bound on this. Bale et al (2005) do so, citing that agreement between observation and the Peredo model “also places an upper limit on the size of any large scale ‘ripples’ on this surface”, but whether the geometry & scale of planetary bow shocks in general place limits upon the supportable wavelengths has not yet been investigated.

Bow Shock Summary

We have established that shock surfing is not viable at the Earth’s Bow Shock, since each of the conditions necessary for surfing ions to exceed the shock drift ceiling is violated. The cross-shock potential is too weak and disjointed to produce a sustained electric field sufficient to restrain ions from being swept downstream by the Lorentz force. Ripples along the front are too severe to permit ions to exceed the shock drift velocity ceiling, even for the very steepest potential ramp gradient. Shock reformation or ‘breathing’ time is too short for the convective electric field to accelerate ions above the shock drift ceiling before the shock overturns or equivalently, the ramp becomes too broad for surfing. Upstream turbulent waves may evacuate a substantial fraction of the incident pickup ion population from the ‘surfing region’ of velocity space, severely limiting the flux of accelerated ions.

As to energetic particle observations, field aligned beams have been detected with energies of only one or two thousand electron volts. However, the energy of incident thermal solar wind protons is of the same order – suggesting that these particles have undergone at most a single reflection. The upstream trajectories of the ‘crossing’ and single-bounce’ particles seen in our baseline simulations are consistent with both the

observed energy and the upstream excursion distance. Hence there is no direct evidence for shock surfing at all.

The Outer Planets

Beginning in 1972 seven missions have been launched to explore the outer planets – Pioneer 10 and 11, Voyager 1 and 2, Ulysses, Galileo and Cassini. Jupiter is the largest and innermost of the gas giants, but only the Galileo Mission was sent expressly to sojourn, observe and collect data there over an extended period of time. The other probes all executed ‘flybys’ using the planet’s gravity in slingshot trajectories on their way to other destinations, taking advantage of the close proximity to make relatively brief investigations during their passage.

The Galileo Mission launched in 1989, reached Jupiter in 1995 and began a 20-month, 11-orbit tour of the Jovian system, including 10 close encounters with Jupiter's major satellites. The mission was extended with 24 more orbits providing additional encounters with all four of Jupiter's major moons, and it continued to investigate the planet’s surrounding magnetic environment for several years until its programmed impact in 2003.

The primary objectives, especially of the early missions, were the investigation of planetary atmospheres, characterization of satellites and rings, and an understanding of the gross features of the planets’ magnetospheres. Together, these probes have provided a rich body of knowledge, along with stunning detailed images of objects so distant that they were once little more than photographic smudges from Earth-bound telescopes.

However, unlike the vast improvements in data acquired by Cluster when compared with earlier missions, the outer heliosphere has not yet been explored with such precision. Numerous examples in the literature highlight following points:

- Multiple probes are essential in providing detailed information on the features of dynamic structures. NASA had originally planned a "Grand Tour" employing a four spacecraft 'Cluster' to explore the outer planets, but this was canceled in 1972 and replaced by the Voyager missions.
- Due to limited instrument sensitivity, especially in the ability to measure the electric fields of potential gradients, the locations of shock crossing events described in the literature from the early missions were inferred from the detection of upstream plasma waves or electron heating rather than by direct observation.
- Absolutely no direct measurements of cross shock potential have been reported from any mission to the outer planets. Neither has there been any direct evidence of accelerated protons where they were expected to be found. Whether this lack of data is due to insufficient instrumentation, absence of any real physical phenomena at all, or simply due to prevailing SW and shock conditions, has not yet been determined.
- Another disadvantage of single-shot probes is their mechanical fallibility in the presence of hostile environments. Data from the first Jovian mission was lost or corrupted due to high intensity levels in the planet's radiation belt. Telemetry failures also interrupted data transmission and acquisition from the first Termination Shock crossing by Voyager 1.

Such single probe missions making transient observations may also fall prey to periodic lapses in fortuitous observing conditions. The Cassini spacecraft was launched to study Saturn and, during its passage ‘skimmed’ along the dusk-side flank of the Jovian magnetosphere crossing the shock (as reported) more than forty times before continuing outbound, while early missions had traversed the dawn-side. Unfortunately, while some new data was acquired, the solar wind at the time was highly active and variable, with strong magnetic fluctuations superimposed on the shock. But variability of the shock itself is not necessarily be due solely to influence from the solar wind, which in turn may suggest long term constraints upon particle acceleration by SS. Zsego et al (2003) noted the absence of reflected upstream protons detected by the probe, and argued that the Jovian bow shock may not always be in a steady state.

In spite of the mission shortcomings, we now examine samples from the body of robust data acquired and show that observations generally support our position that magnetospheres of the gas giants are necessarily highly dynamic and the shock ramps either too shallow or too turbulent to support shock surfing.

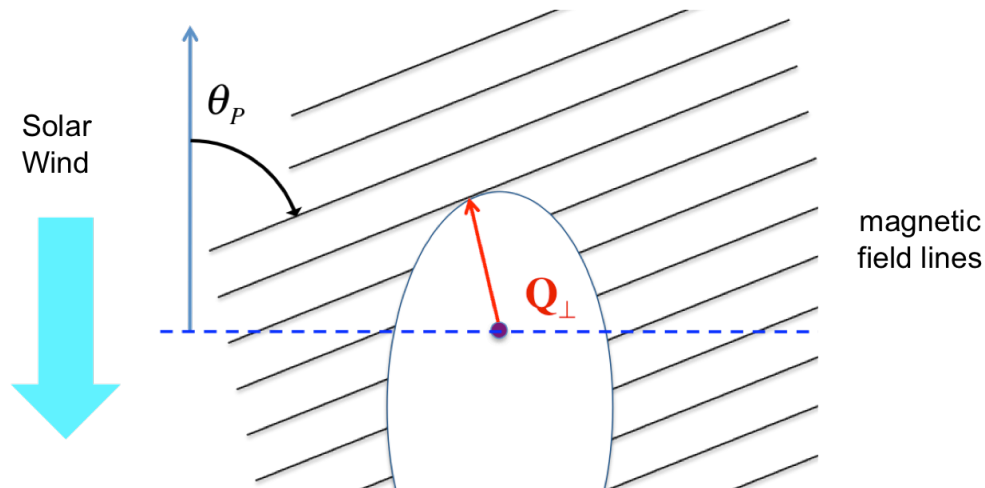


Figure 6-12 Parker Angle at an Outer Planet

The steeper angle places the quasi-perpendicular region near the nose.

Common Morphology

Several features distinguish the planets in the outer heliosphere from those within. The Parker angle at the outer planets, as sketched in Figure 6-12 below, becomes increasingly more perpendicular with solar distance and may thus seem more favorable by comparison with the inner planets. The quasi-perpendicular shock surfing regions should lie relatively close to the noses of the planetary magnetospheres, while the shock fronts along the flanks – at least in the ecliptic – are parallel. Aside from the actual nose region, this leaves only a small ribbon of points where the shock surface intersects planes perpendicular to the ecliptic. At this locus of points along the flanks, the off normal velocity vector effectively reduces the cross-shock potential, and the convective electric field acquires a component normal to the front.

Under stable and otherwise ideal conditions, these two factors alone become unsuitable for supporting multiple reflections, so that successful shock surfing ions would of necessity need to accelerate quite rapidly near the nose before being swept along the

shock surface by the solar wind. But, as has become, these shocks are far from being stationary and quiescent.

Shock Variability

The density and magnetic field strength of the solar wind decrease with distance from the sun, while the intrinsic planetary magnetic fields of the gas giants are much stronger than those of the inner planets. This combination leads to vastly larger solar standoff distances and overall scales, which produce magnetospheres much softer and more susceptible to external driving forces as compared with the rather stiff and temporally robust size and shape of the bow shock found at Earth. Hence, as deduced by Szego et al (2003) from the Voyager mission, the Jovian magnetosphere is easily compressible, which should to some extent be characteristic of the other gas giants as well.

Jupiter in particular possesses an exceptionally strong magnetic field, but a magnetosphere so immense that, if visible, could easily be seen from Earth subtending a solid angle greater than the Moon's. The Jovian standoff R_N has been variously measured between $\sim 130R_J$ and $\sim 215R_J$ where $R_J \simeq 7.1 \times 10^4 km$ is the Jovian diameter (Wenzel 1993), which is two orders of magnitude greater than $\sim 10R_E$ at the Earth, where $R_E \simeq 6 \times 10^3 km$.

Because of Jupiter's great size, the boundaries are highly sensitive to both externally and internally driven motions (Kivelson 1997). Jupiter possesses a strongly co-rotating magnetosphere, which responds not only to variations in SW pressure, but also to strong tidal effects from its large satellites. Galileo first encountered the Jovian bow shock in 1998 and data from missions thereafter continued to reveal an asymmetric,

flattened and wobbling magnetodisk with a period roughly that of the planet itself $T \sim 10$ hr (Huddleston et al 1998).

The authors also note that the Galileo mission confirmed the Jovian standoff varies with solar wind dynamic pressure as $R_N \propto P_{SW}^{-\gamma}$ where $\gamma = 0.22$ in contrast to the terrestrial exponent $\gamma = 1/6$. This previously known fact has been attributed to the presence of hot plasma and centrifugal stretching in the Jovian magnetodisk that lessens the pressure gradients in the outer magnetosphere.

Due to their great size, transit time scales are much larger at the outer planets' magnetospheres. We contrast, for example, typical solar wind and MHD wave transit time down the flanks of the Earth's bow shock (minutes) with those of Jupiter (hours to days). Comparing the latter with the time scales of solar wind fluctuations, one may conclude that the Jovian bow shock may seldom stay in a steady state (if it is ever reached at all) simply due to its great size (Szego et al 2003).

The outer planets also possess short rotational periods and large satellites. The effects from the latter are most easily visible from ring features attributable to orbital parametric resonances, while data show that the coupling with the planet's magnetic field lead to flattening of the magnetospheres and a substantial temporal 'wobble' of the resulting magnetodisks. Coupling is also evident from *in situ* images showing radial striations in the rings, which indicate driven resonant axial modes.

Jupiter has a volcanically active moon, Io, which emits hot plasma into its magnetosphere. Galileo observations confirmed that the volcanism of Io is episodic, possibly leading to variable mass loading of the Jovian magnetosphere and thereby adding another source of variability to its size. Expanded shock boundaries were observed early on Galileo's inbound pass, perhaps in part due to hotspot activity at Io in

the fall of 1995 prior to its encounter, while observations within the magnetosphere are consistent with externally driven motion of the boundaries (Huddleston et al., 1998.)

Neptune's moon, Triton, is also volcanically active. Its thin atmosphere – composed mainly of nitrogen with small amounts of methane – most likely originates from Triton's volcanic activity, which is driven by seasonal heating by the Sun. Besides Venus and Earth, these are the only bodies in the solar system besides known to be volcanically active at the present time.

However, Saturn's moon Titan has a very dense atmosphere, which some surmise bleeds away into the Saturnian magnetosphere. Recent observations by Cassini have confirmed the existence of 'ice volcanoes' with speculation that Titan's atmosphere may in this way be continuously replenished. Hansen et al (2006) reported from the closest Cassini flyby of Saturn's moon Enceladus in early 2005, evidence "for ongoing endogenic activity" in a region centered on satellite's south pole. While these are not strictly volcanic in the strict sense, any source of gas and dust may provide a particle influx to materially affect Saturn's magnetosphere, and may be another feature common to the outer planets.

Finally, all four of the gas giants rotate on their axis very rapidly, with remarkably similar rotational periods. Coupling between these planets, their substantial atmospheres and their magnetospheres should not be surprising. Aside from shaping the coarse features of their bow shocks such as general shape, this strong coupling may provide yet another factor in determining the relative stability of finer features including surface ripples, shock reformation and limits upon sustained potential gradient scales.

Taken together, these common features of the outer planets are sufficient to characterize their bow shocks as highly active and variable. However, in addition to

external influences on the Jovian magnetosphere, simulations by Quest (1985, 1986) and Leroy et al (1982) have shown that, at the high Mach numbers observed, normal energy dissipation mechanisms that stabilize Earth's bow shock might not be adequate. Consequently, the high variability of the Jovian bow shock in comparison with ours should not be surprising, and may explain the absence of shock-reflected ions. Indeed, the drastically changing position of the Jovian bow shock has been observed not just by Cassini but also by all previous missions.

We have already suggested from scaling arguments that the Jovian bow shock as well as those of the other gas giants must lack the temporally rigid structure of our own. We suggest by analogy comparing the stiff response of a small well-inflated party balloon with that of a large, as yet underinflated, weather balloon at launch – and picture the 'flabby' undulations in response to even the gentlest local wind gusts. Szego et al (2003) take this further, and suggest that this characteristic is connected with the low ratio of sheath density to that of the local solar wind. Their intuitive model is remarkably similar – a simple spring:

A softer but longer spring makes larger excursions but oscillates more slowly than a shorter and stronger one (the spring length is connected with the bow shock obstacle distance and the softness is connected with the sheath density).

The Potential Ramp

Remarkably, electric field data for the cross-shock potential are not available from any of the missions to the outer heliosphere, and shock front locations and length scales have therefore been deduced by other means. For example, Voyager plasma wave measurements at Jupiter noted by Scarf et al (1981) found intense broadband

turbulences in a thin layer of the shock, preceded by intense bursts of electron plasma oscillations upstream. They also inferred from magnetometer data that the Magnetosphere boundaries vary widely in response to changes in the solar wind pressure.

Every modern probe has carried a basic suite of instruments to measure magnetic field, plasma waves and ion fluxes. For example, instrumentation aboard Cassini consisted primary of magnetometer, radio and plasma wave science instrument (RPWS), and the Cassini plasma spectrometer (CAPS). The main observational objectives were magnetic field data, wave detection and charged particle fluxes. The CAPS instrument in particular was more extensive than its predecessors, having three independently operated sensors: the ion mass spectrometer (IMS) to analyze ion composition and plasma dynamics, the electron spectrometer (ELS), and the ion beam spectrometer (IBS) to measure narrow, beam-like distributions without mass separation. Historically, magnetic field measurements have been particularly successful, with reliable data gathered from the earliest missions forward.

As previously discussed, it has often been customary to assign similar scales to the electrostatic potential gradient as compared with the magnetic foot and ramp. Regardless of the accuracy of magnetic data, we have shown from Cluster observations that the overall electric potential gradient at the Earth's bow shock, despite containing fine features visible as 'spikes', is no smaller than the magnetic ramp, and cyclically 'breathes' or expands out to the scale of the foot.

Although we have no unimpeachable observational evidence to the contrary, we strongly believe that the difficulty in detecting the presence of cross-shock electric fields together with the complete absence of reflected protons, strongly indicate that shock

surfing was not taking place during the passage of any mission probes. This is not to assert that no cross-shock potential exists, but rather that the difficulty in its detection most probably follows from large gradient scales and consequently weak normal electric fields. In other words, we believe that the relationship between potential and magnetic field scales seen in the foot and ramp at the Earth's bow shock should be found at the outer planets as well.

Case Studies

If the arguments we have so far presented do appear sufficiently compelling, we now offer several examples of anecdotal evidence. Tsurutani et al (1993) cite what appears to be an odd event from the Ulysses mission in which Balogh et al (1992) and Bame et al (1992) report a shock crossing where the magnetic field was directed radially along the Jupiter-Sun line. This observation was made downstream in the magnetosheath, but is nevertheless quite strange – and suggests extreme redirection of the solar wind plasma in the region, with associated bending of the field lines. If this interpretation of the data is correct, then the bow shock was quasi-parallel (surprising for that region along the flank) and upstream waves consequently attributed to reflected ~ 2 keV solar wind protons. However, the authors attribute this unusual morphology to low solar wind ram pressure associated with the trailing portion of a high-speed stream. Even so, the energy ceiling of 2 keV (based on the piston energy) is consistent with a single reflection that might easily occur at a parallel shock in the presence of sufficiently large SLAMS (Small Large Amplitude Magnetic Structures) with scales greater than several ion gyroradii and with energy density a few times that of the solar wind magnetic field. Once again, the scale of these objects was not directly observed, but inferred from

pulse durations assuming a convection speed consistent with that of the local magnetosheath. In any case, even if the shock obliquity were misinterpreted, the observed proton energy would still be consistent with a single ‘reflection’ by way of the shock drift mechanism.

Zsego et al (2004) presented from Cassini yet more evidence of the highly active nature of the Jovian bow shock. Down the flanks, the spacecraft ‘skimmed’ along the front producing several distinct sets of shock crossing data. Particularly noteworthy were the shock-normal magnetic field angles, ranging from parallel $\theta_{BN} = 2^\circ$ through oblique, to quasi-perpendicular and back within the span of a few days – effectively representing a small region of the magnetospheric flank. They reasoned that the shock front was convoluted in the vicinity, certainly an extreme deviation from the transition expected at a smooth locally planar front. Although the spacecraft’s trajectory lay close to the solar equatorial plane, far from the locus of perpendicularity described above, these features may be endemic and characteristic of the shock in all regions.

Where the shock was perpendicular, the authors deduced from the frequency of Langmuir waves $\nu \sim 10 \text{ kHz}$, the derived plasma parameters $n \simeq 0.5 - 1 \text{ cm}^{-3}$, $T_e \simeq 2.6 \text{ eV}$, $B \simeq 0.8 \text{ nT}$ and $V_A \sim 32 \text{ km/s}$. The SW Mach number was $M_A = 10$ while the shock-normal component was $M_{A\perp} = 7$ with a shock compression ratio $R_C \simeq 2.74$. Based upon these parameters and the simulations of Leroy et al (1982), the authors estimate the nominal quasi-perpendicular foot scale to be around the proton gyroradius $r_g \sim 1500 \text{ km}$ and the shock width roughly the ion inertial length $\lambda_{pi} \sim 300 \text{ km}$. Once again, the Cassini shock crossings were inferred from the plasma frequency and the rise in ion temperature rather than from direct field profile measurements. However, if the

values are reasonably accurate and we assume a corresponding electrostatic potential scale, then the ramp is far too wide to permit surfing ions to exceed the shock drift velocity ceiling. This assumptions of the cross-shock potential scaling is further supported by the reduction in expected cross-shock ram pressure, as seen from the shock-normal Mach number derived above.

Uranus & Neptune both have magnetic fields that are inclined from their polar axes by approximately 60° and with rotation periods of 11 hr and 16 hr respectively. As a result, their magnetospheres do not maintain a quasi-static aspect but rather induce constant motion in the shapes of their shock fronts. Indeed, their auroras are extremely faint and almost undetectable, strongly suggesting a lack of temporal constancy and thus another impediment to shock surfing.

Cairns et al (1991) compared the Neptunian magnetic field data at the spacecraft along with two separate observations of intense Langmuir waves and determined that the standoff distance changed significantly within a period of 85 minutes, from $\sim 40R_N$ to $\sim 35R_N$. They assert that this was not in response to a change in the solar wind ram pressure, but rather controlled by the rotation phase of Neptune's oblique, tilted dipole.

Krimigis et al (1988) discuss the observations of Uranus' bow shock by Voyager 2. Although they identified energetic populations of both electron and ions, they concluded that characteristics of the upstream events "strongly favor leakage from the magnetosphere as the source of the particles and do not support predictions of *in situ* Fermi acceleration".

Observations by the low energy charged particle (LECP) experiment on the Voyager 2 spacecraft also indicate coupling between Neptune and its satellite Triton. Mauk et al (1991) discuss the largely anisotropic population of heavy ions (probably N^+)

co-rotating with Neptune, possibly due to interaction with a neutral torus. Although particles with energies greater than 20 keV were occasionally detected, the authors identify these fluxes with Triton rather than with acceleration processes at or near the shock front.

Moses et al (1990) compared plasma wave measurements from the Voyager 2 crossing of Neptune's bow shock with those from the bow shocks of Earth, Jupiter, Saturn and Uranus. They conclude that, when normalized to the local solar wind ion thermal energy density at each planet, wave amplitudes above $10^{-2}\omega_p$ are significantly higher at the outer planets and the electric field energy density exhibits exponential dependence on ion thermal Mach number. This is consistent with a trend that has recently been shown to occur out to Neptune's bow shock, and may be yet another contributing factor to the activity and variability seen in the shocks of outer planets.

Summary & Conclusions

We have shown from direct observational evidence that shock surfing cannot function effectively at the Earth's bow shock. The broad potential ramp, ripples at the front, shock variability and reformation time scales each are sufficient to suppress the mechanism. In combination, they assure its failure as verified by the lack of any significant detection of accelerated ion flux.

We have argued, in the absence of compelling bona fide evidence, that theoretical objections prohibit shock surfing from functioning at the outer planets. We have cited a body of anecdotal evidence that we believe, taken *en masse*, is sufficient to explain the complete lack of any significantly accelerated ions detected at the gas giants. While future missions may provide superior data resolution, especially unambiguous

measurement of the cross-shock potential, many common shock features at these planets such as high variability, broad scales, rippled fronts and temporal reformation have been verified by data from multiple missions. The consistency and agreement between these observations are sufficient for us to assert that shock surfing is not viable at the outer planets.

We now turn our attention to the Termination Shock at the boundary of the heliosphere, and apply the same criteria to assess shock surfing viability there.

Chapter 7 The Termination Shock

We have established, by simulations and theoretical arguments, a set of conditions necessary for shock surfing to function effectively, and have shown by abundant observational evidence that each of these is absent at the Earth's bow shock. It is difficult to be completely confident in drawing similar conclusions about the morphology and behavior of outer planet bow shocks, especially in light of the relatively few observations as compared with those made locally. Although data from the gas giants are less comprehensive and therefore less compelling, we have still seen sufficient evidence to conclude that at least some if not all of the conditions necessary for surfing are also violated, and may therefore assume with reasonable certainty that shock surfing is not viable there.

From the various exploratory missions, we have also seen that the abundance of sustained, high-resolution *in situ* measurements profoundly diminishes with distance. This is due to the smaller number of probes, the shorter *in situ* duration of fly-by versus orbital missions, and the epoch of available onboard technology as determined by the time-of-flight required to reach more distant target objectives. All this naturally limits both the scope and fidelity of investigations even farther out from the sun, and investigators face substantial challenges at the Termination Shock where only a few well-defined shock crossings have been observed.

Recap & Overview

In order to assess the Voyager data, we recapitulate specifically those conditions essential for shock surfing. Only very nearly perpendicular regions of the shock are viable. We have shown that, for magnetic field deviations from the perpendicular much

beyond about $\pm 10^\circ$, both the surfing velocity ceiling and flux are critically suppressed. A sufficiently steep and strong cross-shock potential is required to confine pickup ions at the front against the Lorentz force that tends to drive their escape downstream. The shock must be smooth and free of corrugations that disrupt and thus prematurely terminate the surfing mechanism. Finally, these conditions must be sustained without interruption long enough for the convective electric field to accelerate ions substantially beyond the shock drift velocity ceiling. We stress that, while shock overturning and reformation play an obvious role in constraining sustained acceleration, other time scales such as periodic 'breathing' or other temporal variation in the ramp width may produce the same result while the potential ramp is too broad to support shock surfing.

Due to the paucity of data from the Termination Shock and especially the absence of electric field data, many shock features were not directly measured but rather derived or inferred from available data. Some shock parameters were deduced immediately from observed data by way of the Rankine-Hugoniot relations and are therefore reasonably trustworthy. Other features however have only been proposed well after post-mission analyses from simulation results performed in an attempt to fill in the missing information. While some of these investigations are compelling especially in light of the care taken in matching the simulation parameters to those actually observed, the results cover a broad range of specific interests.

As an incidental note, we remark that the TS is concave upstream and thus one other critical problem is absent at there – that of convex planetary bow shock geometry and the issue of ion advection and escape along the flanks with accompanying acceleration time constraints. At the TS, there is nowhere for ions to escape except

through the shock front itself. Aside from this one advantage, we now offer several reasons to doubt that conditions necessary to support shock surfing actually exist there.

Shock Description

The solar wind flows outward from the Sun forming a ‘bubble’ of solar material in the Local Interstellar Medium (LISM). Three boundary features characterize the structure – the heliopause, where the total internal pressure (magnetic, thermal and momentum flux) matches the external pressure of the much sparser LISM, and the closer interior TS where the solar wind changes from supersonic to subsonic. Outside the heliopause, a bow shock exists due to a possible combination of the non-Keplerian motion of the LISM plasma and the ‘peculiar’ motion of the solar system as it oscillates about the galactic plane. The intervening region is called the heliosheath. The structure of the heliosphere differs from planetary magnetospheres primarily insofar as planets do not emit winds, and lack the momentum component of the internal pressure.

Interest in probing the TS extends beyond simply understanding its composition and nature, as it was long believed to be a site of Anomalous Cosmic Ray (ACR) production. Due to the winding up of the solar wind magnetic field, the Parker angle near the ecliptic plane was thought to be very nearly perpendicular, offering promising conditions for both efficient Diffusive Shock Acceleration (DSA) and Shock Surfing (SS).

Voyager Observations

Two spacecraft have encountered the TS, Voyager 1 (V1) in December 2004 and Voyager 2 (V2) in September 2007. Data from the V1 crossing was unfortunately not recorded, since mission funding limited the monitoring periods, so that the actual crossing was deduced as having taken place between two data sets (bracketing a

telemetry gap) where the change in plasma parameters so indicated. Instrumentation aboard these probes was somewhat primitive by current standards, due to the technological state of the art when they were launched in late 1977 – V2 on August 20, and V1 roughly two weeks later.

After gathering data from the outer planets, the V1 and V2 continued their mission outward to observe the TS, which they encountered at 94 AU (latitude 34.1°) and 84 AU (latitude -27.5°) respectively. Richardson (2008) argues that the difference in distance measured by the two spacecraft is far too great to be attributable simply to temporal changes in SW pressure (at most 2-3 AU), but rather indicates a blunt structure as illustrated in Figure 7-1 which follows from the two spacecraft encountering the shock at different locations along the front.

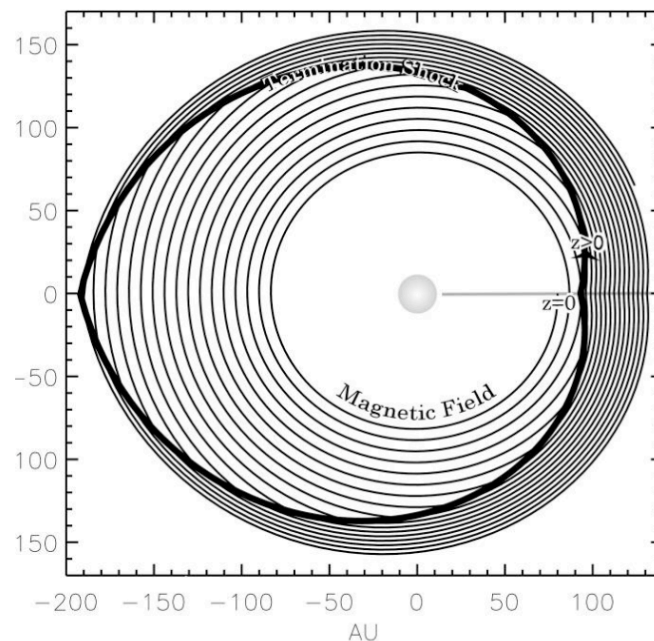


Figure 7-1 Magnetic field angle model of the Termination Shock
 The plot represents a slice through the ecliptic plane, and the right-hand boundary the region explored Voyager showing a 'blunt nose'
 After Schwadron et al (2008)

The heliospheric model evolved over the past two decades with the acquisition of observational data and, in the past few years, has changed substantially. The Interstellar Boundary Explorer (IBEX) mission revealed an extremely bright and narrow ribbon of energetic neutrals as shown in Figure 7-2, which challenges the previously held notion of a comet-like structure moving through the “galactic wind”.

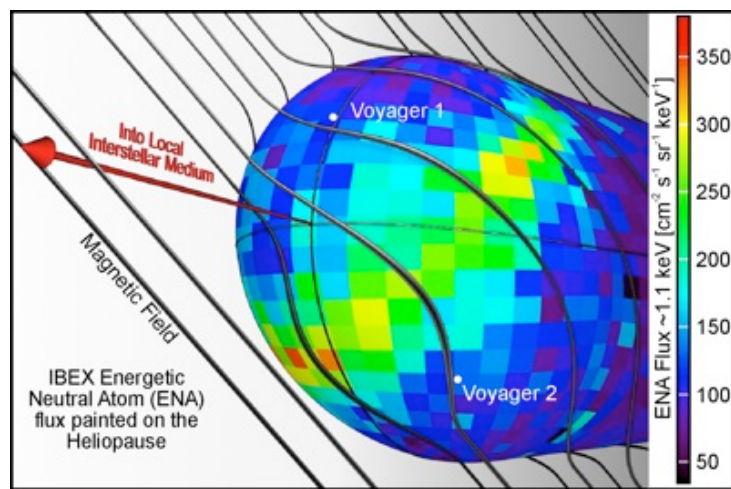


Figure 7-2 IBEX Neutral Atom Map

A bright ribbon of neutrals is superimposed on the heliopause
After McComas et al (2009)

This ‘bubble-like’ model was further supported from observations by Cassini’s Ion and Neutral Camera (MIMI / INCA), while measurements by the Solar Wind Anisotropies (SWAN) instrument onboard the Solar and Heliospheric Observatory (SOHO) revealed a distorted heliosphere, likely under the effect of the local galactic magnetic field (Kerr, 2009). Interest in the heliospheric bow shock has also gained additional momentum from images gathered by space-based platforms that show clear evidence of such structures

at nearby stars as shown in the examples in Figure 7-3 below. In short, the grand structure of the heliosphere and the physics that govern it currently remain open questions.

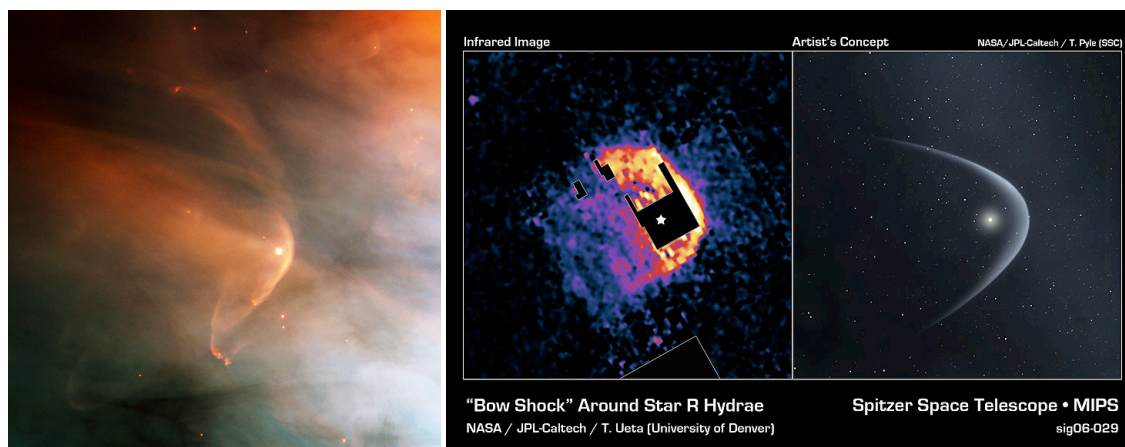


Figure 7-3 Bow Shocks of Nearby Stellar Systems

The left panel shows LL Ori as observed by Hubble (1995) and the right panel R Hydrae by Spitzer as a composite image with artists conception.

(NASA, JPL, and Toshiya Ueta)

Primary Data & Surprises

Aside from the large-scale structure deduced by comparing observations from the two Voyager missions, the data from V2 provide somewhat more detailed information about the shock front structure itself. The spacecraft followed an essentially radial trajectory outward so that, rather than ‘skimming’ the shock as did Cassini along the Jovian flank, only one shock crossing was expected. Surprisingly, five front crossings were inferred from the data, but only three provided data sufficient for more comprehensive and detailed analysis of the structure. Listed below in Table 7-1 are shock parameters derived from data gathered at two of these distinct shock crossings.

Since only two crossings yielded sufficiently unambiguous data for post reduction analysis, interpretations drawn from them must be made with great care, especially since the shock was not in stasis, and its observed state may not be typical or even representative.

Table 7-1 Termination Shock Parameters

The values are derived from Voyager 2 data assuming locally planar MHD and applying the Rankine-Hugoniot conditions. Richardson (Nature 2008)

Parameter	Termination shock motion	
	TS-2: outwards	TS-3: inwards
East-west shock normal angle	$188.0^\circ \pm 4.0^\circ$	$5.8^\circ \pm 10.3^\circ$
Shock speed	$94.0 \pm 3.4 \text{ km s}^{-1}$	$67.9 \pm 17.3 \text{ km s}^{-1}$
North-south shock normal angle	$2.0^\circ \pm 6.2^\circ$	$-4.6^\circ \pm 19.2^\circ$
Angle between shock normal and magnetic field	$82.8^\circ \pm 3.9^\circ$	$74.3^\circ \pm 11.2^\circ$
Compression ratio	2.38 ± 0.14	1.58 ± 0.71
Solar wind fast-mode Mach number	4.9 ± 0.1	8.8 ± 1.2
Heliosheath fast-mode Mach number	1.1 ± 0.1	2.8 ± 0.4

Unfortunately, V2 was incapable of directly measuring quasi-static electric fields associated with the potential gradient, but the probes were able to detect plasma waves. These observations were correlated with those from the magnetometer to identify the shock crossings. Gurnett et al (2008) assert that two “well-defined bursts of broadband electrostatic waves... corresponded to shock crossings identified in the magnetometer and plasma data” noting that similar bursts have been identified as signatures of planetary bow shocks. Several features are immediately apparent from these two front crossings.

First, the shock was apparently in motion. Whether the displacement was due to periodic oscillation in the front, ripples propagating along the surface, overturning and reformation, or other dynamic processes remains in our view an open question. The correct interpretation is germane to our investigation, however, since both reformation

time and ripple length scales in particular play crucial roles in suppressing the shock surfing mechanism.

Second, while the shock front itself remained nearly perpendicular, the magnetic field angle did not, owing to the blunt and asymmetrical structure of the nose as shown previously in Figure 7-1. The field angle at crossing TS-2 was measured at 82.8° or roughly 7 degrees from the perpendicular, and therefore theoretically capable of supporting reasonably robust confinement, given a sufficiently strong and steep cross-shock potential gradient. But the TS-3 crossing at 74.3° is more than 15 degrees off normal, which as we have shown, is enough to suppress the maximum surfing velocity by 75%, all other conditions being ideal.

Third, the strength of the shock as defined by the compression ratio was only weak to moderate, the outward TS-2 being comparable to the Earth's bow shock with compression ratio $R_C \approx 2.38$ and the inward TS-3 having a meager $R_C \approx 1.58$, although the compression ratio makes a smaller contribution to the cross-shock potential than does the Mach number.

Finally, the most surprising observation is that the downstream plasma velocity appeared supersonic with respect to the local plasma parameters. Upon initial analysis, this raised fundamental questions about the nature of the shock itself, with several consequences for shock surfing as well. Both ion and electron temperatures confirmed that the piston energy was not going into heating SW ions, suggesting to many observers that most of the energy must be transferred to PIs or other energetic particles. Yet data from neither Voyager mission showed any such evidence of back-streaming ions from the shock front. While Voyager 2 was equipped with a plasma instrument capable of directly measuring suprathermal ions, it could not discriminate between solar

wind (SW) and pickup ions (PI), and consequently no direct observations about the partition of the dissipation energy were possible. As noted by Wu et al (2008), this is most likely due to the fact that, since solar wind ions gain velocity $\Delta v_{\perp} \sim u_{SW}$ through a single reflection before being swept downstream, it is “difficult to observationally separate reflected solar wind ions from pick-up ions downstream.”

We mention one caveat in passing that, as noted in Stone et al (2005), during the crossing made by V1, a local source of anomalous cosmic rays was not observed. However, there is insufficient evidence from this to conclude that particle acceleration is necessarily unlikely at the TS, since the paucity of data may simply be consistent with the probes having passed through the region during a period of quiescence. Nevertheless, Schwadron et al (2008) advanced a model where “strong nose-to-tail asymmetry” disfavors ACR production at the ‘nose’, and argued that the observed modulation ‘bump’ in the power spectrum is consistent with acceleration taking place, not through DSA, but “farther out in the heliosheath by a distributed stochastic process (Fisk et al. 2006)”.

Time and Length Scales

Knowledge of both time and length scales is essential in determining the viability of shock surfing. The ramp width affects the strength of the retarding electric field, while the wavelength of possible ripples along the front affects the local curvature of the surface. Both of these impact the maximum velocity attainable by surfing ions. The time scales of large fluctuations in the front structure, whether they be undulations, overturning and reformation, or secular displacements of the front itself are necessary for our analysis as well. As we have shown, these not only affect the velocity ceiling, but

also the flux of accelerated ions. Since Voyager 2 was a single probe which encountered the region in one pass, some of these scales were not directly observable, but may be deduced by suitable interpretation of the data, which must be performed with care, especially when several alternative models may present themselves.

Table 7-2 Summary of derived front crossing parameters
TS-2 shock is moving outwards, TS-3 inwards. The quantity r_s is the shock compression ratio, w_s the shock width and τ the temperature jump ratio. After Richardson et al (2008), Wu et al (2008).

Table 1. Voyager 1 (V1) and Voyager 2 (V2) Termination Shock Encounters

encounters	$r(AU)$	r_s	u_u (km/s)	u_d (km/s)	θ_{Bn}	w_s (km)	τ	T_d (k)
V1	94	$2.6^{+0.4}_{-0.2}$	200	100	-	-	-	-
V2(TS - 2)	84	2.38 ± 0.14	325	150	$82.8^\circ \pm 3.9^\circ$	300,300	10	10^5
V2(TS - 3)	84	1.58 ± 0.71	250	150	$74.3^\circ \pm 11.2^\circ$	100,000	10	10^5

[Richardson et al., 2008; Stone et al., 2005; Decker et al., 2005; Burlaga et al., 2005]

We refer to the example in Table 7-2 above, which reiterates some of the data from the previous table, but includes estimates on the widths of the two shock crossings TS-2 and TS-3. Based upon general agreement among several authors, the ion inertial length at both shock crossings should be approximately $\lambda_{pi} = c / \omega_{pi} = 6000$ km . Under this assumption, the table gives normalized values between 15 and 20 ion inertial lengths, which is extremely wide indeed. However, this is not an unreasonable estimate for the entire shock front structure including foot, ramp and overshoot, consistent with the upstream cyclotron radius of a proton reflected once from shock at the given Mach number. Recall that the foot alone has length scale roughly $L_{foot} \sim \lambda_g = c / \Omega_g \sim M_A \lambda_{pi}$. However, this is in contradistinction to the relative widths when compared to the shock Mach numbers shown for TS-2 and TS-3 in Table 7-1, and may therefore be somewhat suspect, unless some shock features actually steepen with Mach number.

Cross-Shock Potential

Several post-mission attempts have been made to explain the conundrum of supersonic downstream flow, now generally attributed to incomplete SW ion heating resulting in higher than expected magnetosonic speed. Since the solar wind is not sufficiently heated, some additional form of dissipation is needed to explain the existence of the shock structure, and energy diversion into the pickup population has been proposed. Prior to the Voyager observations, predictions abounded regarding both expected shock parameters and arguments promoting the viability of shock surfing, advanced independently by Zank et al (1996), Lee et al (1996) and by Lipatov and Zank (1999). However, these were based upon the results of numerical simulations applying parameters assumed to be likely at the termination shock and, in particular, an extremely small length scale for the potential ramp – on the order of a Debye length. Since observations have revealed a cool solar wind with $\beta \approx 0.05$, this implies a scale much smaller than an electron inertial length, which is not physically reasonable. At the same time, the implications of the data are in contrast to the earlier work of Liewer et al (1993) whose one-dimensional hybrid simulations showed that SW ions should provide most of the dissipation, even with relatively high pick-up ion density.

Although observations failed to detect accelerated ions, Burlaga et al (2008) attributed the weakness of the TS-3 encounter to the possible presence of “suprathermal pickup protons”, while Burrows et al (2009) have again argued in favor of shock surfing, provisionally assuming that “for a strong enough cross-shock potential and sufficiently narrow shock ramp MRI [Multiply Reflected Ions] acceleration can account for the ‘missing’ energy of the downstream solar wind plasma.” Due to the complete absence of direct field measurements, it is essential to determine the scale of the shock ramp

electric field component by some other means, since this is of paramount importance to shock surfing.

Following the Voyager data, Wu et al (2009) and Winske et al (2009) appear to have resolved both of these questions by independently performing self-consistent numerical simulations, varying pickup ion densities with respect to the background solar wind. Since credible results profoundly depend upon an accurate realization of the shock structure, it is notable that the modeling parameters were chosen so as to reproduce the magnetic field profiles and plasma temperatures both upstream and down as actually observed by the Voyager instruments. Although the two authors address slightly different objectives, their combined results not only go a long way towards solving the velocity and dissipation conundrum, but also coincidentally provide valuable insights for our investigation as well. We shall focus primarily on a specific selection of their results, the implications in determining the electric potential gradient scale within the ramp, and the consequences for shock surfing.

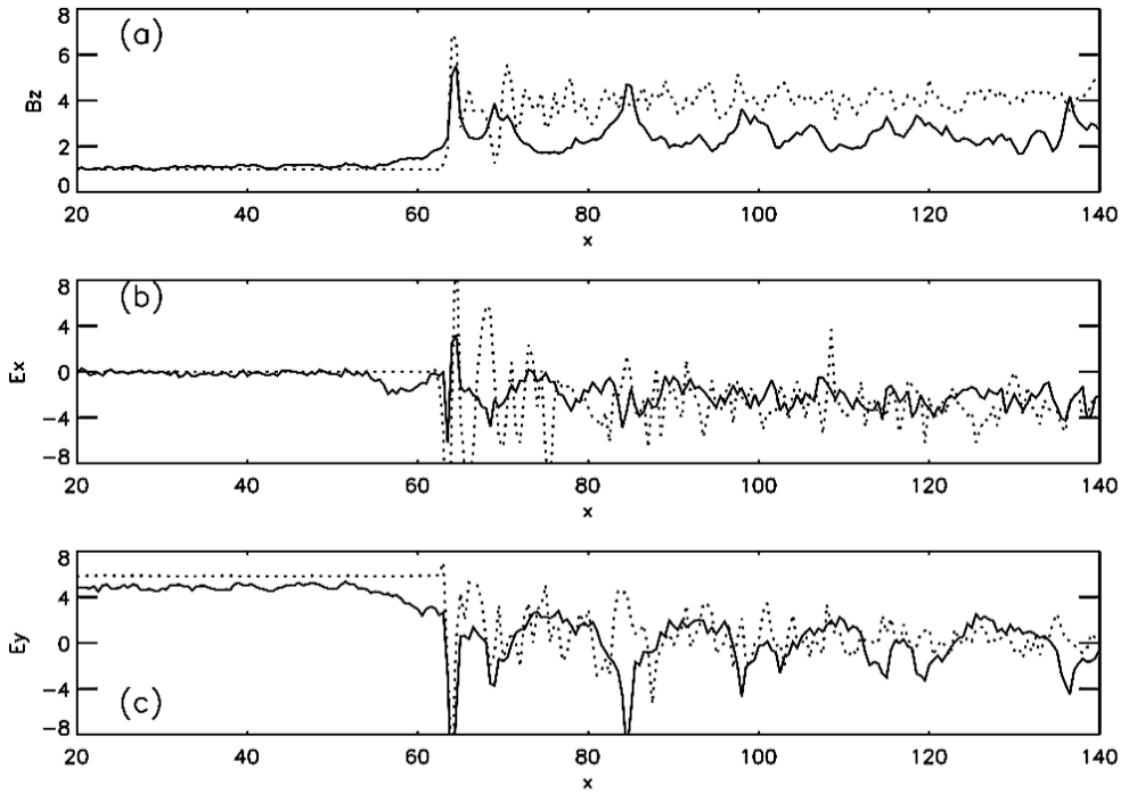


Figure 7-4 Simulated TS Field Profiles

Solid curves represent the case where PI comprise 30% of total flux, dotted curves 0%. Note the absence of a foot with no PI. The scales of most features broaden with increasing PI component. The potential ramp scale even without PI is of order $3 - 4 \lambda_{pi}$. After Winske et al (2009)

One important result is that shock features broaden with increasing PI density, which can be seen in Figure 7-4 above. Panel (a) shows the magnetic field normalized to the upstream field B_0 while panels (b) and (c) show the electric field components normalized to $B_0 v_A / c$. The two superimposed curve sets represent 30% PI population (solid) and no PI at all (dashed). Note that the PI-moderated field profiles at the ramp are also of lower amplitude. Of particular interest to our investigation is the scale of the electric potential component within the ramp, which is of order a few ion inertial lengths,

and no finer scales are observed within the region. While Wu et al concluded that, due to the large PI population, “the environment of the termination shock is very different from that of the terrestrial bow shock”, high-resolution data gathered by the Cluster mission showed that the cross-shock electric field scales were at least those of the magnetic field, and there is no reason in principle to assume that this is not a general feature applicable at the TS.

Hence, based on our velocity ceiling criterion established in baseline studies, this component of the ramp is too broad to support shock surfing. Simply applying the confinement condition $q|E_x| > qv_i B / c$ and substituting the maximum ramp field values confirms this conclusion. In reduced units employed from the figure, $B_{\max} \sim 5.5B_0$ and $|E_{x\max}| \sim 6 B_0 v_A / c$, whence $v_{\max} = \frac{E_x c}{B} \simeq 1.1v_A$ well below the shock drift velocity ceiling.

We should also point out that our trapping condition was derived assuming zero potential ahead of the ramp, while the potential gradient evident from the electric field in the foot region further degrades PI trapping at the front owing to the reduced electric field (remaining balance of the total cross-shock potential) available within the ramp itself.

There is no extended ‘foot’ in the absence of pickup ions, clearly demonstrating that the key role they play in populating that region, with scale roughly that of the upstream PI ion cyclotron, in agreement with Liewer et al (1993). Note also that the magnetic field component of the foot is roughly $\sim 7\lambda_{pi}$ in length and, with a maximum field value $\sim 7B_0$ at the base of the ramp, provides ample opportunity for shock drift to expand the PI velocity shell and thereby effectively reduce the population of surfing candidate ions.

While the studies by Wu et al focus on energy dissipation by preferential heating of pickup ions over those in the solar wind, this is not acceleration per se, as can be seen from the distribution function shown in Figure 7-5. Further, the extent of the velocity space plots shows a ceiling approximately twice the solar wind speed, consistent with that attainable by shock drift ions. Indeed, it would be surprising if evidence of shock drift acceleration were entirely absent, especially in light of the mechanism's robustness and insensitivity to variability in the electrostatic potential and shock ramp width, as we have demonstrated.

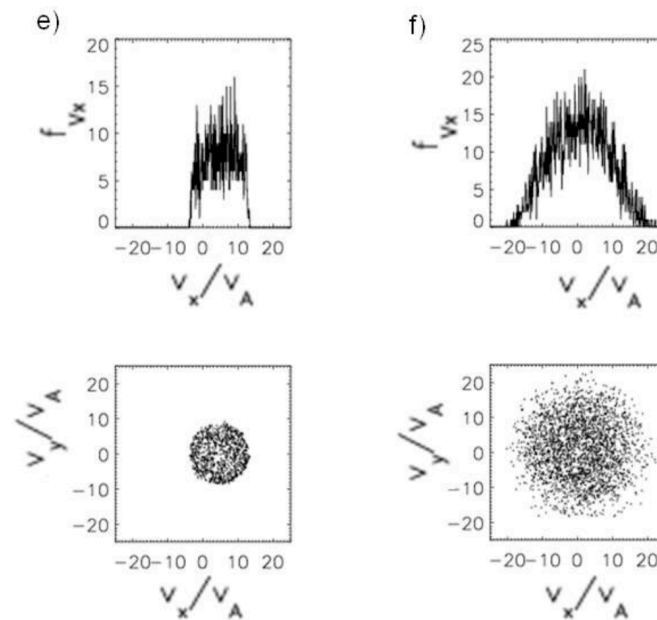


Figure 7-5 PI Distribution Functions and Velocity Space Plots

Pickup Ion distribution functions are shown in the top pair and velocity space plots below. The left panels (e) show upstream populations and the right panels (f) downstream. The shock Mach number is roughly $M_A \approx 8$ so that maximum downstream velocity is only slightly greater than twice the piston velocity. PI density is 20%. After Wu et al (2009)

Perhaps most important for our purposes is that Wu et al clearly demonstrate by varying the fractional component of PI from 0% through 30%, that the scales of critical shock features appear to soften with increasing pickup ion populations. Since the efficacy of shock surfing depends directly upon the PI fraction and inversely upon ramp width, this demonstrated dependence is diametrically opposed to efficient surfing. We are therefore in agreement with the conclusion of Wu et al that “specular reflection does not apply to pickup ions, suggesting that a different physical interpretation is needed to describe the energy gain by these more energetic ions.” The authors also reach the conclusion that “the termination shock is not a very strong Alfvénic shock such as the familiar terrestrial bow shock; rather, as our analytic model presents, the presence of the pick-up ions raises the magnetosonic wave speed so that the termination shock behaves more like a gas kinetic shock [Fisk, 1996].”

Indeed, the authors’ choice of 17% pickup density resulted in approximately 90% of the ram energy going into heating the downstream PI population, consistent with Voyager 2 observations near the shock and sufficient to damp the shock beyond the limits required for multiple ion reflection. This fact was further validated by the work of Winske et al (2008), who concluded that:

Contrary to previous models based on pickup ion energy gain by repeated crossings of the shock front (shock surfing) or due to a reforming shock front, the present simulations show that pickup ion energy gain involves a gyro-phase-dependent interaction with the inhomogeneous motional electric field at the shock.

Wu et al (2008) argued that the ramp is independent of the electron inertia length, based on earlier work by Scudder et al (1986) who claim a scale roughly that of the upstream ion gyro-radius, and the authors validated this estimate with simulations containing no

pick-up ions at all. They continue, “Intuitively, we expect that with the adding of pick-up ions at the termination shock, the shock will weaken and the ramp should widen. Further, Winske et al. [2009] demonstrated that the electric fields (and so the shock potential) at the termination shock front are steeper and stronger where there are no pick-up ions, as compared to when there are 30% pick-up ions.”

Shock Variability & Time Scales

Several time scales have either been directly observed or inferred from the data. Both the paucity of shock crossings and total absence of electric field data have necessarily led to broad speculation on at least some of these scales as well as proposals on candidate mechanisms to explain surprising observations. After waiting thirty years for the Voyager probes to reach their destinations, it is understandable that some investigators may perhaps have been somewhat ambitious and creative in drawing their inferences. Certainly a variety of opinions are extant, and follow from different interpretations of the raw data.

Shock Reformation

We reproduce below in Figure 7-6 the serial data set of magnetic field and solar wind velocity measurements encompassing what may have been three shock crossings. Burlaga et al (2008) argue from the difference in character between the profiles of TS-3 and TS-4 that the structure had significantly evolved during the intervening time of 2.7 hours or approximately $10 \Omega_C^{-1}$. They do the same for the TS-2 to TS-3 crossing interval, noting structural transformation during the period of 3.9 hours, and also describe soliton-like ‘enhancements’ in the ramp that were separated in time by roughly a single ion cyclotron period $1 \Omega_C^{-1}$. For completeness we note that, while the field profile at the TS-2

crossing is unusual, a similar structure was observed at the Earth's bow shock as reported by Bale and Mozer (2007).

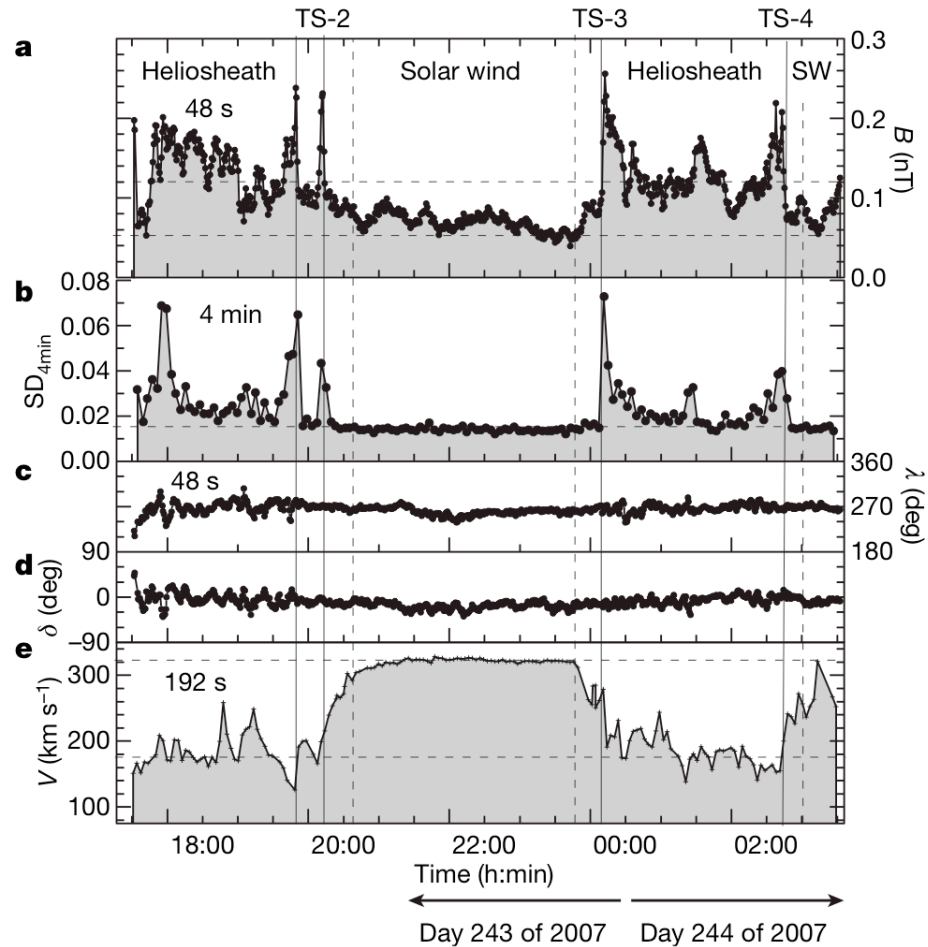


Figure 7-6 Three identifiable crossings of the Termination Shock
The top (a) and bottom (e) panels show the magnetic field and SW velocity respectively. Apparently, TS-2 and TS-4 moved inwards, while TS-3 moved outwards. After Burlaga et al, Nature 2008

It is easy to understand why there has been some disagreement in the literature with regard to the shock reformation time scale. Richardson (2008) describes the heliosheath as a “region with large fluctuations on time scales minutes to days”,

observes that the structure of TS-3 and TS-4 appear very different “only hours apart”, and notes that the latter “might be in the process of reforming”. He asserts that the plasma parameters are volatile and “the rapid variations are likely driven both by changes in the upstream solar wind and by variability in the TS”. This same analysis leads Burlaga et al (2008) to describe the TS as being “a complex, rippled, quasi-perpendicular supercritical [MHD] shock of moderate strength undergoing reformation on a scale of a few hours”. However, there is no reason in principle to assume that the shock was not reforming elsewhere and thus unobserved during the interval. Recall that shock reformation timescales at planetary bow shocks and in numerous simulation trials have been observed at or near the ion cyclotron period Ω_C^{-1} . In the same article, in fact, the authors assert “reformation is a patchy cyclic shock reformation process with a characteristic time of the order of the downstream gyroperiod”, even shorter than that upstream.

Winske et al (2009) argue that weak, low beta shocks such as the TS have been shown to be inherently unsteady, while Chapman et al (2005) concur, asserting “kinetic simulations of supercritical, quasi-perpendicular shocks yield time-varying solutions that cyclically reform on proton spatio-temporal scales”.

Shocks do not necessarily reconstitute themselves in such a way that ramps reform in the same place. During our own attempts to produce very steep ramps using finite electron mass kinetic hybrid codes, we saw instabilities manifest in the overall structure, including large secular displacements when the shock overturns and reforms. Not only did we observe shocks weaken, dissipate and reform but also simply slide upstream and down with no discernible temporal constancy. Larger choices for the anomalous resistivity appeared to destabilize the shock into more erratic modes,

although our simulations did not include self-consistent pickup ions that may have damped these effects. Hence, aside from the fact that the termination shock dynamically responds to temporal changes in the solar wind pressure, actual displacement as seen by the reappearance of the ramp may possibly have been misinterpreted.

In any case, the lack of sustained and comprehensive data at the front itself prevent us from stating with confidence that shock reformation time scales should be as expected from previous encounters within the heliosphere. Indeed, fundamental questions remain open in light of recent work on this subject, which has raised new questions not only about cyclic time scales, but also about the nature of the mechanisms involved (see, for example, Hellinger et al, 2007).

Secular Shock Motion

One often observed shock feature overlooked with respect to pickup ions is secular motion not related to undulating ripples in the front. This motion may be attributable to temporal changes in the SW pressure associated with solar variation, or to cyclic shock reformation whether quasi-periodic or otherwise. By examining the canonical pickup ion velocity shell, it should be clear that if PI perpendicular velocities within the population remain fixed, the shell position in velocity space must translate according to the piston or 'ram' velocity in the shock stationary frame, as shown below in Figure 7-7. Outward shock motion lowers the incident fluid velocity, shifting the shell to the left. All other parameters being equal, this effectively moves the surfing candidate region on the shell. Inward motion increases the ram velocity, shifting the shell to the right, resulting in an empty intersection between the surfing candidate region and the shell itself.

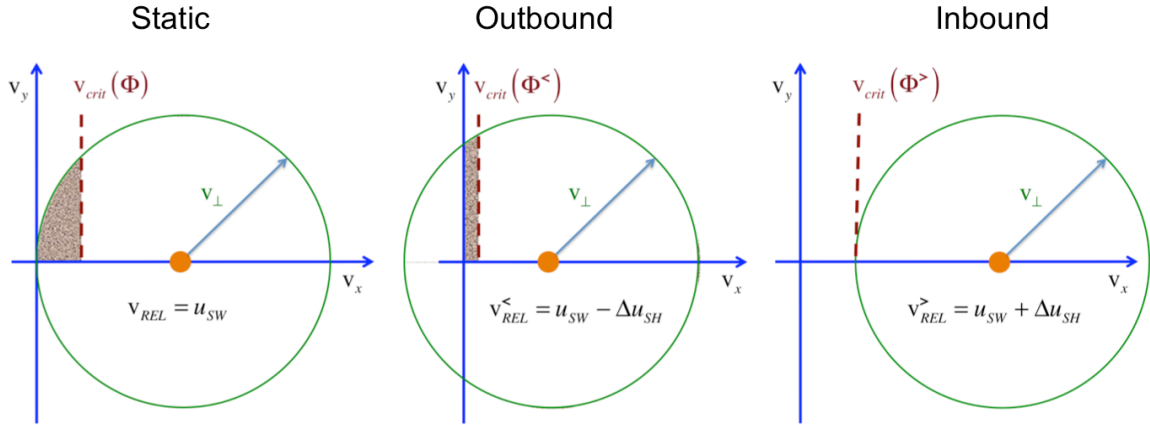


Figure 7-7 Displacement of the PI velocity shell.

Assuming constant SW velocity in the solar frame, changes in relative ram velocity due to shock motion shifts the shell position. This affects the surfing candidate (shaded) region, degrading the flux.

However, both the cross-shock potential and its gradient weaken for the outward TS-3 shock, making the ramp scale requirement even more stringent. This is most easily seen for the case of a step-function ramp, where the reflection boundary is the vertical line $v_x = v_{crit}$. Recall that the maximum normal incident velocity for ion reflection is

determined by $\frac{1}{2}mv_{crit}^2 = q\Phi$ while the potential jump condition is calculated from

$$q\Phi = \frac{1}{2}mu_{RAM}^2(1 - R_C^{-2}) \text{ so that } v_{crit} = u_{RAM}(1 - R_C^{-2}) \text{ where } R_C \text{ is the shock compression}$$

ratio, always greater than unity. Hence the shell displacement is accompanied by a fractional shift in the critical reflection boundary as well, so that secular motion of the shock front in either direction may be detrimental to shock surfing.

Citing shock displacement velocity, Richardson (2008) notes that radial front position speeds “vary from 60 - 200 km/s and the other plasma parameters are equally

volatile”. Note that the maximum velocity cited is roughly that of the SW itself, while the lower end of the range is still a respectable 20% of that value. Hence the magnitude of the displacement velocity is not trivial, and at least warrants consideration. Note further, that even small amplitude motions of the front may produce a disruptive effect on surfing ions by inducing normal oscillation in their trajectories within the ramp which, as we have shown, increase the likelihood of premature penetration of the front, thus terminating the process.

Shock Front Ripples

As we have shown from our study of corrugated shocks, ripples in the front prematurely terminate the surfing process when the ion cyclotron radius of the free unbound accelerating pickup ion exceeds the local radius of curvature.

We next examine the fine structure within the ramp as shown by the magnetic field profile in Figure 7-8 for possible evidence of such ripples. Burlaga et al (2008) suggest that this may be comparable to the “oscillating ramp containing substructures with a scale of $0.2 c / \omega_{pi}$ ” as observed at the Earth’s Bow shock. However, since this data comes from a single probe moving radially outwards through the solar wind, there may be some ambiguity in interpreting the data.

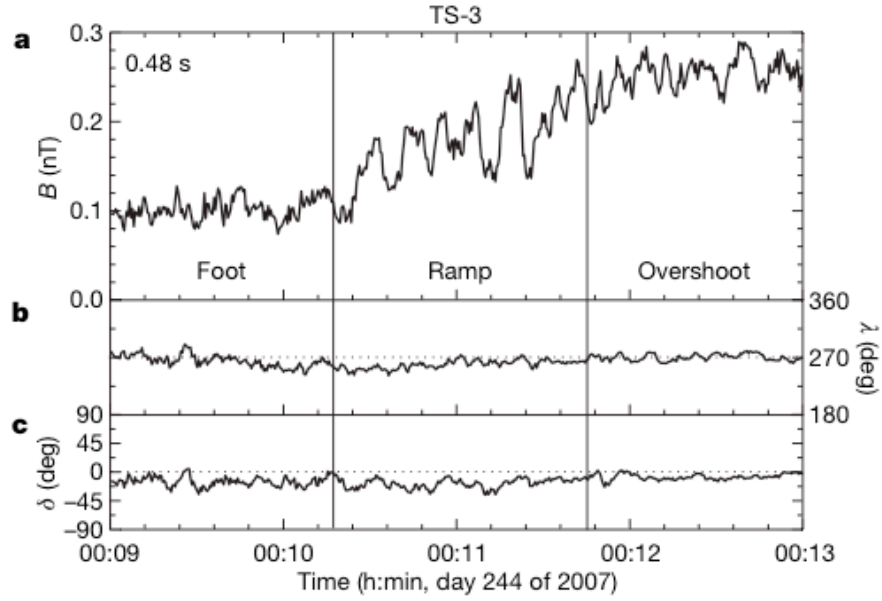


Figure 7-8 Magnetic Structure of Shock Crossing TS-3

The top panel (a) is the magnetic field strength and the bottom two represent the orientation directions. After Burlaga et al (2008)

One possibility is that the front in the region was quasi-planar, and the entire shock was oscillating as the spacecraft moved through it. Using the length scale $\lambda_{\delta B} \approx 0.2\lambda_{pi}$ and the estimated time scale $\tau \sim \Omega_c^{-1}$ yields a shock-normal oscillation velocity on the order of $0.2 v_A$, which is not unreasonable, but may nevertheless be sufficient to induce front-normal oscillation in surfing ion trajectories. On the other hand, the time scale of the broadest feature seen in the figure is roughly 12 seconds, on the order of $\tau_{jiggle} \approx 10^{-2} \Omega_C^{-1}$. This yields a front oscillation velocity $20 v_A$ or roughly three times the ram velocity, possibly inconsistent with the physical dynamics of the constituents. Such motion would be more than ample to disable any surfing ion from following the front, since the normal oscillation induced in the ion's trajectory would be

sufficient for it to penetrate the shock, while the lower estimate can still attenuate the mechanism.

Another possibility is that the fine structure itself was quasi-static, carried along by the shock motion or moving slowly through it. This ‘serrated’ magnetic ramp may be either dynamic or quasi-static, i.e., standing in the shock frame, and the potential gradient may not necessarily mimic this structure. Such fine features were observed in the electric field at the bow shock, but not in the magnetic field. Features of this character also appeared in our own kinetic-hybrid simulations but, as noted by Hellinger et al (2002) and others, re-evaluation of such results from one-dimensional simulations may be warranted. In any case, we have shown that the scale of such fine internal features cannot be used to establish criteria for shock surfing, since it is the length scale of the entire potential gradient that determines the shock-normal electric field.

Finally, the mean shock front velocity may be steady, but with ripples propagating laterally along the surface. This seems to be the most popular paradigm in the current literature. Again, due to the lack of ‘Cluster-like’ observations, we have sparse data and therefore few options for calculating the ripple parameters. Perhaps the most straightforward is to invoke time and length scales as put forth by Burlaga et al (2008). They derive the cross-shock length scale, $\lambda_{ripple} \simeq 0.15 \lambda_{pi}$ which, in conjunction with the probe’s transit time, gives a period $\tau_{ripple} \simeq 15$ s. Assuming the ripples propagate along the front at the Alfvén speed $v_A \sim 40 \text{ km/s}$, we obtain $\lambda_{ripple} \simeq v_A \tau_{ripple}$, or approximately $600 \text{ km} \simeq 0.1 \lambda_{pi}$ with corresponding wavenumber $k_{ripple} = 2\pi / \lambda_{ripple} \simeq 63$. This, as we have previously shown, is not a physically reasonable value and may be dismissed out of hand. Another alternative, perhaps more appropriate in the absence of more detailed

observational data, would be to assign the ‘typical’ wavelength scale often advanced in the literature of an upstream ion cyclotron radius, on the order of $r_g \sim M_A \lambda_{pi}$ or about 8 in units normalized for TS-3. This gives the normalized wavenumber $k_{ripple} \approx 0.8$.

We calculate the ripple amplitude under the assumption that the wave speed is much greater than the relative transit speed of the probe; otherwise we would not see several complete cycles. From the peak-to-peak amplitude of the fluctuations at approximately a fourth the total magnetic ramp, we obtain ripple amplitude

$a_{ripple} \approx 0.13 \lambda_{pi}$. Using the criterion we previously derived for shock corrugations, we calculate the maximum attainable velocity in units normalized to the shock parameters, $\tilde{v}_{max} = (mak^2)^{-1}$. Applying the ‘reasonable’ estimate for the wavenumber value, we obtain the velocity ceiling for surfing protons of $\tilde{v}_{max} \approx 13 v_A \approx 1.6 u_{SW}$, which is well below the shock drift ceiling $\tilde{v}_{SD} \approx 2.3 u_{SW}$ under the same conditions. We reiterate that this was based on conservative estimates, and should therefore reflect the most favorable case possible for shock surfing under the observed conditions.

Further support for this interpretation is indicated by Richardson (2008) who notes from analysis of the magnetic field orientation that changes in “shock directions were probably due to small-scale fluctuations in the shock”.

Summary and Conclusions

Since the Voyager missions – unlike Cluster – did not provide exhaustive, high-resolution data sets over sustained periods of time, we must be cautious in drawing conclusions from the data, especially in light of conflicting interpretations made by some observers. Consequently, although no evidence of energetic ions – either reflected or

transmitted – was detected above velocities comparable to the shock drift ceiling, we cannot conclude that unfavorable conditions must necessarily persist. However, we may at least say with some degree of certainty from the following observations and conclusions that shock surfing is unlikely to be viable at the Termination Shock:

The observed magnetic field ramp is no less than one ion inertial length, with an extended foot region several times greater, consistent with features seen at planetary bow shocks. Although no direct electric field measurements were made, simulations have confirmed that the presence of a substantial pickup ion population, comparable to that actually observed at about 17%, broadens all shock features, including the potential component of the ramp. This is also consistent with Cluster observations showing an overall cross shock potential scale no less than that of the magnetic field ramp. Further, the foot becomes more pronounced with increasing PI density and part of the potential gradient resides in that region as well. Taken together, these indicate that the cross-shock potential is too weak and the ramp too broad for shock surfing to promote ions above the shock drift velocity ceiling. Perhaps more detrimental still is the possibility that the shock-normal electric field profile may mimic the observed magnetic field profile oscillation within the ramp which, as we have shown, further degrades confinement of pickup ions at the front.

The oscillating magnetic field profile observed within the ramp seen at the third crossing TS-3 likely indicates the presence of a rippled shock front surface, and has often been cited as such in the literature. Here again, this is consistent with observations elsewhere in the heliosphere as well as results from numerical simulations. Based upon reasonable and conservative interpretation of the data, wavelength and amplitude scale

estimates show that these perturbations are incapable of supporting surfing velocities above the shock drift ceiling.

Finally, we assert that the most reasonable estimate for the periodicity of at least some variable shock features is on the order of an ion cyclotron period. Whether this represents complete cyclic overturning and reformation, ‘breathing’ (periodic broadening) of the ramp structure itself, or simply oscillation in the position of the front, any one of these alone is sufficient – on that time scale – to terminate surfing acceleration before ions can reach the shock drift ceiling.

Each of these factors separately seriously compromises the shock surfing mechanism. The low compression ratio coupled with a shallow ramp implies a weak retarding electrostatic field, the lack of a temporally rigid structure places constraints on the acceleration time, and ripples suppress both the velocity ceiling and flux of accelerated ions. Quite simply, taken together, these conditions are fatal to the surfing mechanism.

We have concluded with a high degree of confidence that shock surfing does not take place at the Earth’s bow shock and, with reasonable certainty, that it does not function at the outer planets. Although the data from the termination shock is sparse, we believe that sufficient evidence exists coupled with results from comprehensive numerical simulations to conclude that surfing is not viable there either. While this evidence is not entirely conclusive, one can at least say with a high degree of probability that Shock Surfing does not take place anywhere in the Heliosphere.

We now turn our attention outside the heliosphere and investigate whether shock surfing may play a role in acceleration processes at supernovae and other violent objects believed to be the sources of the most energetic cosmic rays.

Chapter 8 Supernovae

We have shown conclusively from high-resolution observations coupled with in-depth numerical simulations that shock surfing is not viable at the Earth's bow shock, since each of the conditions absolutely necessary for the acceleration mechanism to function is absent. We have also demonstrated that at least some of these are violated at the outer planets, and the rest are most probably so. Finally, we have argued from the few observations made that the functioning of this mechanism at the heliospheric termination shock is, if not impossible, at least highly improbable. In short, we believe that shock surfing does not function anywhere within the heliosphere. We now ask, can it function elsewhere?

Supernovae (SN) are among the most violent and energetic events in the universe. Stellar deaths can produce, for a time, energy fluxes equivalent to those of entire galaxies. Typically, after an initial neutrino burst of energy $E_\nu \sim 10^{53} \text{ erg}$ lasting about 10 seconds, an expanding shock wave drives ejecta outward with kinetic energy $E_K \sim 10^{51} \text{ erg}$. When the shock reaches a radius where the material becomes optically thin, a photon flux or Shock Breakout (SBO) lasting several days is released with total energy $E_\gamma \sim 10^{49} \text{ erg}$ (Carroll & Ostlie 1996).

Besides their enormous energy, SN shocks are believed to possess very high Mach numbers and steepen with time, making them attractive candidates for the production of cosmic rays by Diffusive Shock Acceleration (DSA). Since the DSA mechanism is believed to require a population of suprathermal ions (the so-called 'injection problem') several studies over the past decade have investigated relativistic Shock Surfing (SS) at these fronts. However, due to the power and scale of these

events, they are likely accompanied by extreme turbulence which, although beneficial for DSA, may disturb the shock front to the point where surfing ceases to function. In other words, the strong turbulence that favors DSA may disfavor SS.

Moreover, shock surfing investigations have assumed profiles and parameters similar to those locally observable without addressing their plausibility, especially in light of recent SN observations. There is no reason in principle to assume that shocks produced by SN are comparable to those in the heliosphere, at least insofar as the characteristics germane to the study of SS is concerned. Indeed, planetary shocks and the Termination Shock are standing structures, while SN shocks are rapidly and continuously expanding, and therefore change in nature over the course of time. Further, for example, a quasi-perpendicular shock is defined by the upstream shock-normal magnetic field angle that, in the case of SN fronts, lies in the exterior region of the structure, as opposed to those in the heliosphere.

While recent research has advanced our understanding of many SN features and mechanisms, none has specifically addressed either Shock Surfing or directly investigated the conditions necessary for it. Due to the absence of direct observational evidence, especially of early stage SN shock structure, we shall see that this inquiry raises many questions but yields few answers.

Although SN have been divided into several classes, for our purposes we shall divide them into two types – those resulting from the core-collapse and explosion of super-massive stars, and those associated with white dwarfs. However, regardless of the shock source, the same conditions necessary for SS to function in the heliosphere are still required. We reiterate these conditions and investigate whether there is any evidence to support them.

1. A population of upstream ions (not necessarily Pickup Ions) with small incident normal velocities in the shock frame is essential. These have been observed in the heliosphere, and the means by which they are created is fairly well understood. However, when and in what manner such a population might be created at SN is not clear.
2. The shock must be nearly perpendicular, since we have shown that fairly small deviations in the magnetic field angle suppress both the maximum attainable velocity and the flux of promoted ions. This constraint limits the regions of the shock where SS may take place.
3. A sufficiently strong and steep electrostatic potential gradient is needed, i.e., a shock-normal electric field capable of restraining ions from being swept downstream by the Lorentz force. The high Mach numbers of SN shocks ensures a large potential step, but the gradient strongly depends upon the characteristic length scale of the shock ramp, which has never been observed.
4. Since shock surfing is a coherent process, the front must be smooth and free of physical corrugations or other features that could prematurely terminate the acceleration mechanism.
5. All of these conditions must be sustained over enough time for the convective electric field to accelerate candidate ions above velocities attainable by other mechanisms, especially shock drift.
6. Since Supernovae evolve over large distance and times scales, several unique stages are apparent and the above conditions must be evaluated at each stage independently.

Light Curves, SN Types & Shock Breakout

In order to study Shock Surfing at supernovae, it is essential to know their morphologies. This involves not only the mechanics of the internal explosion and character of the expanding shock wave but, even more importantly, the medium into which it propagates. Of the few observable sources of data, light curves – integrated spectra plotted against time – as shown in the graphic shown below, have been used to classify and study SN in the early stages of their evolution.

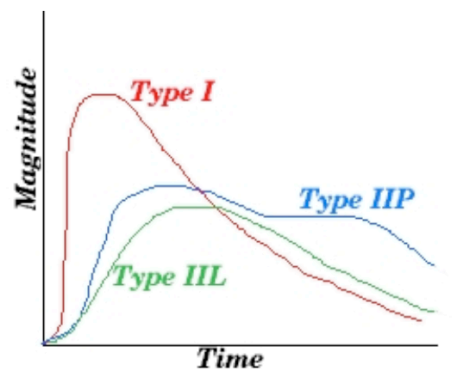


Figure 8-1 SN Light curves

As such, aside from a natural fascination with the fundamental physical processes that drive them, interest in supernovae was further stimulated by their utility as ‘standard candles’ in determining astrophysical distances. However, it is not the predictive utility of their light curves that concerns us, but rather the progenitors that govern the features of the shock front structure and the environments into which they propagate. The state of the art in observation has improved to the point where many SN have been detected at inception and, even more fortuitously, nearby events have recently occurred and their progenitors identified. As a result, it has been possible to model numerical simulations with parameters appropriate to these observed events and

be fairly confident in proposing physical models when the results agree with continuing observations. Also notable are detections by neutrino observatories which, when compared with the time delay of broadband photon observations, have provided constraints of spatial scales during the early stages of these events.

Despite the various classes into which SN have been divided, we shall concentrate on two general categories based on their progenitors – those resulting from the collapse of Massive Stars on the one hand versus explosions of White Dwarfs on the other. Prior to their demise, these progenitors create significantly different environments into which their shocks propagate, and we shall treat these separately, beginning with massive stars including the most common Super-massive Red Giants (SRG).

Temporal Stability

Since SN shocks are so different from those locally observable in the heliosphere, we cannot assume similarities between them such as reformation time scales. Indeed, the existence of precursor ‘foot’ features has been attributed primarily to the presence of an upstream pickup ion population that may not even exist at SN. To resolve this question, we examine the studies by Lee et al (2004) who note that recent observations of “X-ray and γ -ray spectra from supernova remnant RX J1713.7-3946 ... can only be explained by accelerated ions.” Under the assumption that a suprathermal ion population is required and therefore consistent with some form of pre-injection mechanism, the authors conducted “high-resolution particle-in-cell (PIC) code simulations, for parameters relevant to SNR shocks”. Of essential interest to us is that “... the structure in our PIC simulations is fundamentally time-dependent: for the

parameters of interest here, no time-stationary shock solutions exist, and instead we observe cyclic reformation on ion cyclotron time scales.”

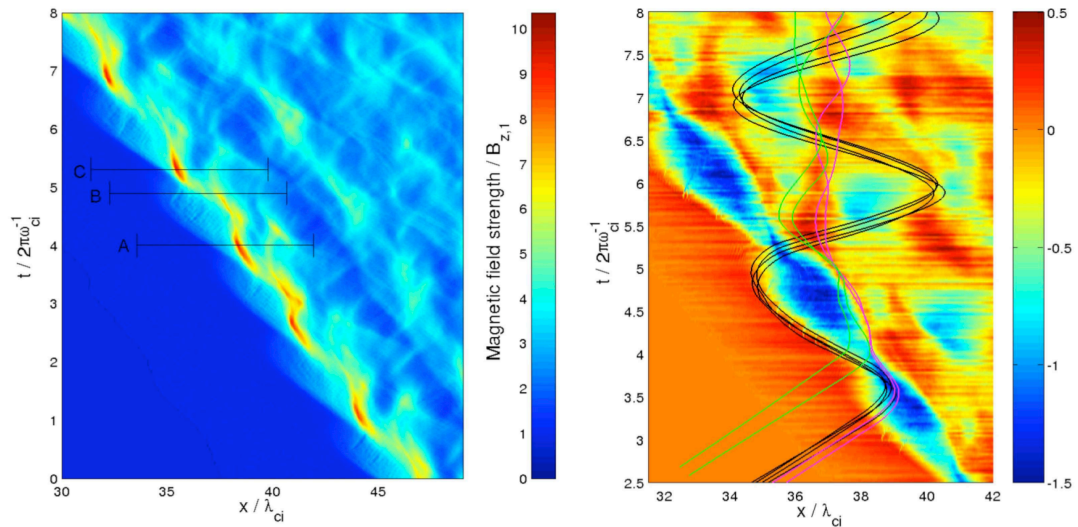


Figure 8-2 Magnetic Field (left) and Electric Potential (right)

The shock front propagates from right to left, with time on the vertical axis and shock position on the horizontal. The shock overturns and reforms on the time scale of the ion cyclotron period. Note the cyclic formation and disappearance of an extended precursor or ‘foot’. After Lee et al (2004)

It is also interesting to note that the shock velocity varies from nearly stalled to three times the average speed consistent with the Rankine-Hugoniot conditions. This maximum velocity is roughly that of the shock drift ceiling, and suggests some coupling mechanism similar to those observed at lower Mach shocks in the heliosphere, as well as those seen in numerous simulation schemes employed to model them. If there is in fact a common mechanism, we may then deduce with reasonable certainty that shock reformation times and length relaxation scales normalized to the particular shock parameters are otherwise Mach independent, and should therefore apply generally to all quasi-perpendicular shocks.

As we saw from simulations of the Termination Shock by Winske et al (2009) and Wu et al (2009), shock broadening and especially extended features such as the foot appeared only in the presence of substantial pickup ion populations. It is therefore of special note that no pickup ions were employed in these simulations. In other words, some other mechanism may act at these high Mach number shocks to enable reflection of thermal ions in the upstream medium when the shock is sufficiently steep.

While a strong cross-shock potential is to be expected at very high Mach numbers, it is the shock-normal electric field that is responsible for constraining shock surfing ions at the front, and therefore of particular interest. The composite plot shown in Figure 8-3 below illustrates, at two different stages of the cyclic shock, the magnetic field and electrostatic potential profiles, along with the ion kinetic energy and velocity space plots. The magnetic field and particle velocities are normalized to their upstream fluid values, while the electric potential and ion kinetic energy are normalized to the piston energy. As can be seen in the second row, even when the shock is at its steepest the potential is not particularly strong, surprising for a high Mach shock. As a result, the maximum velocity of accelerated particles is on the order of the shock drift ceiling, consistent with that due to a single reflection from the front, and also with the maximum shock compression ratio we have previously demonstrated. We first conclude our discussion of temporal stability and its implications for shock surfing, and then address this conundrum.

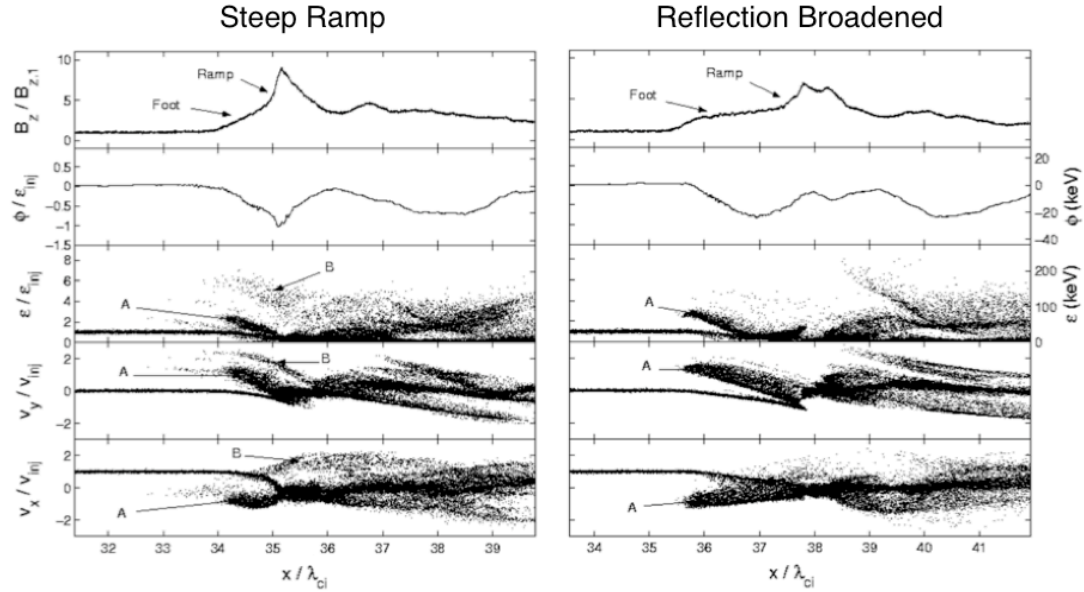


Figure 8-3 Shock Profile During Reformation

Representative cross section ‘slices’ corresponding to lines ‘A’ and ‘B’ in the previous figure. The panels show magnetic field strength, electrostatic potential, energy and velocity space plots. All shock features including the foot broaden to a few ion-cyclotron radii due to reflected ions.

After Lee et al (2004)

These simulations strongly suggest that cyclic reformation may be a general feature of high Mach shocks. When the shock is sufficiently steep to reflect upstream ions, they form extended precursors regions that weaken the potential gradient and ultimately degrade the potential step as well. This means that the shock-normal electric field cyclically weakens and reforms in response to such reflection, making the overall process self-limiting to the point where shock surfing cannot function beyond a single reformation period, if at all.

As we have previously shown, shock reformation or even cyclic ‘breathing’ of the ramp with period Ω_c^{-1} impose acceleration timescales sufficiently small to limit the

maximum surfing velocity to that of shock drift acceleration – at most a few times that of the upstream plasma velocity. Hence, even with an abundant and suitable population of candidate ions in the presence of an otherwise smooth shock front, the temporal constraints imposed from the simulation results are, by themselves, enough to render surfing untenable at SN shocks.

Cross-Shock Potential

Observations of the Doppler shifted emission lines of both young remnants and supernovae in progress indicate a typical early stage shock velocity on the order of $v_{shock} \sim 10^{-1} c$ or less, so that the shock itself is not relativistic. In order to determine the shock Mach number, we must know not only the absolute shock velocity, but also the velocity of the existing circumstellar medium (CSM) into which it propagates. We show in the figure below, results from simulations conducted by Vikram Dwarkadas (2006) who modeled the winds of massive stellar progenitors over their lifetimes. Until the very end, these winds themselves propagated outwards at approximately $v_{wind} \sim 10^{-2} c$. The figure also shows the significant mass loss that these stars undergo and, taken together, the two plots suggest that just prior to the core-collapse and resulting explosion, the ram pressure may be nearly constant in time. Since the extent of the wind-blown ‘bubble’ is on the order of a parsec, it is clear that shock will rapidly overtake the nearby region where the velocity and density are not constant. At one-tenth the velocity of light, the shock will reach the dense shell of the CSM within approximately 10 years, and its expansion thereafter will evolve according to the parameters of the interstellar medium beyond.

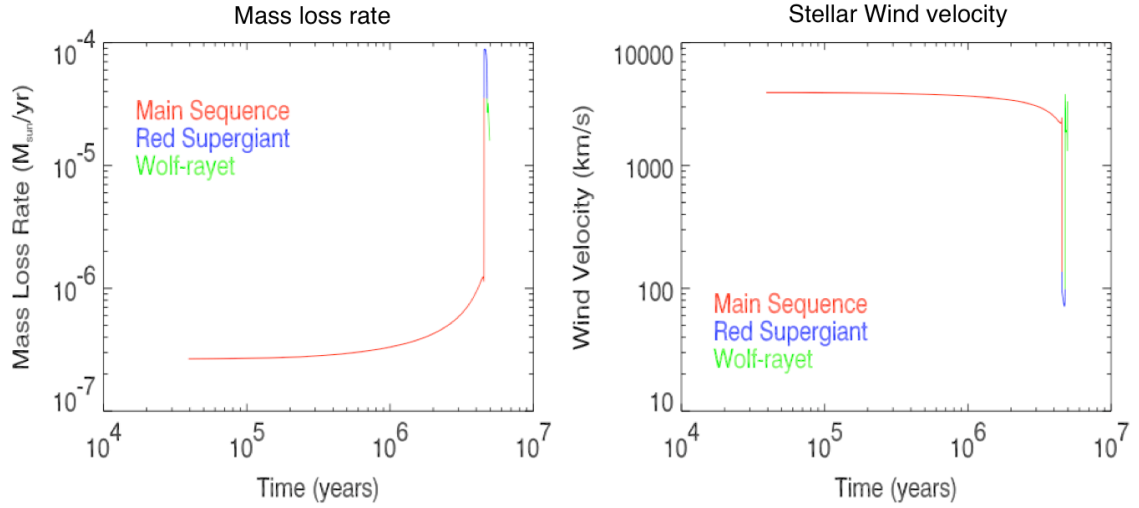


Figure 8-4 Pre-Supernova Progenitor Wind Parameters
Mass loss rate (left) and Stellar Wind velocity (right) are shown
After Dwarkadas (2006)

Since the wind speed within the entire CSM is so far below the shock velocity itself, we may treat existing stellar wind as stationary in calculating the shock potential. However, the Mach number requires the magnetic field strength in order to compute the Alfvén speed and, to the best of our knowledge, this value remains an open question. However, let us proceed with an estimate that the Alfvén velocity in the regions of interest lies in the range $10^{-4}c < v_A < 10^{-3}c$, which gives Mach numbers between $10^2 < M_A < 10^3$.

Recall that the electrostatic potential may be estimated from the kinetic energy difference between the upstream and downstream fluid flows. Since the shock compression ratio R_C is a weak function of the Mach number, we may further approximate the potential simply by $\Delta\Phi \approx \frac{1}{2}M_A^2$ in units normalized to the local plasma parameters. This implies an immense electrostatic potential – on the order of

$10^4 < \Delta\Phi < 10^6$, whose electric field can easily overcome the Lorentz force, given a sufficiently thin ramp. Recall that the condition necessary for surfing ions to exceed the

shock drift velocity ceiling is $qE_{\perp} > q\frac{v}{c}B_0$ which, along with the SD velocity ceiling

$v_{SD} \sim 3M_A v_A$ gives $L_r < \frac{1}{6}M_A \lambda_{pi}$. However, we note from the examples given by Lee et

al (2006) that the scale over which the majority of the cross-shock potential occurs at SN

shocks is on the order of the cyclotron radius λ_{ci} of reflected upstream ions, which is

proportional to the Mach number $\lambda_{ci} \sim M_A \lambda_{pi}$. Hence the ramp width $L_r \sim \lambda_{ci} \sim M_A \lambda_{pi}$ is

almost an order of magnitude too large to satisfy the surfing constraint criterion.

We also have reason to question whether such high Mach numbers can actually

be attained. Recent theoretical work supported with observational evidence by Sorasio

et al (2006) has demonstrated that an upper limit may exist for the shock Mach number.

They argue that the charge separation that forms at the front cannot be sustained for all

electron temperature and pressure gradients. They conclude that the system is

supportable only when “the electron pressure exceeds the ion pressure along the shock,

and if both coincide in value and slope at the [shock-normal front] boundaries” while at

the same time, “if the electrostatic potential exceeds the critical value $\Phi = M_A^2 / 2$, the

ion pressure becomes imaginary” and the wave breaks. Hence they assert the existence

of a maximum attainable Mach number $M_{\max} \simeq \frac{3(\Upsilon-1)}{\Upsilon} \sqrt{\frac{\pi\Theta}{8}}$ where $\Upsilon = \frac{n_{e1}}{n_{e2}}$ is ratio of

electron densities and $\Theta = \frac{T_{e1}}{T_{e2}}$ is the ratio of electron temperatures.

In addition to this, analysis from recent observations has suggested that the picture of a SN shock as a dense expanding shell with a hot, sparse interior may be flawed. Rather, it has been suggested that, while the shock front does sweep up some of the CSM through which it passes, the aggregation is not complete but rather leaves a substantial portion of the heated material behind. If this is so, then the cross-shock potential cannot be calculated in the manner we describe above, and we must take this material into account as if it were a different species. Perhaps this may be observable in the future, if it becomes possible from the spectra of thermal emissions to distinguish between the blue-shifted component of the expanding shell and the material left behind.

Finally, we claim that during the ‘shock breakout’ stage, the intense photon flux may significantly broaden the ramp. As we show below, later stages in the evolution of SN shocks exhibit severe turbulence and deformation of the shock front, leaving only the early stages smooth enough to allow shock surfing. Although derivation and discussion of the pertinent parameters are beyond the scope of this investigation, such broadening of ionization fronts has been observed in the simulations of Whalen and Norman (2008).

Smoothness

We have conclusively shown that ripples or corrugations in the shock front disrupt the surfing mechanism, which terminates when the cyclotron radius of the accelerating ion exceeds the local radius of curvature. Therefore, it is essential to know the extent to which instabilities or asymmetries may perturb the shock front and so disable shock surfing. It is well known that SN explosions are commonly asymmetrical, since many of the compact objects formed from the cores are ‘ejected’ with substantial peculiar velocity relative to their remnants. Some SN have a preference for emitting

ejecta along the polar axis, while others are believed to be aspherical due to uneven nuclear burning during the explosion. That many supernovae shocks are asymmetrical is well established and attributed to several independent factors. Radio observations of many SN remnants show "barrel shaped" morphologies (Choe & Jung), while observations of Wolf-Rayet stage progenitors, e.g., Eta-Carinae, clearly show bipolar distributions of their rapid mass loss. Many White Dwarfs are known to emit ejecta preferentially along the rotation axis, and any SN shock front expanding from these progenitors along the polar axes must be influenced by the asymmetry.

Blondin et al (2006) have addressed the observed distribution of pulsar rotation rates, not adequately explained by previous theoretical models. The authors performed numerical simulations to investigate the evolution of SN accretion shocks, discovered a 'robust instability' and concluded that a "new mechanism for the generation of neutron star spin... weakens, if not breaks, the assumed correlation between the rotational periods of supernova progenitor cores and pulsar spin." Of particular interest to our investigation of shock surfing is the influence that such instabilities might have on the symmetry of the expanding shock wave, and especially should they induce corrugations upon the smoothness of the front. The figure below shows the evolution of one such SASI or "Stationary Accretion Shock Instability" propagating along the shock surface.

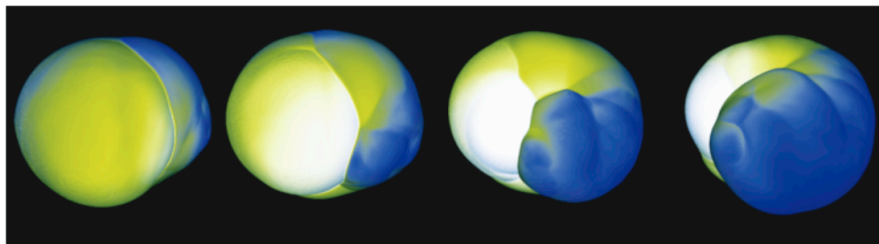


Figure 8-5 SN Accretion Shock Instability
After Blondin et al (2006)

Complex structures have long been observed in mature remnants along with more recent observations of young SN remnants, and four such examples are shown in the figure below. On the top left is Tycho with Kepler's beneath. Both appear almost spherical but upon close inspection exhibit considerable turbulence in their structure. On the top right is N49 with The Crab nebula below. The complex features of these two are more apparent, especially Rayleigh-Taylor 'fingers' emanating from within.

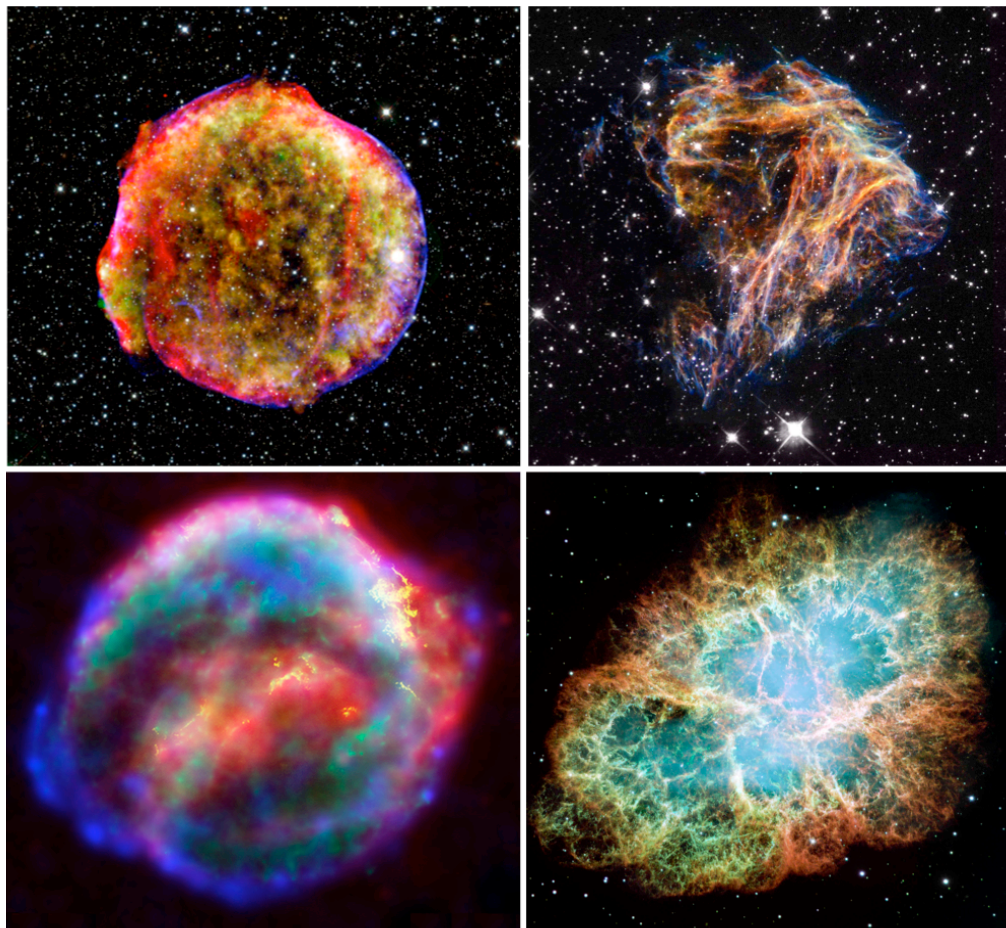


Figure 8-6 Supernova Remnants

Clockwise from top left are Tycho, N49, Crab Nebula and Kepler
After Various Sources

Such cases are numerous enough to avoid dismissal, and it cannot be assumed out of hand that a simple spherical structure will provide a suitable model for a generally applicable morphology. In addition to large-scale geometric features, evidence of turbulent deformation abounds from observations of these remnants. Whether they form at early or later stages raises questions regarding the mechanisms responsible, and whose solutions are pertinent to shock surfing.

With the advent of substantially more powerful computational resources, investigators have directed their attention towards numerical studies of these questions. Recent improvements in observational fidelity and detail not only provide the means for validating numerical results, but also the data required for estimates on fundamental progenitor and shock parameters.

Early numerical simulations indicate that initial perturbations in Type II SN lead to clumping and fast knots (Luo & Chevalier, 1994). In contrast to low compression-ratio shocks, Jun & Norman (1996) show interior fingers and bumps, while high compression-ratio shock simulations by Blondin & Ellison (2001) show that “enhancement of the radial magnetic field from Rayleigh-Taylor instabilities... extend all the way to the shock front, and this may help explain the slight dominance of radial fields long seen in polarization measurements of young remnants like Tycho.” These turbulent structures are illustrated in the figure below, which were obtained using different values of the polytropic index, while the following figure validates the existence of these instabilities independent of the dimensionality of the simulation.

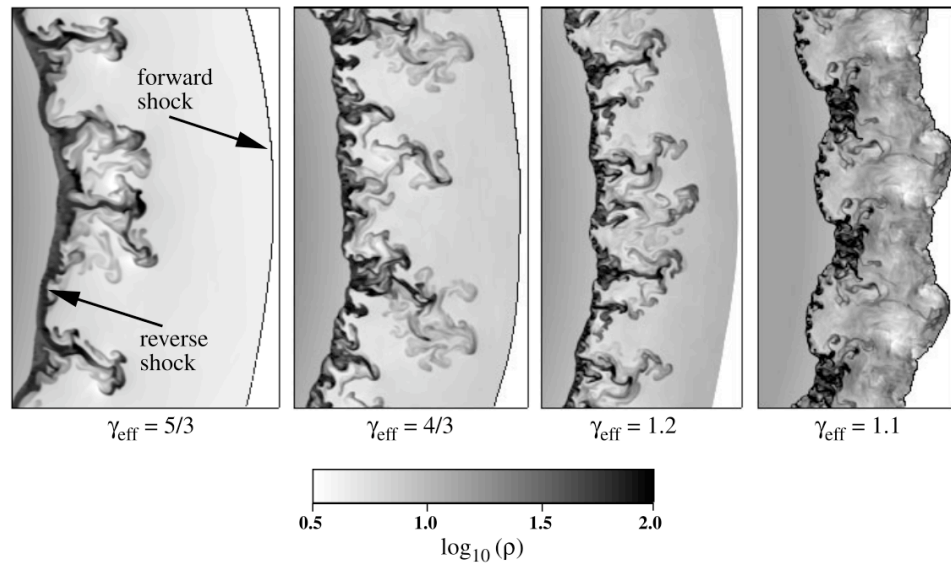


Figure 8-7 Density Plot of Rayleigh-Taylor Instabilities

Shocks propagate outward towards the dense shell of the circumstellar medium. Each panel shows the simulation result for a different value of the polytropic index. After Blondin et al (2001)

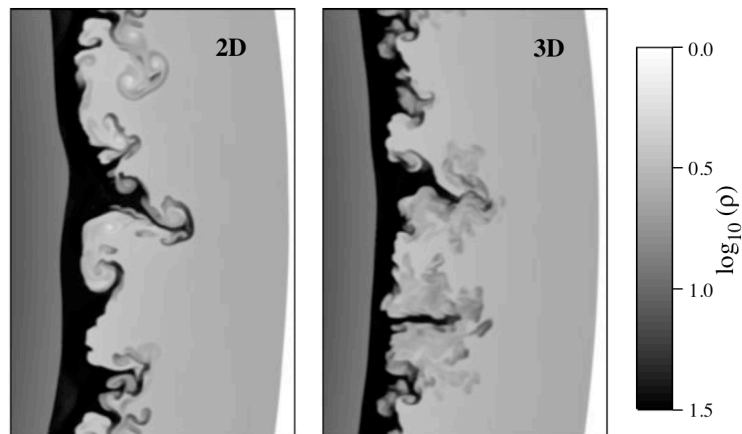


Figure 8-8 2D versus 3D Simulation Comparison

Dark areas represent the shocked ejecta while light areas to the right indicate the unshocked pre-existing circumstellar medium into which the shock propagates. After Blondin et al (2001)

More recent simulations of standing accretion shocks by Blondin et al (2003) show that “relatively small-amplitude perturbations, whether they originate from inhomogeneities in the infalling gas, aspherical pressure waves from the interior region, or perturbations in the postshock velocity field can excite perturbations in a standing accretion shock that lead to vigorous turbulence and large-amplitude variations in the shape and position of the shock front.”

Some of the most comprehensive studies to date have been performed by Vikram Dwarkadas, whose multi-dimensional simulations reveal “the presence of several hydrodynamic instabilities... The turbulent interior, coupled with the large fluctuations in density and pressure, gives rise to an extremely corrugated SN shock wave.” In addition, reverse shocks are generated which then bounce off the core and propagate outwards again ultimately forming a family of transmitted and reflected shocks within the CSM. Density fluctuations within the initial wind bubble ultimately break whatever spherical symmetry may have originally existed, leading to non-radially intersecting shock fronts and creating severe wrinkles in the structure.

Owing to the long duration of SN events, the distinction between early and late stages of the shock evolution may be essential in determining when, if at all, surfing is possible. This study is particularly valuable for our purposes since it shows that instabilities inherent even in the core collapse result in conditions throughout the lifetime of the shock that are unfavorable for shock surfing. This is illustrated in the next two figures, which show the earliest emergence of the shock from the stellar envelope, through an intermediate “Wolf-Rayet” stage, and finally the later stage of its evolution towards the outer shell of the CSM.

Figure 8-10 graphically illustrates Rayleigh-Taylor ‘fingers’ of extremely large extent evident from the shell radius at 80 parsecs. These create critical problems in terms of identifying the local magnetic field orientation and possible locations of quasi-perpendicularity needed for shock surfing. While we are not able to quantify the length scales of small perturbation amplitudes or wavelengths, especially at the late stages, Figure 8-8 demonstrates that such structures should persist at smaller length scales. Recent work by Hartigan et al (2009) comparing 3D simulations with experimental data has confirmed that fine turbulent substructures exist around such fingers that are not seen computationally.

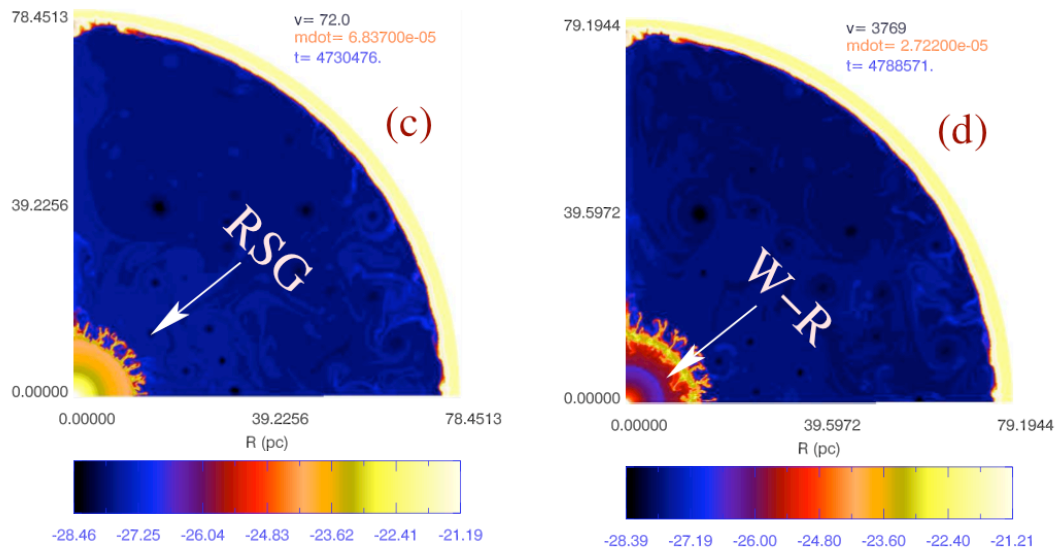


Figure 8-9 SN Shock Wave Evolution

The shock propagates into the progenitor’s circumstellar bubble of radius $R \sim 80$ pc. The left panel shows the shock emerging from the Red Super Giant envelope, while the right shows evolution into the Wolf-Rayet stage.

After Dwarkadas et al (2007)

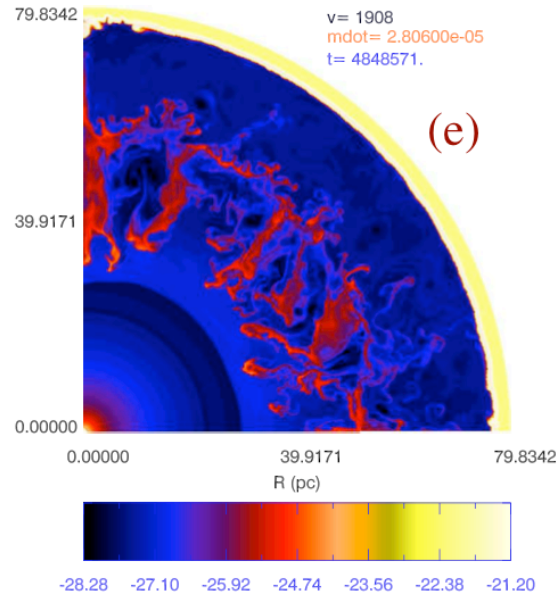


Figure 8-10 Late stage SN Shock Wave

The shock turbulence is fully developed, and the Rayleigh-Taylor ‘fingers’ are about to pierce and completely disrupt the circumstellar shell.

After Dwarkadas et al (2007)

We believe that the results of these studies demonstrate that turbulence at all stages of the SN shock evolution is alone sufficient to prohibit shock surfing. The validity of these simulations is most clearly evident in the morphologies of SN remnants; the asymmetry and filamentation visible in remnants such as the Crab Nebula show that the features seen in the figures above continue to evolve, and that turbulent conditions unfavorable for shock surfing persist well after the shock has evolved beyond the CSM shell. We refer the reader to the website of Vikram Dwarkadas, <http://astro.uchicago.edu/~vikram/sncsm.html>, and links therein for additional examples and discussion of Cas A, SN1987A and W44.

Surfing Candidate Ions

We have given several reasons to doubt whether conditions conducive for shock surfing may briefly exist during the evolution of a supernova shock. However, in the spirit of completeness, we now ask whether a population of Pickup Ions or other suitable surfing candidates might exist.

It should be patently evident that, in contradistinction to the standing shocks within the solar CSM, pickup ions produced by photo ionization and carried away in the progenitor's stellar wind could never be stationary in the reference frame of the rapidly expanding supernova shock front, so that some other mechanism would be required to produce such a candidate population.

One possibility is that some of the ejecta driven outward beyond the shock front may assume guiding center trajectories suitable for SS injection. Another is that the large photon flux emitted during the early stages when the shock first becomes optically thin (Shock Breakout) may be capable of accelerating ions in the shock precursor region to sufficient velocity by means of photon pressure. A third possibility is that accelerated electrons believed to stream outwards through the front may accomplish this purpose either through induced heating, Coulomb collisions or waves generated by their passage through the precursor.

Certainly shock surfing would be limited to very early stages but the detailed observational evidence required to evaluate these alternatives is sadly missing. However, some conjecture may be possible from recent advances.

Shock Breakout

Simulations show and observations confirm that, as the radiation-dominated shock wave from core collapse approaches the stellar surface, the optical depth of the plasma ahead of the shock decreases until radiation escapes in a burst. This Shock Breakout (SBO) occurs days or weeks before the optical light from radioactive decay peaks and can be used to determine the radius of the progenitor star and possibly its recent mass loss history (Frey et al 2009).

Estimates as low as a few astronomical units (AU) have been made from the delay between neutrino detection and that of optical SBO, such as SN1987A (Carroll & Ostlie). Observations of SNLS-04D2dc with the Galaxy Evolution Explorer space telescope revealed a radiative precursor from the shock before it reached the surface, theoretically fitted UV light curves by Schawinski et al (2008) confirm a Red Supergiant progenitor, consistent with SBO occurring at a radius of several AU.

Since typical peak SN luminosity is of order $L_{peak} \sim 10^{43} \text{ erg s}^{-1}$, we calculate the energy flux at a radius of several AU to be on the order of $10^{12} \text{ Watt / m}^2$, a few orders of magnitude greater than at the sun's surface. However, Sironi (2009) has shown that the initial flux consists of thermal X-rays which, at $\lambda = 1500 \text{ \AA}$ have insufficient energy to ionize atomic hydrogen. These same studies and related simulations indicate that the shock front initially remains shallow and only steepens much later as multiple internally reflected shocks converge, while simulations by Ensman & Burrows (2002) showed that shock becomes radiation-damped at peak SBO.

In short, while new insights continue to emerge, many questions remain unanswered. Do these latter studies imply that shock surfing is not viable at early SN

times? Do the results of instability studies disfavor later stages? If the developed turbulence is chaotic, does it possess a fractal nature so that observed filaments also appear at smaller scales?

We showed in baseline studies that, for SS to function efficiently – all others factors being favorable – the upstream magnetic field into which the shock propagates must be very close to the perpendicular. Whether such regions exist depends primarily upon the magnetic field structure of the pre-shock stellar wind and its composition. What is the magnetic field strength of the stellar wind? Does the CSM possess a current sheet? What is the ionization state?

These are serious questions we have not fully addressed but mention in passing – all these and more must be answered in order to determine with certainty whether the necessary parameters are suitable for shock surfing.

Summary

Although SN produce extremely strong shocks, the conditions necessary for shock surfing do not appear likely to exist there. We have shown that SN shocks reform cyclically on spatio-temporal scales, and that the period of these oscillations is incompatible with efficient SS acceleration. We have argued that the potential gradient is in all probability too wide to support a strong enough shock-normal electric field even if the Mach number is high, and that early stage opportunities for SS may be quelled by shock broadening via SBO. We have also presented evidence of the growth of violent instabilities, and argued that these are likely accompanied by perturbations in the shock front also unfavorable for shock surfing. Finally, we question the existence of an effective mechanism for creating a population of candidate ions capable of surfing.

Supernovae are not the only energetic objects in the universe. Much attention has been given to relativistic jets from pulsars and active galactic nuclei and the shocks they produce. We have not addressed relativistic effects at SN, simply because the shock themselves are not. However, these are fascinating objects which have been studied both theoretically and numerically, most recently in a series of papers by Niemiec and Ostrowski (2006) who conclude that “emission observed from astrophysical sites with relativistic jets, e.g., active galactic nuclei and gamma-ray bursts, is likely generated by particles accelerated in processes other than the widely invoked first-order Fermi mechanism”.

Indeed, long accepted models such as DSA are being questioned, as in the case of CygA by Stawarz et al (2007), who argue for the existence of different regimes, divided by a critical energy “above which Fermi-type acceleration processes may play a major role, but below which the operating acceleration mechanism has to be of a different type.” The most recent advances in this area of observational astronomy have come from the Fermi Gamma-ray Space Telescope, formerly known as GLAST. New insights into the acceleration mechanisms responsible for creating cosmic rays and the morphologies and physics of the objects that produce them have begun to appear in the literature, e.g., Abdo et al. 2009 and Weltevrede et al 2010.

The fact is that ample evidence exists to identify these energetic objects as the sources of cosmic rays. However, there is none to suggest that shock surfing does or even can play a role in the processes that produce them.

Epilogue

We have presented a body of evidence – perhaps in some areas anecdotal – to argue that shock surfing is not supportable at energetic objects responsible for high-energy cosmic ray production. Indeed, due to the strong turbulence believed to exist in these environments, Shock Surfing as a pre-injection mechanism may not be necessary. Rather, it was primarily conceived to explain the conundrum of anomalous cosmic ray production in the heliosphere, and we have demonstrated beyond reasonable doubt that shock surfing is not viable there.

That some mechanism or mechanisms must be responsible for promoting ions to high energies is manifestly evident both from locally detected fluxes in the heliosphere as well as observations of distant objects. We intend to continue our investigations in that area with subsequent studies of compact objects e.g., Cygnus A, which offer interesting possibilities for energization via the Blandford-Znajek process and variations on that mechanism. Jets and beams also offer fertile ground for further study. Perhaps the most fruitful investigation would be fully self-consistent three-dimensional simulations of Supernovae where the merging of multiple internal shock trains may offer opportunities for confinement and compression.

If it appears that we have produced an over-abundance of arguments contrary to the functioning of Shock Surfing, we wish to stress that this was not motivated by malicious intent, but conducted rather in the spirit of providing ample analytical machinery with which to evaluate this mechanism under the manifold conditions that accompany the broad variety of shocks existing in the universe.

Having put Shock Surfing to permanent rest, we cannot help recalling J. Robert Oppenheimer's remark at the first Trinity Test, which we may paraphrase thus – "Now I am become Shiva, Destroyer of Theories". During our investigation of Shock Surfing, we necessarily explored areas far afield from this specific topic and, in so doing raised serious questions about the current models of velocity diffusion, specifically with respect to the Diffusive Shock Mechanism itself.

Rather than continuing on a path of destruction and placing ourselves open to the same criticism, we look forward instead to creating, and intend to continue our investigations with a view towards establishing a new mathematical foundation for turbulent velocity diffusion.

Appendix A The Adiabatic Invariant

We review the theoretical foundations of Shock Surfing (SS) as advanced by Lee et al (1996) hereinafter referred to as ‘Lee’ by comparing our baseline results with their predictions for the case of an infinitely thin shock ramp. Since our particle pusher was written strictly to accommodate finite ramps, we set the width to a value well below that reasonably expected to be found at any real physical shock, roughly half an electron inertial length, $L_r = 0.01 \lambda_{pi} \approx 0.5 \lambda_{pe}$. According to Lee, the adiabatic invariant $\oint v_x dx$ should hold, where v_x and dx are the shock-normal particle velocity and differential displacement integrated over a single bounce period. As seen in Figure A-1, the value is not strictly invariant, but rapidly becomes a slowly decreasing function that asymptotes (as seen in log-log plots) to $\sim n^{-2/3}$ after bounce number $n = 500$.

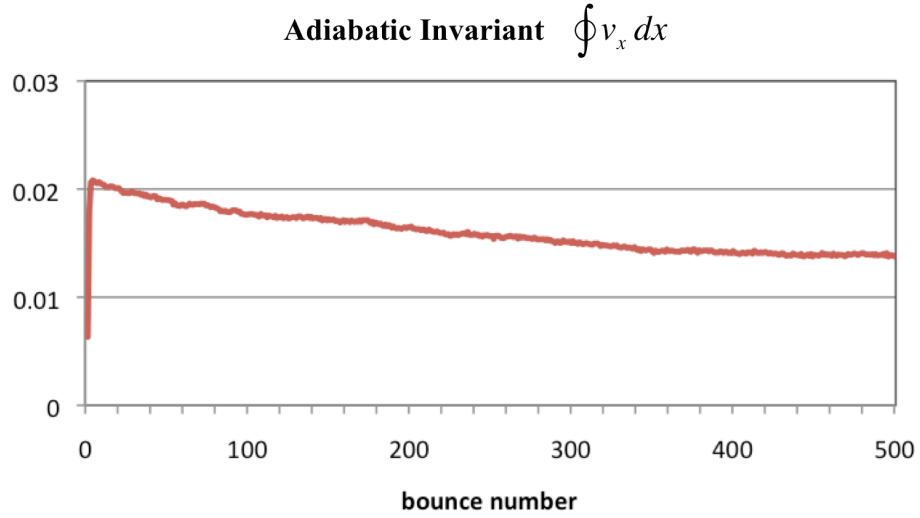


Figure A-1 The Adiabatic Invariant, or normal action integrated over a bounce period. After the first few, the function becomes smooth, gradually decreasing monotonically

We next examine the particle's kinetic energy as a function of bounce number, and see in Figure A-2 that the energy growth is almost linear. The longitudinal displacement Δy along the shock front (in the direction of the convective electric field) between reflections must be nearly constant, as the change in kinetic energy is a direct function of that displacement i.e., $\Delta\epsilon = qE_{\parallel}\Delta y$.

Since these particles represent protons, the energy is normalized to the solar wind piston energy. Initially, the function is linear with $\tilde{\epsilon} = \epsilon / \epsilon_0 \sim 2n$ where the energy gain per bounce is **twice** the piston energy, and the function asymptotes at bounce number $n = 200$ to approximately $\Delta\epsilon \sim 1.5\epsilon_0$. Note that, in the limit of a vanishing ramp width, Lee predicts normalized kinetic energy gain per bounce of unity.

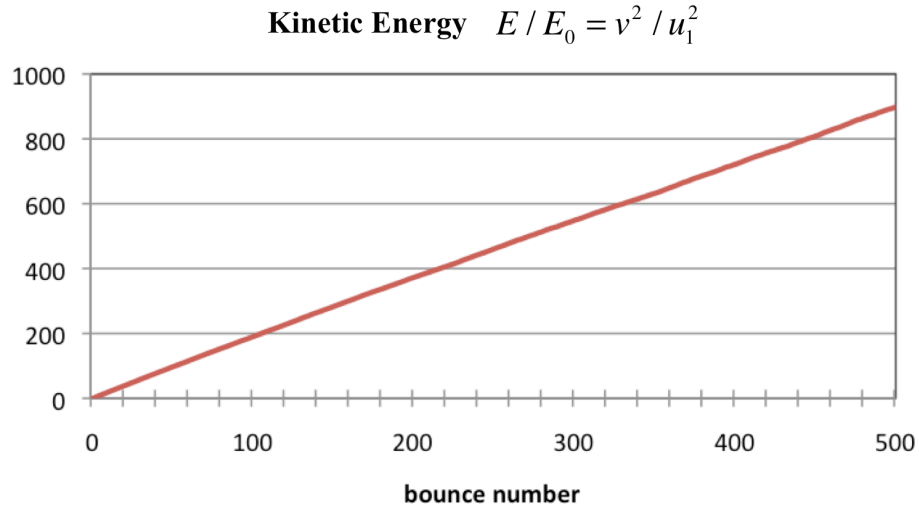


Figure A-2 Kinetic Energy

Energy grows linearly, indicating fixed displacement between reflections.

Another requisite condition for the theory to hold is that the normal bounce amplitude should monotonically decrease with bounce number. We show in Figure A-3 that amplitude normalized the ion internal length c / ω_{pi} and, for the sake of comparison

below, we approximate the asymptotic limit of the function as having roughly negative two-thirds dependence, i.e., $\Delta x(n) \sim n^{-2/3}$.

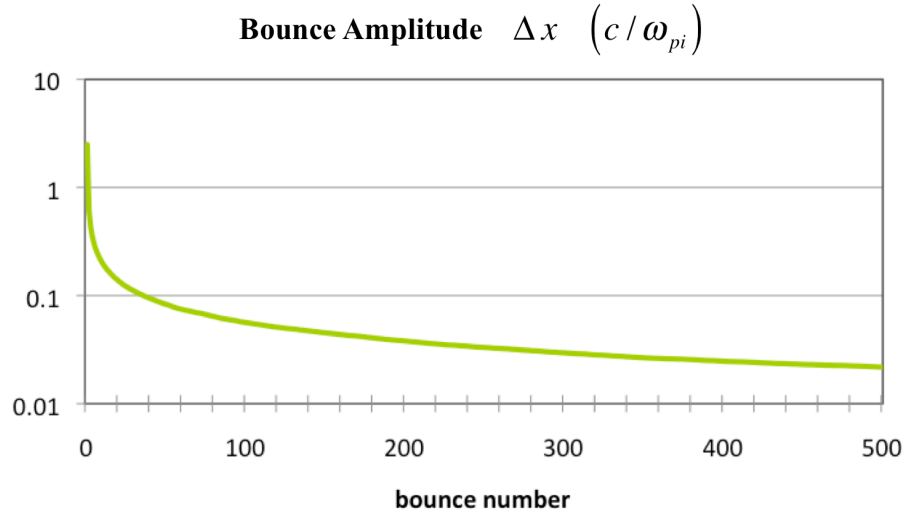


Figure A-3 Bounce Amplitude, or maximum normal excursion after reflection. After settling at 100 reflections, the function asymptotes to

$$\Delta x(n) \sim n^{-0.62} \text{ or roughly a } -2/3 \text{ power law}$$

Finally, we show in Figure A-4 evolution the bounce period. Again, for the sake of ready comparison, we approximate the asymptotic limit of the function as $\tau(n) \sim n^{-1/2}$, where the time is normalized to the ion cyclotron period.

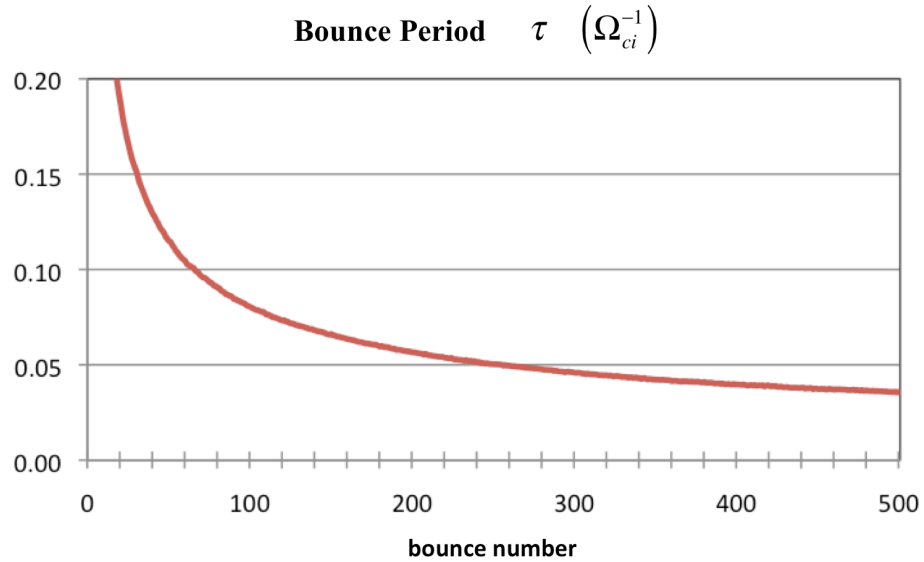


Figure A-4 Ion Bounce Period

After 100 reflections, the function asymptotes to $\tau \sim n^{-0.54}$

Deviation from the Thin Ramp

As can be seen in the combined plots of Figure A-5, the ‘adiabatic invariant’ is still smooth but no longer constant. At least one can say that an asymptotic region must be attained much later with respect to bounce number as the ramp width increases. Note that the range of ramp widths plotted is quite small – less than four electron inertial lengths – in other words, the function deviates from the ideal rather quickly.

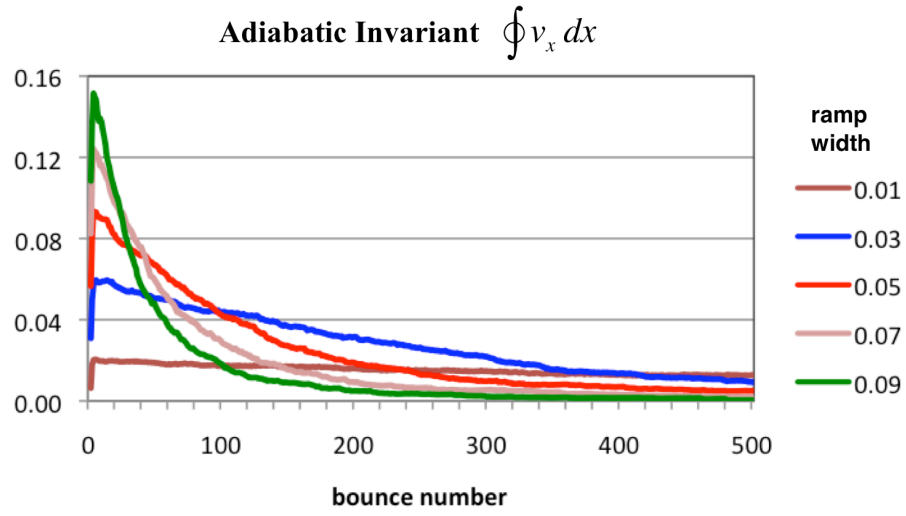


Figure A-5 Deviation of the Adiabatic Invariant

The function diverges from its ideal with increasing ramp width

Figure A-6 further illustrates that the function may be expressed as a power law in bounce number $I = kn^p$ with exponent $p \simeq -2/3$. For the shallow ramps $L_r = \{0.01, 0.03\}$, less than or equal to the electron inertial length c / ω_{pe} , the constant $k \sim 1$, which diverges from unity with increasing ramp width $k(.05) \simeq 0.355$, $k(.07) \simeq 0.126$ and $k(.09) \simeq 0.027$.

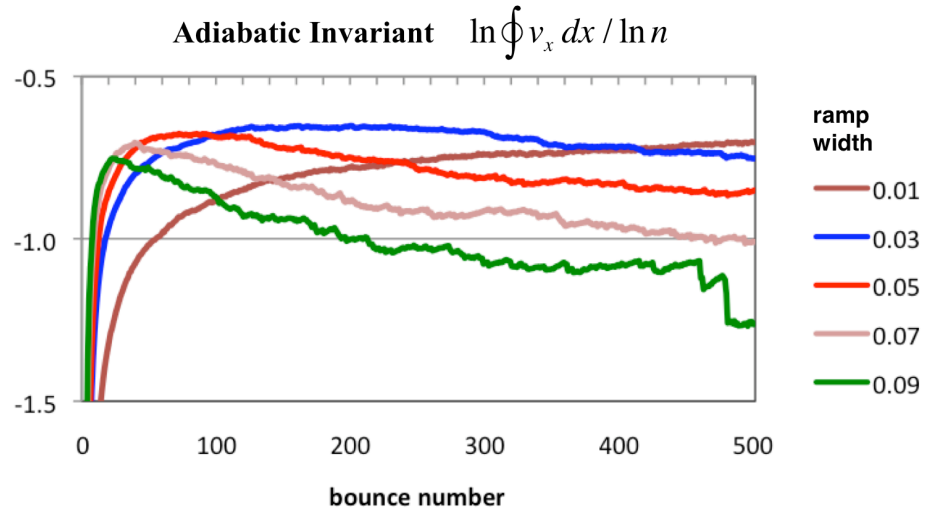


Figure A-6 Power Law Plot of Adiabatic Invariant

Violation begins at widths greater than $L_r = 0.03$ or approximately one electron inertial length

The kinetic energy gain also exhibits significant changes with increasing ramp width as seen in Figure A-7 and the accompanying log plot Figure A-8. The range of bounce values shown is well within the time limit to reach the maximum theoretical velocity ceiling.

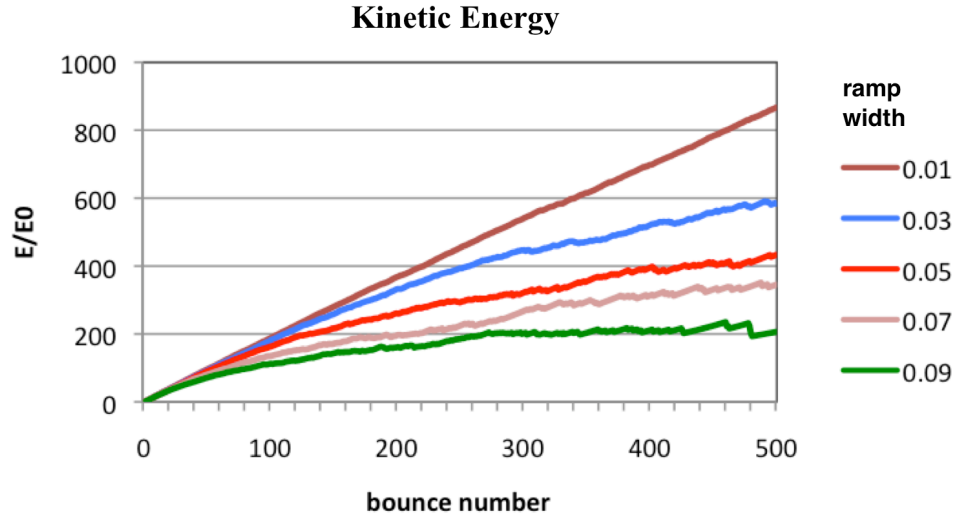


Figure A-7 Kinetic Energy growth

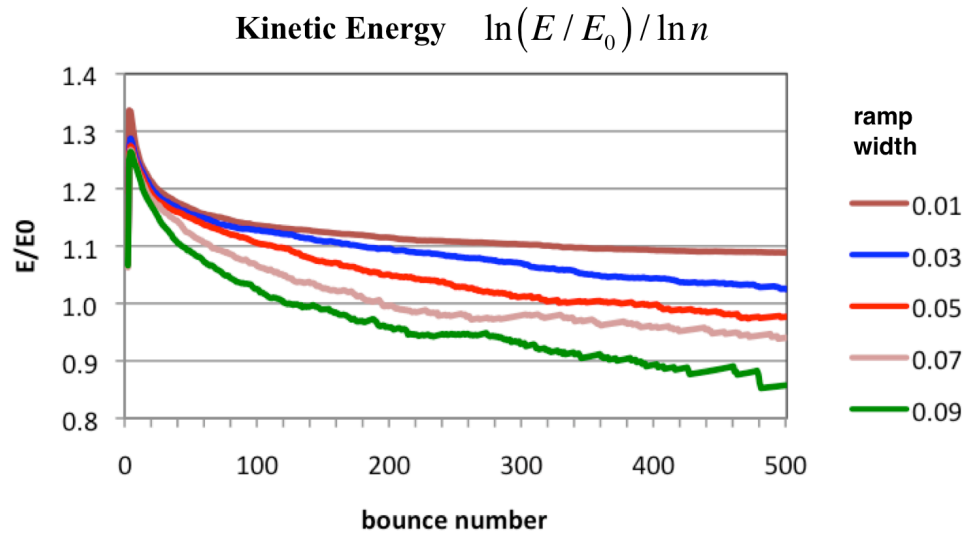


Figure A-8 Kinetic Energy Power Law Plot

As shown in Figure A-9, the period between reflections is only weakly sensitive to ramp width within the range displayed. However, the dependence increases noticeably as the width increases beyond $1/10$ ion inertial length. The erratic appearance of the $L_r = 0.09$ curve is due to the paucity of data for this width at high bounce numbers.

Recall that the ‘support’ region of incident pickup ions shrinks with increasing ramp width, so that it takes much longer for the simulations to fill in these points.

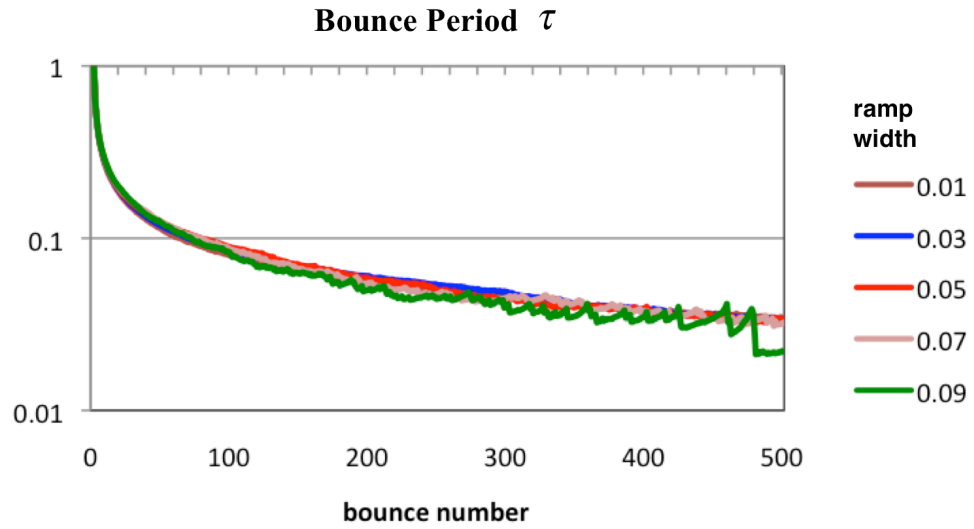


Figure A-9 Bounce Period

As shown in the Figure A-10, below this critical ramp width and after approximately 50 reflections, the bounce period may be expressed as a negative power law $\tau = n^{-p}$ with exponent $p \sim 0.55$.

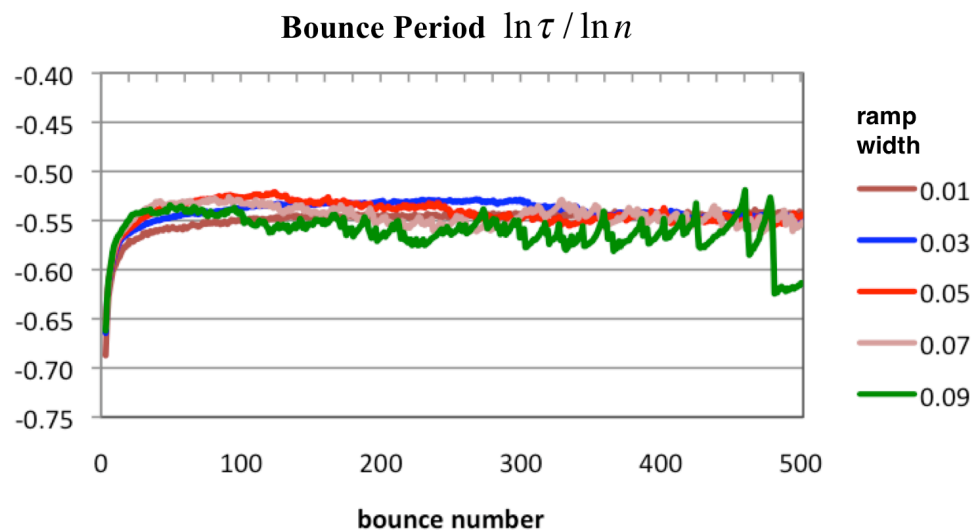


Figure A-10 Bounce Period Power Law Plot

Of particular interest, especially in light of the dependence of kinetic energy growth per bounce on ramp width, is the measured bounce amplitude shown in Figure A-11.

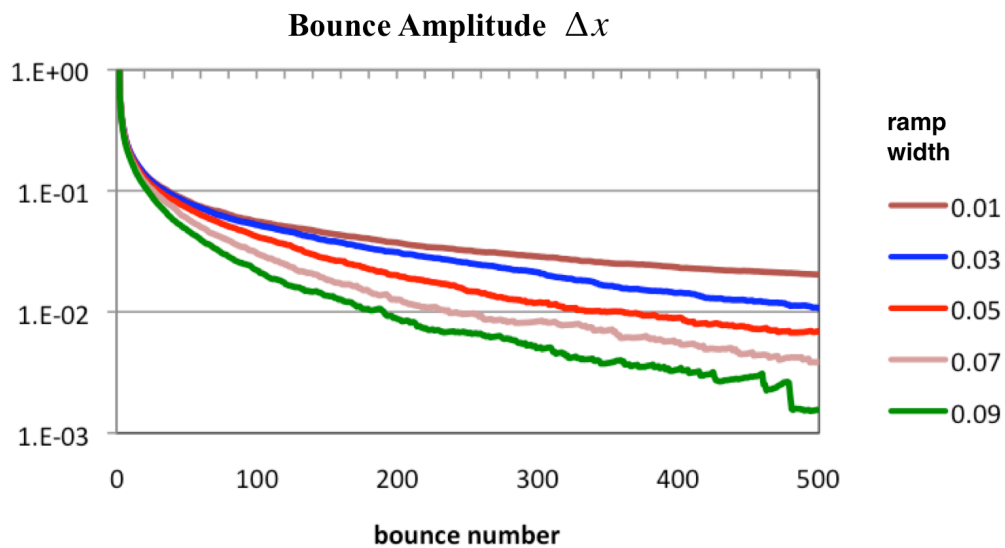


Figure A-11 Bounce Amplitude

As seen in Figure A-12, only the very thinnest ramp admits a power law dependence on the bounce number.

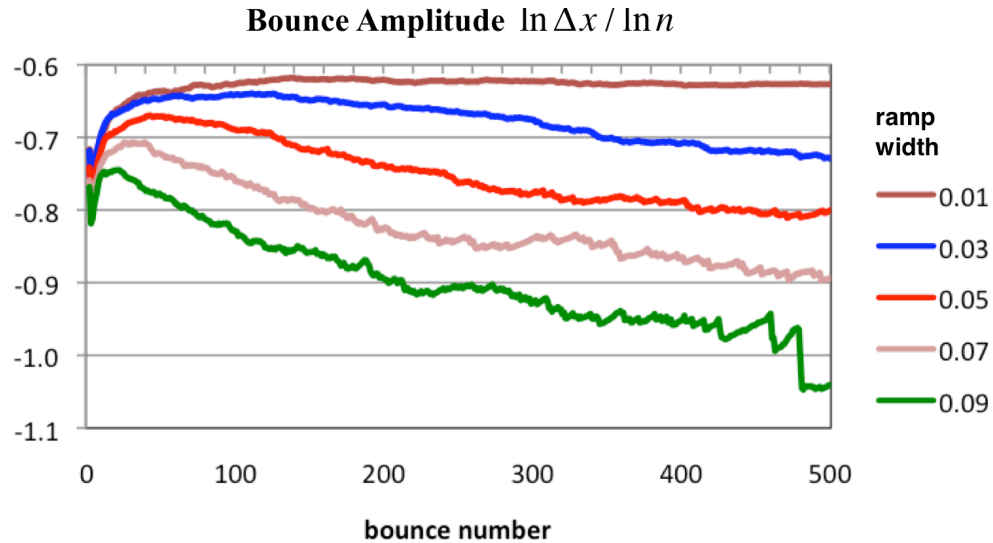


Figure A-12 Bounce Amplitude Power Law Plot

The range of ramp widths shown is extremely narrow compared with those of numerous shocks observed in the heliosphere. That the variables in question deviate so rapidly from theoretically predicted values indicates that this formalism may not necessarily describe the behavior of the shock surfing mechanism in the more general context of the manifold shocks actually found in nature.

Appendix B Plasma Waves

Alfven, Lower Hybrid and Magnetosonic waves in turn were imposed upon the background fields previously employed to drive test particle trajectories. A battery of simulations was performed as in the Baseline studies with parameters to model protons encountering a perfectly perpendicular Mach 5 shock with very thin ramp of width $L_r = 0.05\lambda_{pi}$, equivalent to about two electron inertial lengths. We varied the wavenumber and polarization appropriate to each wave type and explored a region of parameter space that we believed would be both physical and likely to produce significant results. The field amplitudes and wavenumbers ranged from zero to values exceeding those so far observed in situ, and in agreement with theoretically proposed physical limits.

Method

Several schemes for imposing waves upon the background fields were attempted to localize perturbations near the shock front. When upstream wave regions extend beyond a few pickup ion cyclotron radii, wave induced velocity diffusion expands the initial PI velocity shell during transit to the front, thus depleting the surfing candidate region in velocity space. This in turn lowers the surfing ‘efficiency’, which we avoided by limiting the size of the simulation box to include a small injection region upstream of the ramp, and a downstream region just large enough to allow upstream return of the population of transiting ions. Especially for the case of Lower Hybrid waves, this approach was motivated by the presupposition that such waves may themselves be

generated by accelerated surfing ions, but also had the merit of introducing fewer extraneous data due to wave-induced velocity diffusion outside the ramp.

Lower Hybrid Waves

Lower Hybrid waves (LH) are longitudinal electrostatic fluctuations that propagate nearly perpendicular to the ambient magnetic field, within an angle $\Delta\theta < \sqrt{m_e / m_i}$. Since we employ test particles, we are free to choose the wave parameters, with some discretion required. Since we require $\mathbf{k}_{LH} \perp B_0 \hat{\mathbf{z}}$, we may choose $\mathbf{k}_{LH} = k \hat{\mathbf{x}}$ or $\mathbf{k}_{LH} = k \hat{\mathbf{y}}$ which requires that $\mathbf{E}_{LH} = E(x) \hat{\mathbf{x}}$ or $\mathbf{E}_{LH} = E(y) \hat{\mathbf{y}}$ respectively. Since the LH frequency is roughly $\omega_{LH} = \sqrt{\Omega_{ci} \Omega_{ce}} \approx 43 \Omega_{ci}$, and thus much lower than the electron plasma frequency ω_{pe} , we impose static waves. Waves with vectors $\mathbf{k}_{LH} = k \hat{\mathbf{y}}$ travel along the shock front and, while $\mathbf{k}_{LH} = k \hat{\mathbf{x}}$ waves cannot propagate far upstream against the supersonic solar wind unless excited by back-streaming particles, we include these in the spirit of exploring the effects of such waves standing in the near-upstream region of the shock front.

We show below in Figure B-1 the effect of LH waves on the total fluxes of surfing ions, and in Figure B-2 the surfing ‘efficiency’ as defined by the integrated flux of ions promoted above three times the solar wind speed, and in Figure B-3 the surfing velocity ceiling. In each figure, the left panel represents the $\mathbf{k}_{LH} = k \hat{\mathbf{x}}$ waves and the right panel, those propagating the $\mathbf{k}_{LH} = k \hat{\mathbf{y}}$ direction.

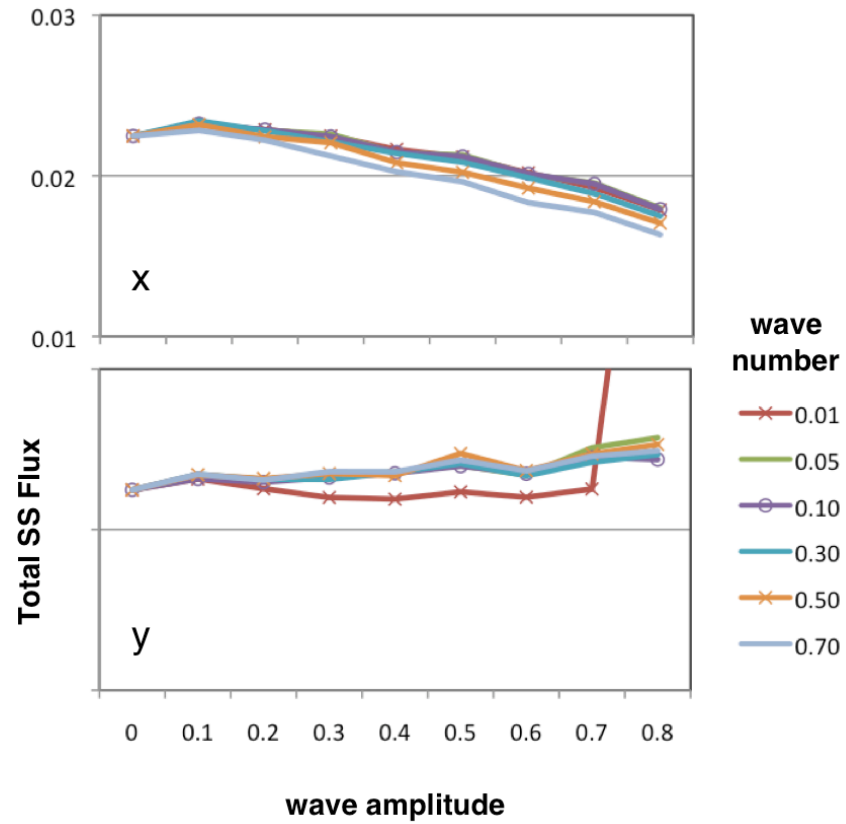


Figure B-1 Effect of LH waves on total SS fluxes

Waves were imposed on protons at Mach 5 shock in the range $0.01 \leq k \leq 0.7$ propagating in both the shock-normal **x**-direction (top panel) and along the front in the **y**-direction (bottom panel) with no appreciable effect on the total SS flux.

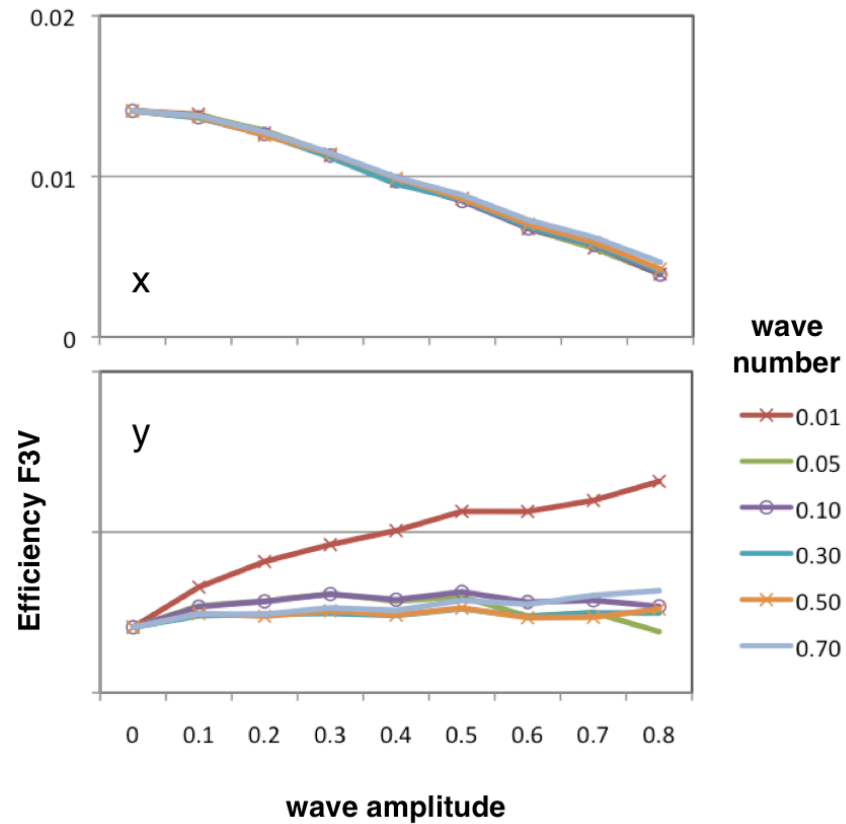


Figure B-2 Surfing efficiency 'F3V' with LH waves

The same parameter range was explored with negligible suppression on the SS flux above the SD ceiling for **x**-waves, and slight enhancement for those propagating in the **y**-direction at the longest wavelength $k = 0.01$.

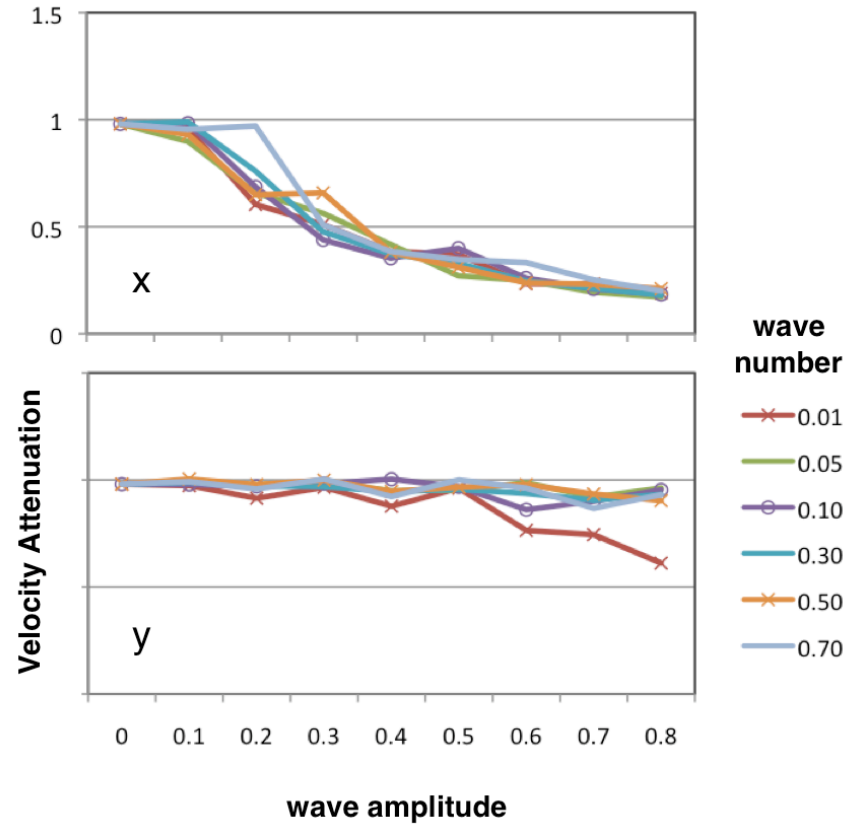


Figure B-3 Maximum Surfing Velocity with LH waves

No significant velocity suppression occurred with y-polarized waves imposed (right panel) while x-polarized waves (left panel) showed significant attenuation at all wavelengths. However, the maximum applied amplitude was several times that of the ambient magnetic field.

No significant effects on surfing ions were observed due to LH Waves propagating along the front, i.e., $\mathbf{k}_{LH} = k \hat{\mathbf{y}}$ even at the highest amplitudes and longest wavelengths applied. Suppression the total surfing fluxes was minor, and no significant effect apparent on either the surfing 'efficiency' or velocity maxima regardless of amplitude or wavelength. The only exception occurred at the very longest wavelength $k = 0.8$, where a small 'spike' appeared in the total flux, possibly attributable to a

resonance since the flux returned to the baseline at higher amplitudes. The efficiency also increased with amplitude for that wavelength, but we argue that the effect is entirely inconsequential when compared with other parametric dependence, especially ramp width.

On the other hand, the $\mathbf{k}_{LH} = k \hat{\mathbf{x}}$ waves attenuated both the surfing ‘efficiency’ and velocity ceiling by as much as 75% for the very highest amplitude waves imposed. However, we again assert from our baseline studies that small deviations in the magnetic field angle from the perpendicular or increases in ramp width far outweigh the observed effects of LH waves. Indeed, these effects entirely disappear when the ramp widens beyond $L_r = 0.1\lambda_{pi}$, still well below those of ramps actually observed in situ as we demonstrate in Chapter 6.

Shapiro et al (2001) argue that LH waves may either be pre-existing or excited by accelerated particles, and that such turbulence may in turn scatter “accelerated particles effectively, removing them from the surfing acceleration.” They add that the wave energy density “can be as large as 3–5% of the solar wind kinetic energy density”, consistent with typical values measured at the Earth’s bow shock by the Prognoz 8 mission in 1983. The kinetic energy described is simply the ram energy $m_p n_p u_{sw}^2 / 2$, which, in our simulation units is simply $M_A^2 / 2$. Assuming the 5% of ram energy as proposed by Shapiro et al gives a maximum field amplitude (square root of the energy density) roughly $0.16 M_A$ while the maximum field amplitude we imposed was normalized to that of the convective electric field, $E_{\parallel} = u_1 B_1 = M_A$. In other words, our field scale (normalized to unity) was six times greater than actually observed, corresponding to energy density almost 40 times larger. Hence, we claim that our Lower Hybrid study is

more than exhaustive, and suggest that realistic effects on shock surfing should lie within the amplitude range $0 \leq a \leq 0.16$ on the plots we have shown.

Shear Alfven Waves

Alfven waves are transverse magnetic fluctuations that propagate along the ambient magnetic field lines. In that sense, they are not unlike transverse waves along a stretched string. Taking the background magnetic field $\mathbf{B} = B_0 \hat{\mathbf{z}}$, two wavevectors (or combinations thereof) are possible, $\delta \mathbf{B} = \delta B(z) \hat{\mathbf{y}}$ and $\delta \mathbf{B} = \delta B(z) \hat{\mathbf{x}}$ but, in our one-dimensional simulation scheme, the latter was not functionally appropriate. We show below in Figure B-4 plots of the total surfing flux in the left panel and the surfing ‘efficiency’ on the right. Although the total flux is somewhat enhanced at higher amplitudes, the flux promoted above the SD velocity ceiling is only marginally so, and then only briefly before falling off as the wave amplitude is further increased to the highest values imposed. The surfing velocity maxima shown in Figure B-5 exhibit only the slightest variation under the imposition of these waves at all lengths and amplitudes.

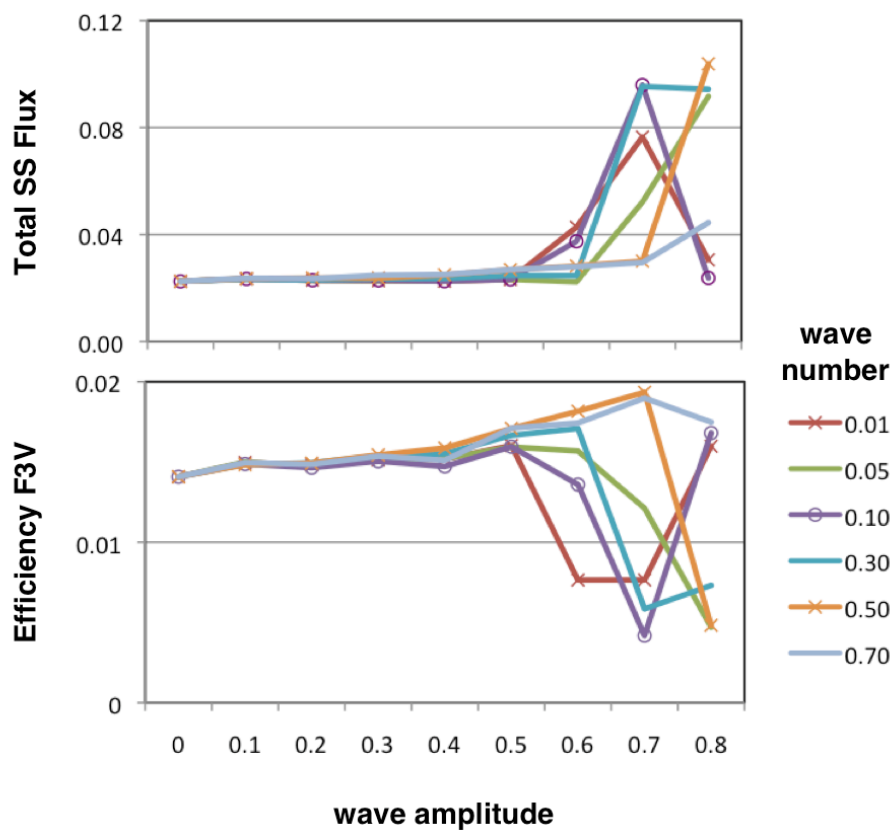


Figure B-4 Fluxes with Shear Alfvén Waves

Some enhancement in the SS flux (top panel) is seen at large field amplitudes. The SS efficiency (bottom panel) increases slightly but drops dramatically at higher amplitudes, suggesting mechanism disruption.

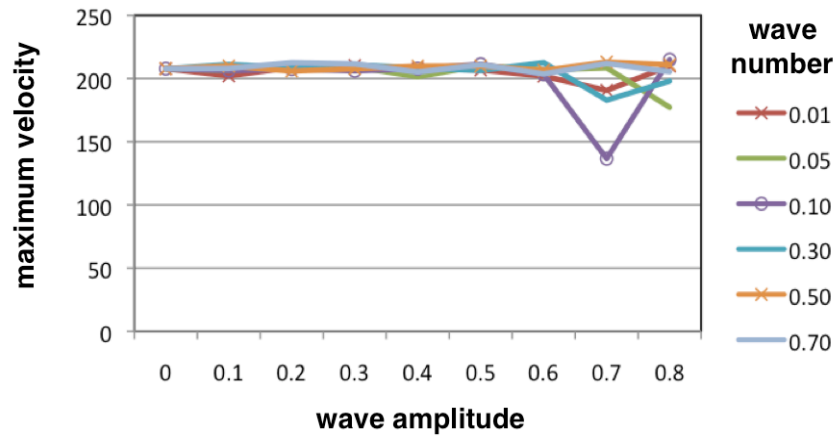


Figure B-5 Surfing Velocity Maxima

The effect at all amplitudes and wavelengths is entirely negligible.

Magnetosonic Waves

Magnetosonic (MS) waves are essentially a combination of thermal and magnetic pressure waves that propagate perpendicular to the background magnetic field. Hence the field perturbations must have the same orientation as the ambient field, but we are free to choose the wavevectors provided that $\mathbf{k}_{MS} \perp B_0 \hat{\mathbf{z}}$, and we chose these in the plane defined by the coordinate axes $\hat{\mathbf{x}}$ and $\hat{\mathbf{y}}$. The effects of these waves on the total surfing fluxes are shown below in Figure B-6 for the two polarizations.

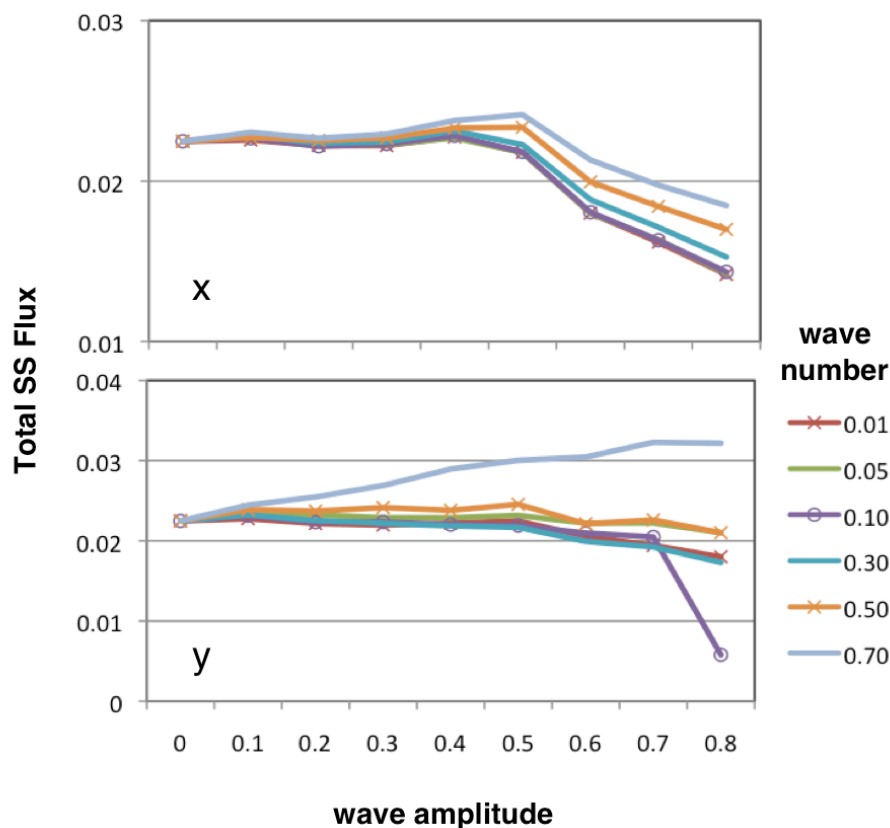


Figure B-6 Total Surfing Fluxes with Magnetosonic Waves
Wave vectors are in the **x**-direction (left) and **y**-direction (right). Some variation in fluxes is seen, but the effect is trivial when compared to variation in ramp width.

The effects of MS waves on the fluxes of surfing ions promoted above the SD velocity ceiling are shown below in Figure B-7. Here we clearly see that the distribution function tails of SS ions are profoundly suppressed by **x**-waves (shown in the left panel), and that the effect is dependent upon amplitude, but nearly insensitive to wavelength. Attenuation by **y**-waves (right panel) is wavelength-dependent, but with mixed effect depending upon the lengths of various waves imposed.

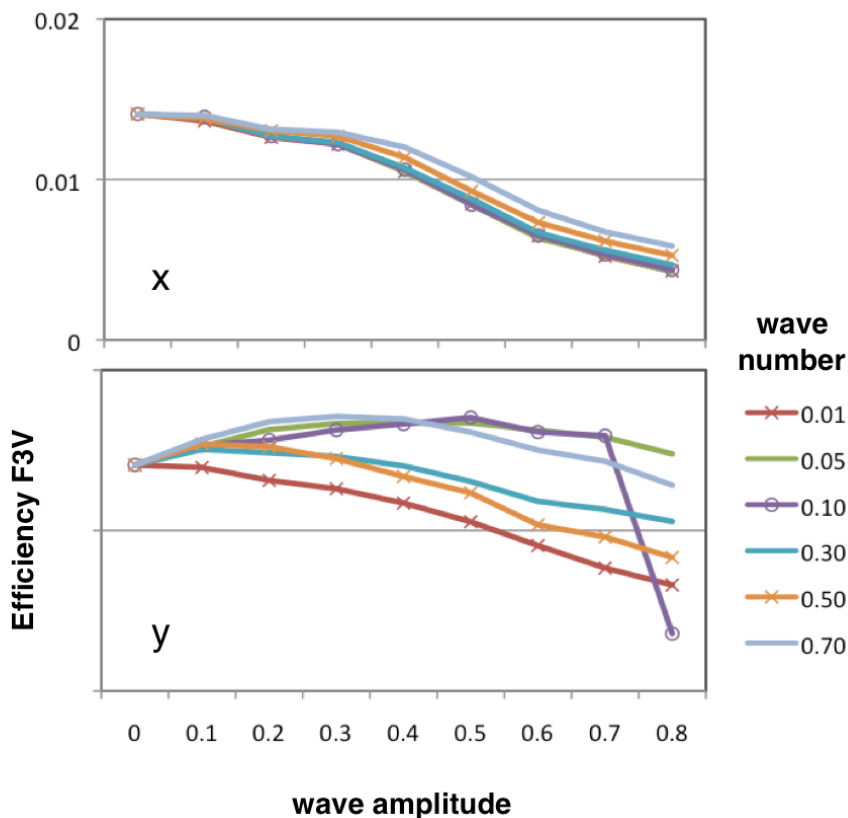


Figure B-7 Effect of MS waves on Surfing Efficiency

Waves traveling in the **x**-direction (left) suppress the accelerated SS flux by more than 60% at the highest amplitude, independent of wavelength.

Waves in the **y**-direction only accomplish this at certain wavelengths.

Finally, we show in Figure B-8 the effect of MS waves on surfing ion velocity maxima. In the case of both wave polarizations imposed, the velocity suppression is approximately 50% at amplitudes approaching the physical limit normalized to unity. Although the velocity attenuation appears wavelength-insensitive, the velocity decreases monotonically with increasing wave amplitude.

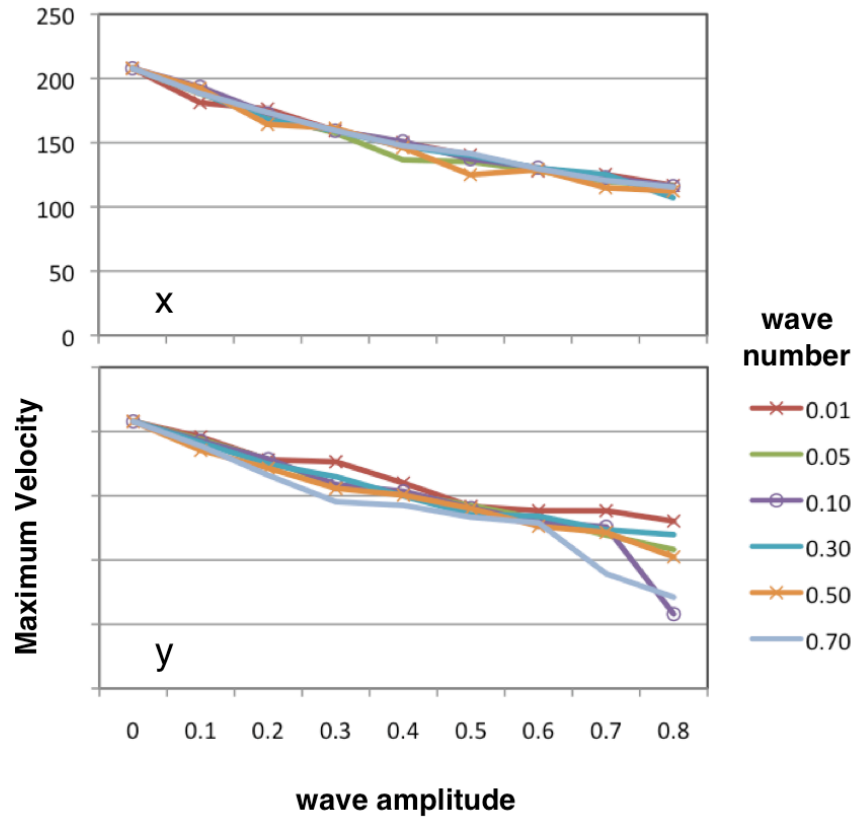


Figure B-8 Effect of MS waves on Velocity Maxima

Velocity ceiling is suppressed by a factor of two, independent of wavelength. The same effect ensues by increasing the ramp width by the same factor from $L_r = 0.05$ to $L_r = 0.10$, well below typically observed values on the order of one ion inertial length.

Our results may be generally summarized thus: extremely large field amplitude is required for any waves to affect the distribution functions of pickup ions accelerated by shock surfing.

Lower Hybrid waves do so only when their propagation direction is normal to the front, and effects from Shear Alfvén waves are only significant at the highest physically realistic energy. However, Magnetosonic waves do attenuate both the fluxes of ions accelerated by shock surfing as well as the velocity maxima attainable by this

mechanism. How and whether waves of this type may actually be produced in strength at shock fronts is essentially moot, since attenuation of the maximum surfing velocity and tail populations are observable only for the very steepest shock ramps. To put the effect of these perturbations in perspective, we note by comparison that the strongest suppression of SS within the parameter range tested (roughly 50%) is equivalent to increasing the ramp width from $L_r = 0.05$ to $L_{ramp} = 0.10$ which is still an order of magnitude smaller than the ion inertial length, commonly observed *in situ*. Hence, even at the stoutest perpendicular shocks where this variability would be considered negligible, the effects of ramp widening on the SS acceleration mechanisms would dramatically outweigh those of the strongest plasma waves.

We now exhibit a complete set of plots from simulations that showed some measureable effect. We show three sets for each type of wave imposed, including orientation or polarization, with individual panels for each wavenumber.

- Total mechanism fluxes where ‘FCC’ is the ‘crossing’ species, ‘FBB’ the ‘single-bounce’ and ‘FSS’ the shock surfing ions.
- Surfing ‘efficiency’ given by three measures, ‘S>C’ represents the flux of surfing ions accelerated above the ‘crossing’ species velocity ceiling, ‘S>B’ similarly for the ‘single-bounce’ species, and ‘F3V’ the flux of surfing ions promoted above three times the upstream wind speed.
- Surfing ion velocity maxima ‘VSS’ along with the same measure for the other two mechanisms, ‘crossing’ (VCC) and ‘single-bounce’ (VBB).

All simulations were run for protons at a perpendicular Mach 5 shock. In all plots, the fluxes are given as a fraction of the total incident pickup ion population. Wave amplitudes

are normalized to the appropriate field strength for the type of wave imposed – the background magnetic field $B_0 \rightarrow 1$ for Alfvén and Magnetosonic waves, and the convective electric field $\frac{u_{SW}}{c} B_0 \rightarrow M_A$ for Lower Hybrid. The wavenumbers are normalized to the ion inertial length via $k_{pi} = 2\pi / \lambda_{pi}$ and ion velocities to the maximum shock surfing velocity attainable (approximately $v_{\max} = 212 v_A$) at the ramp width $L_r = 0.05 \lambda_{pi}$ that we employed for all the simulations in this set. The coordinates for all panels in each figure are identical, with the exception of Figure B-16 where the top left panel representing the smallest wavenumber $k = 0.01$ has maximum value 0.03 for the flux fraction rather than 0.02.

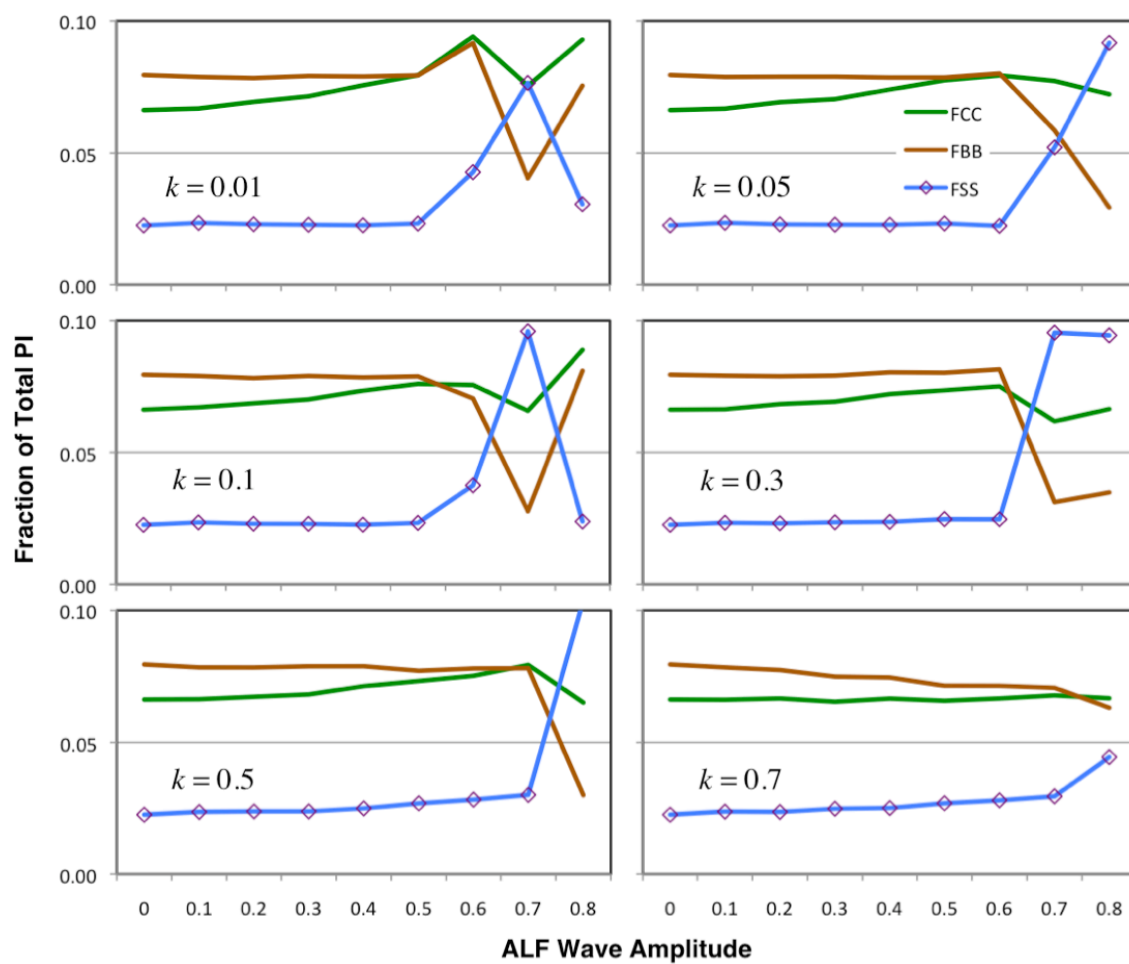


Figure B-9 Alfvén Wave effect on total fluxes

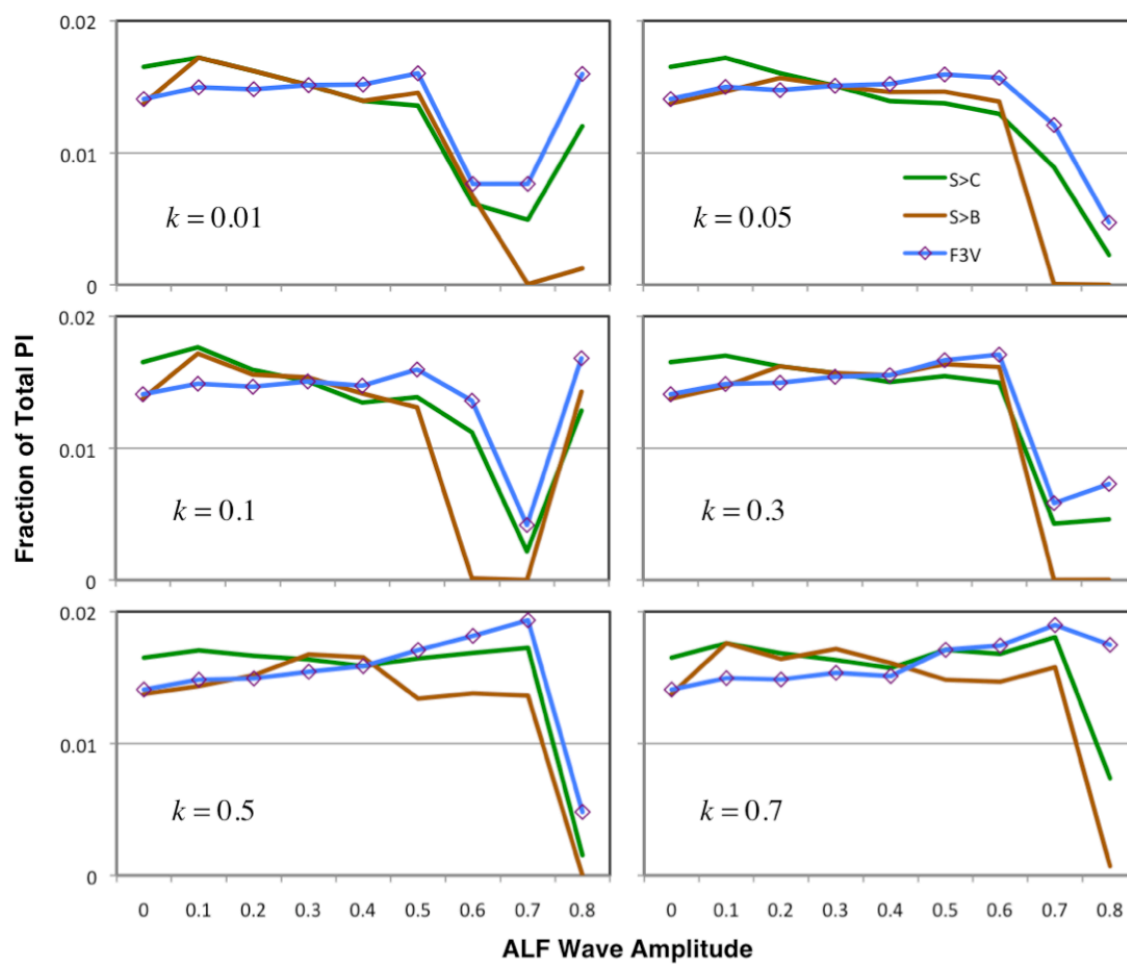


Figure B-10 Alfvén Wave effect on surfing efficiency

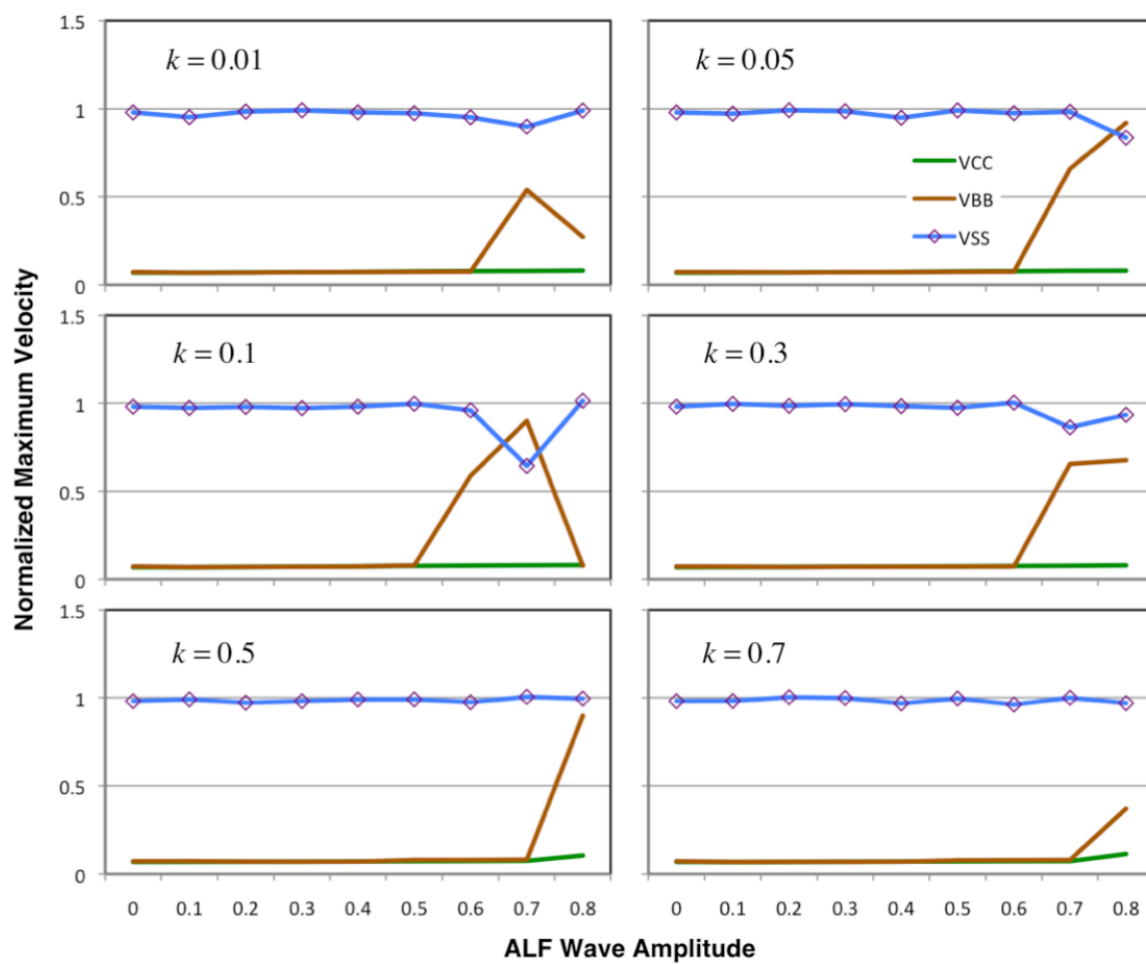


Figure B-11 Alfvén Wave effect on Velocity Maxima

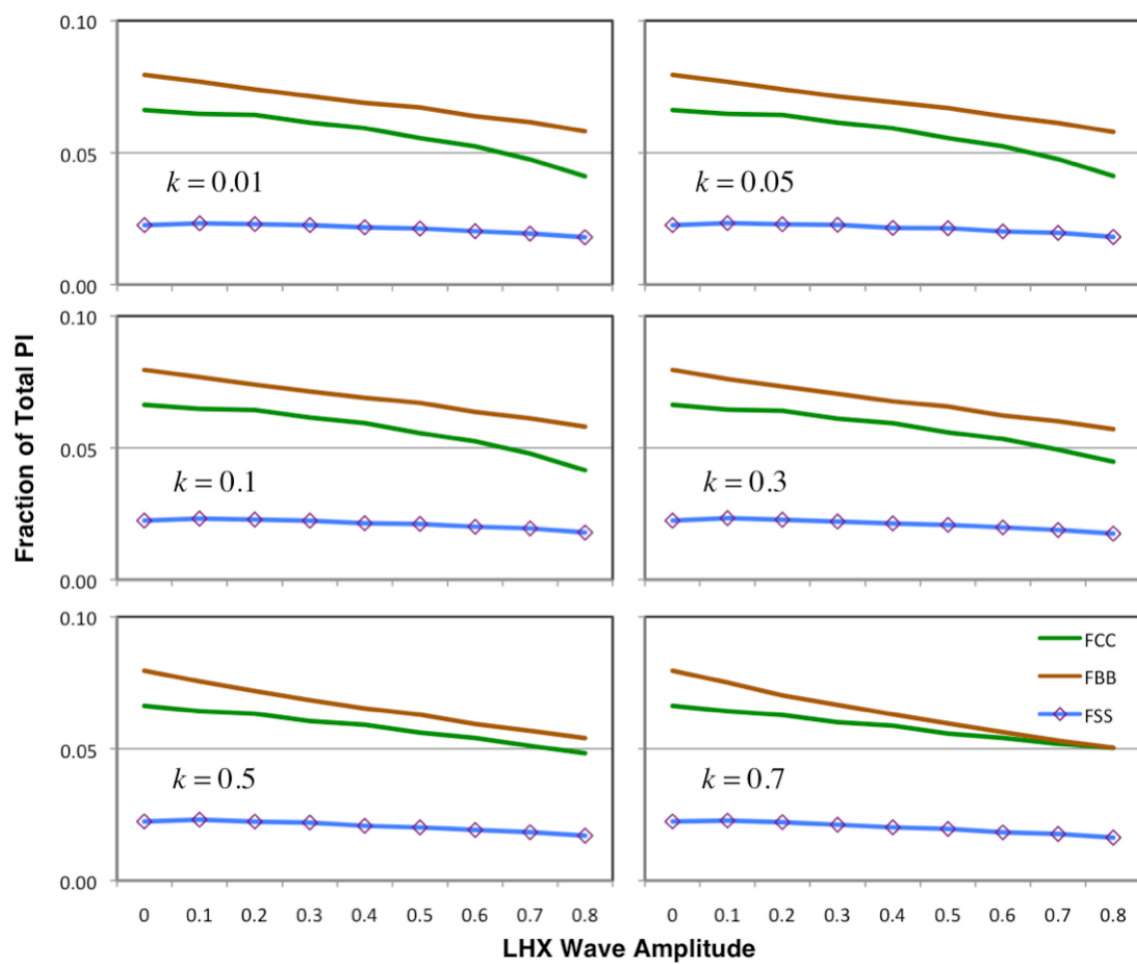


Figure B-12 Lower Hybrid x-Wave effect on Fluxes

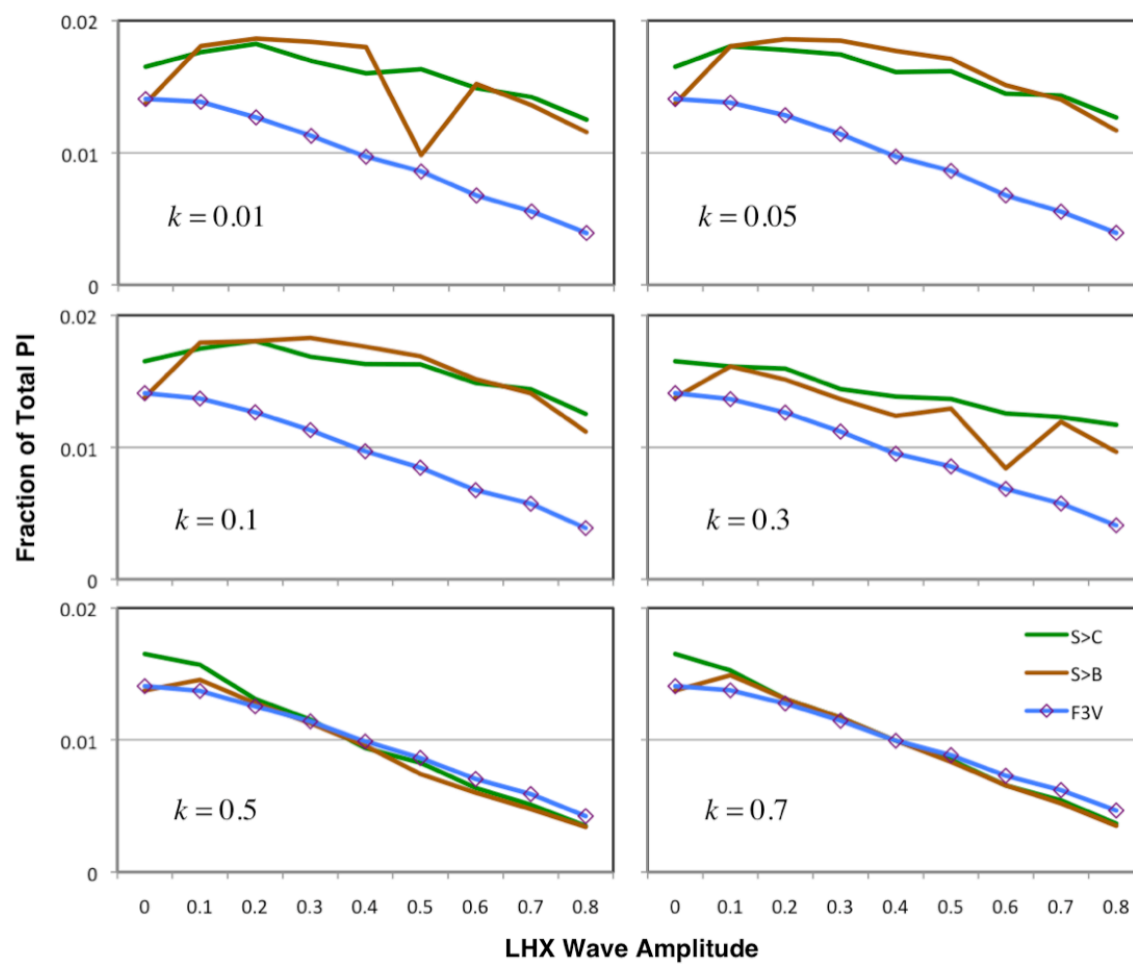


Figure B-13 Lower Hybrid x-Wave effect on Efficiency

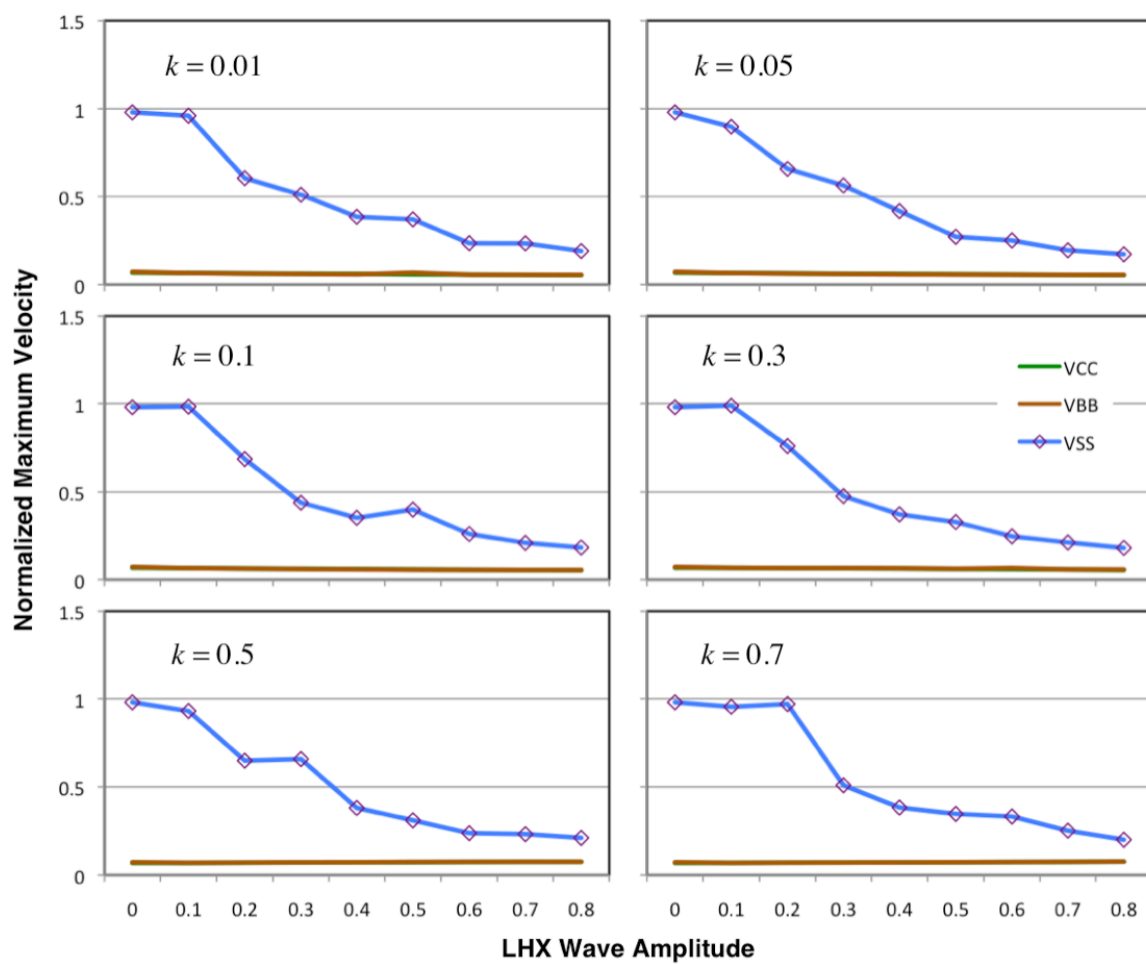


Figure B-14 Lower Hybrid x-Wave effect on Velocity Maxima

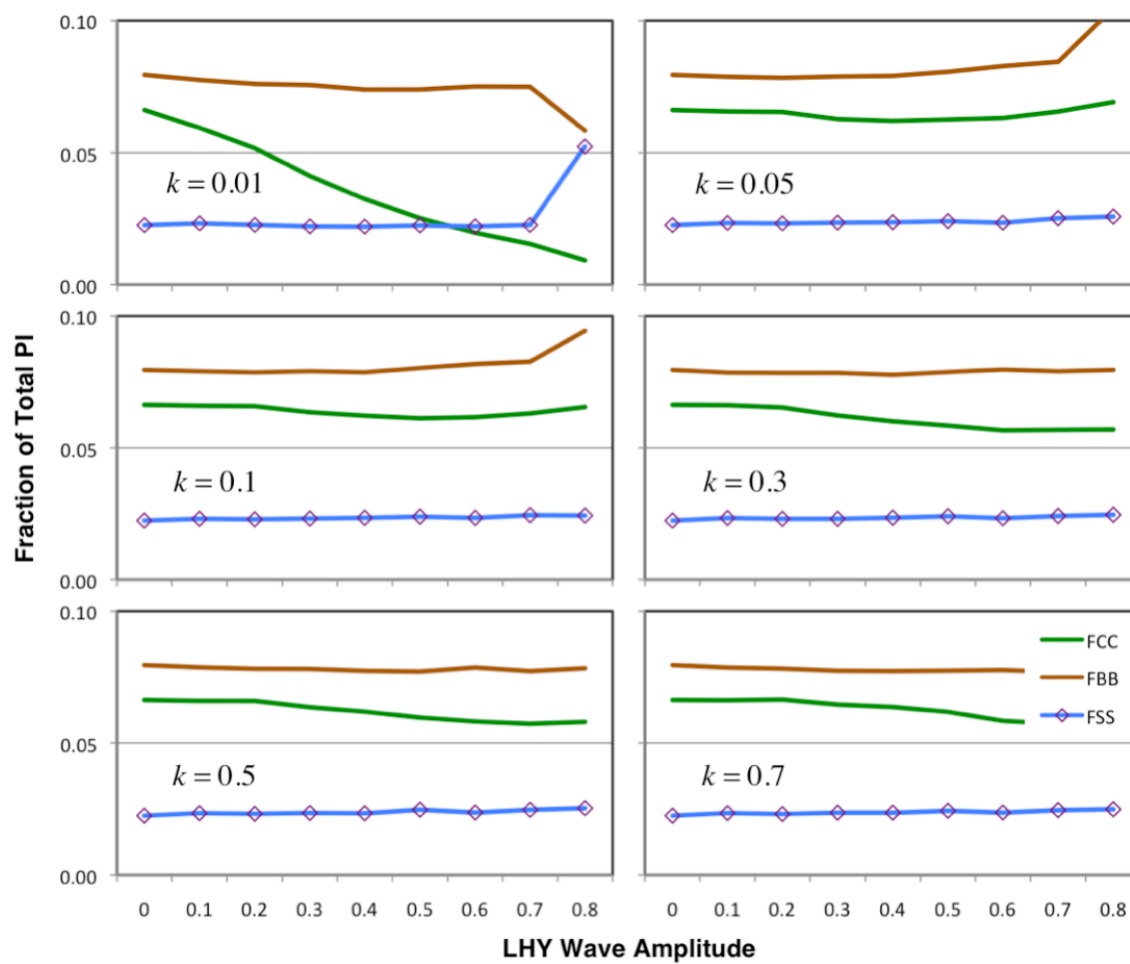


Figure B-15 Lower Hybrid y-Wave effect on Fluxes

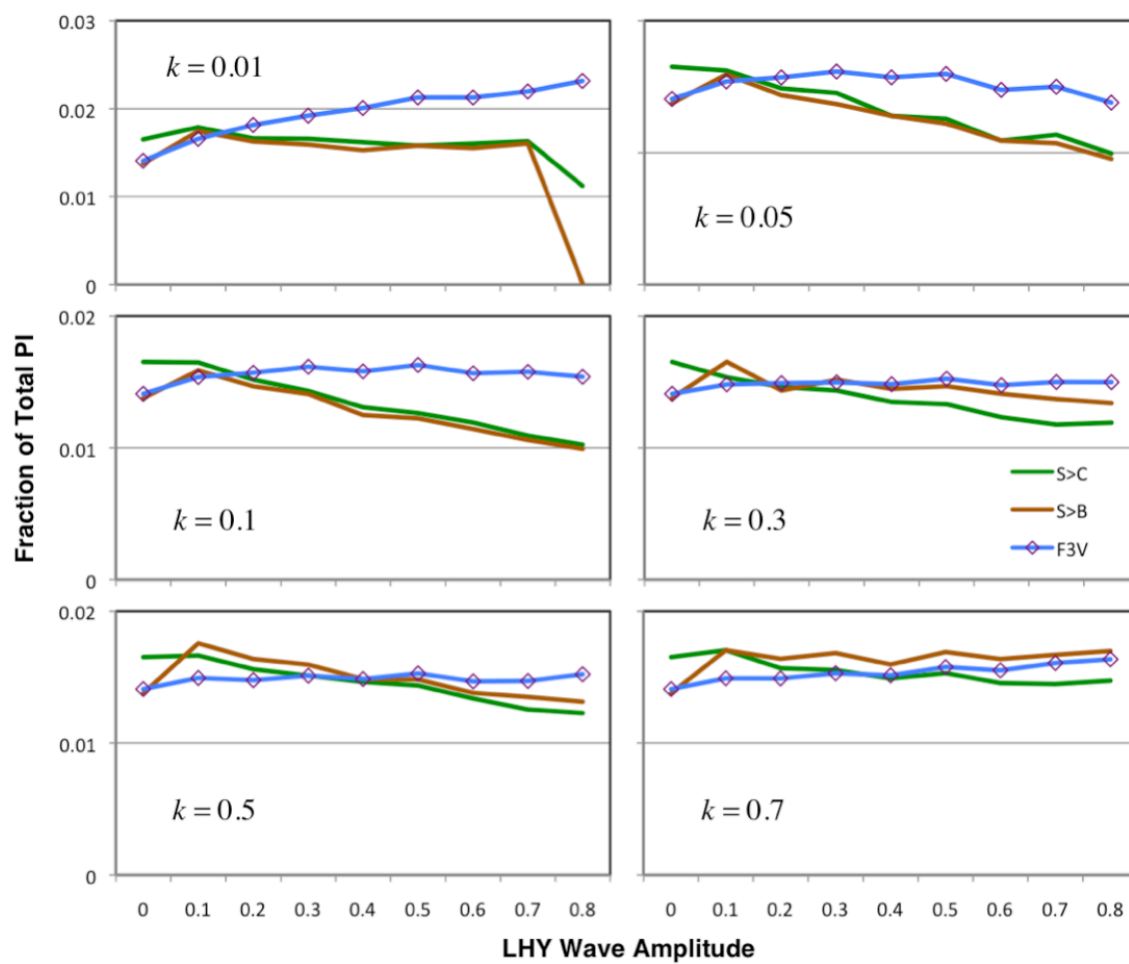


Figure B-16 Lower Hybrid y-Wave effect on Efficiency

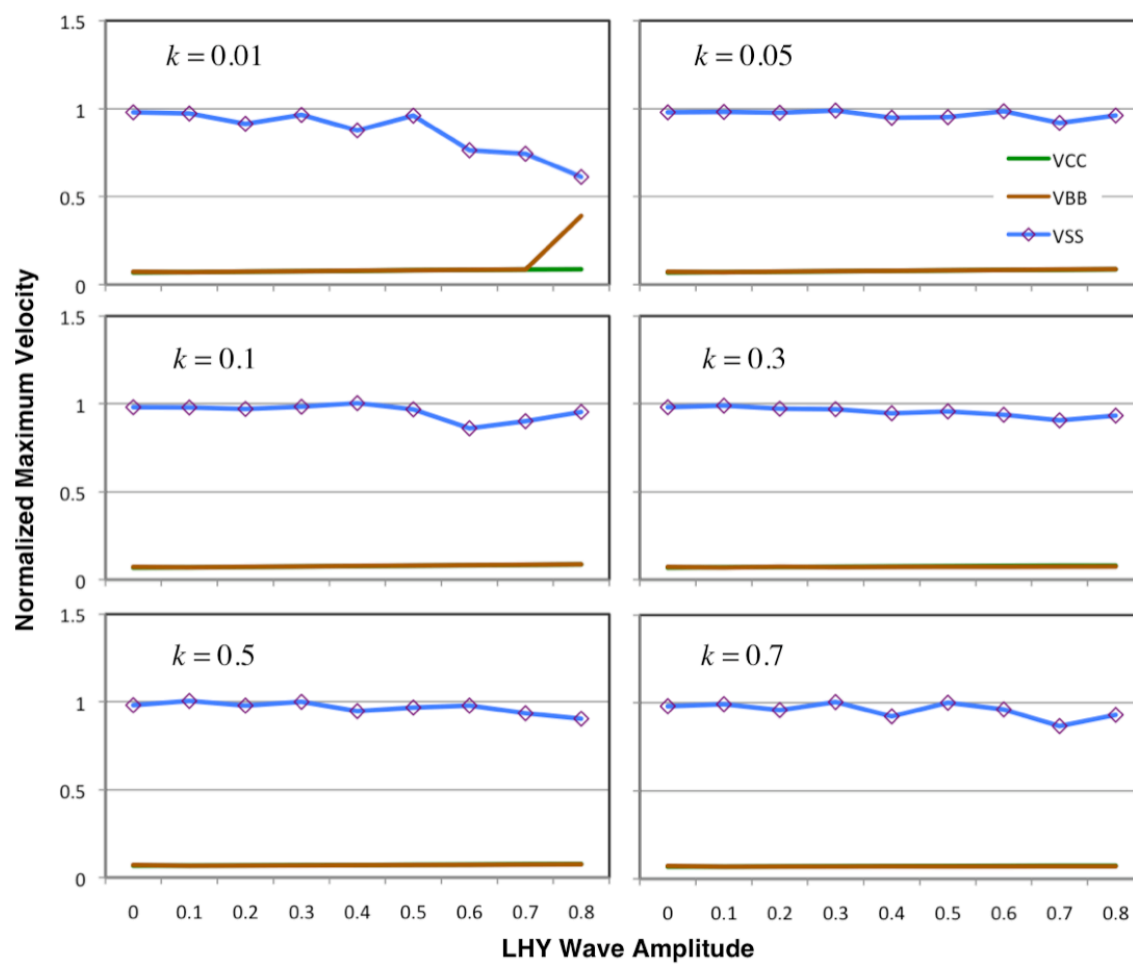


Figure B-17 Lower Hybrid y -Wave effect on Velocity Maxima

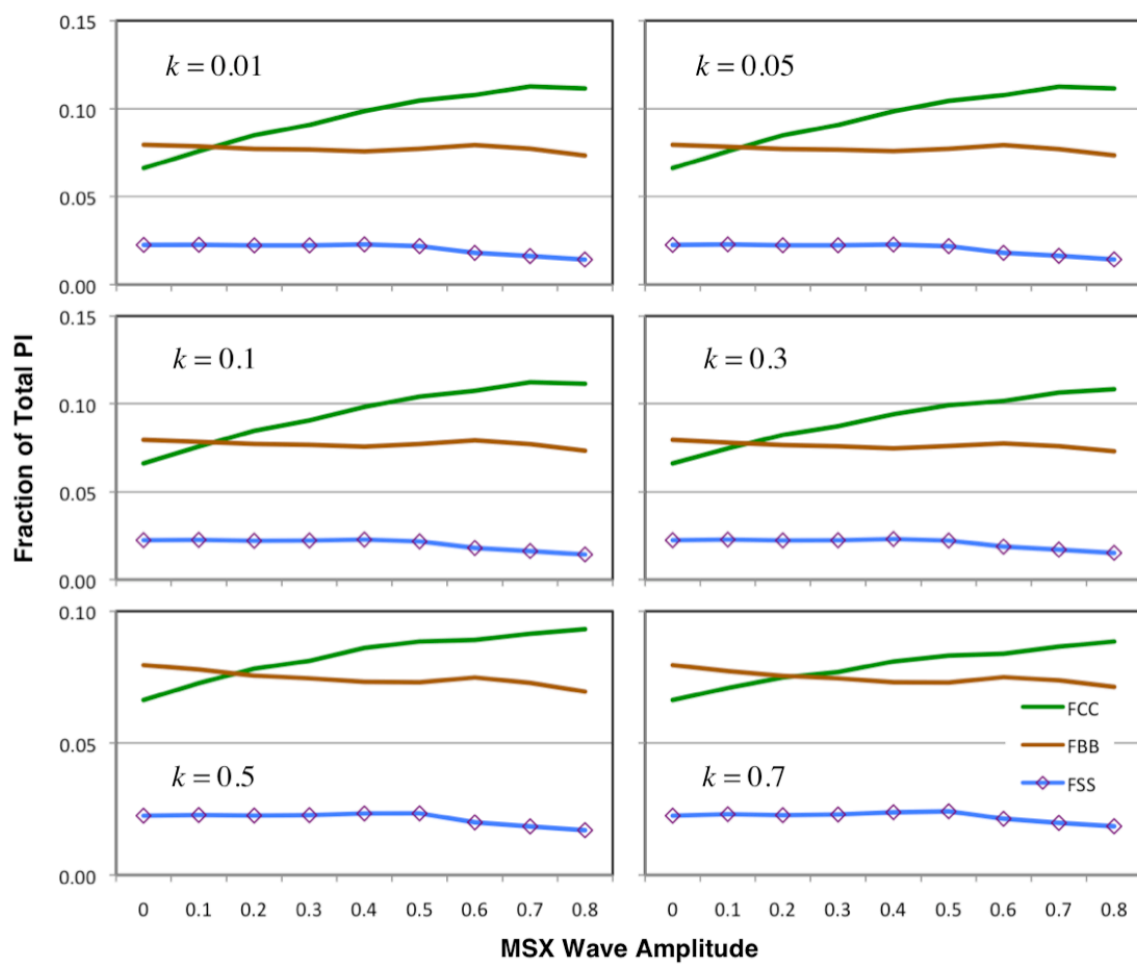


Figure B-18 Magnetosonic x-Wave effect on Fluxes

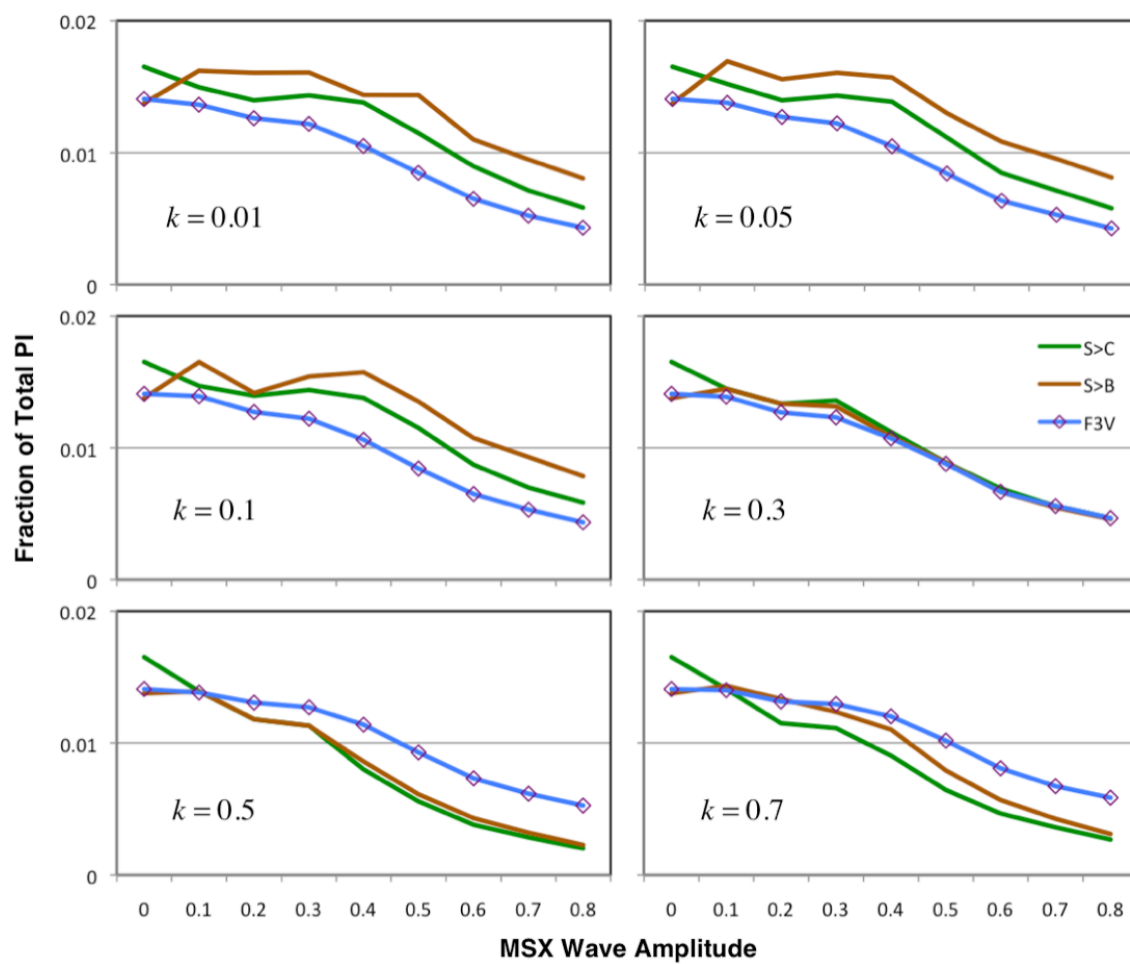


Figure B-19 Magnetosonic α -Wave effect on Efficiency

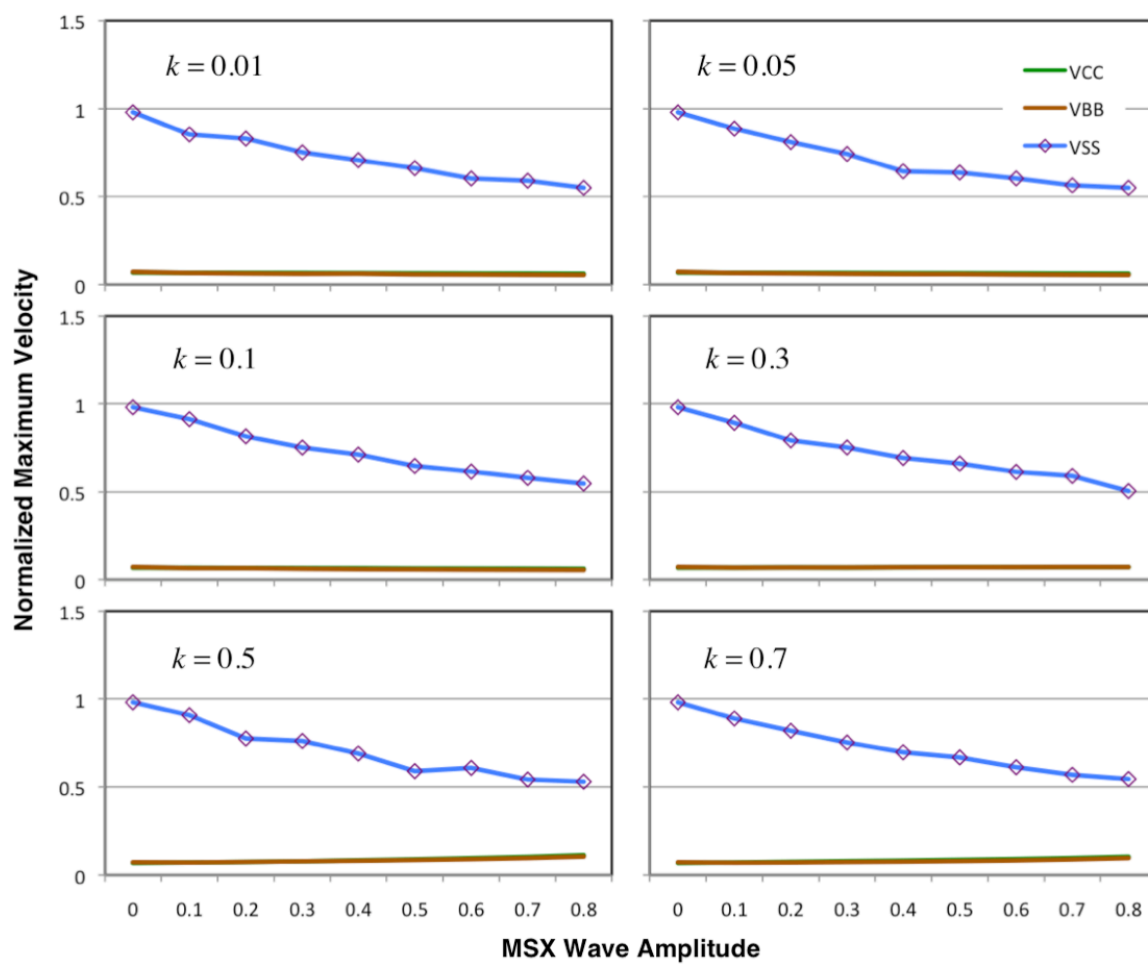
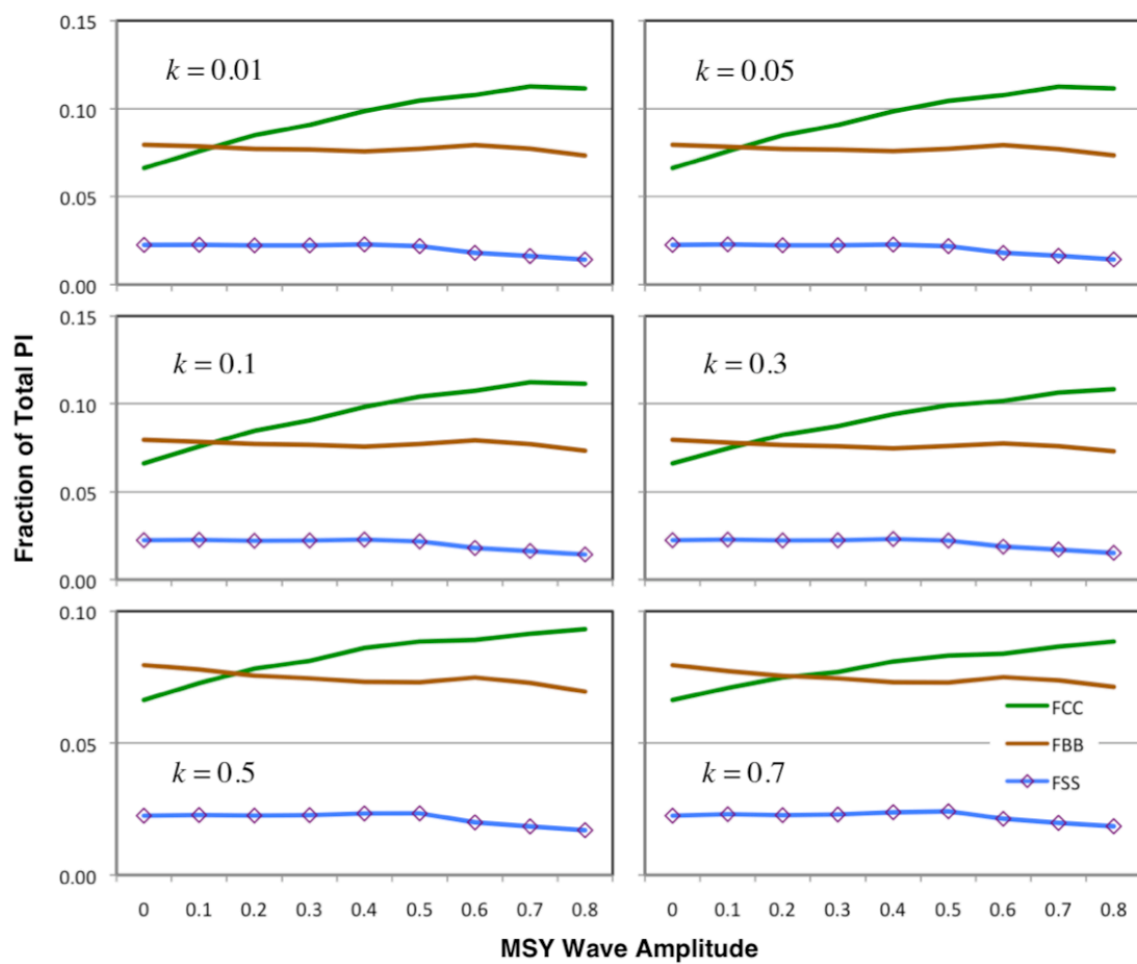
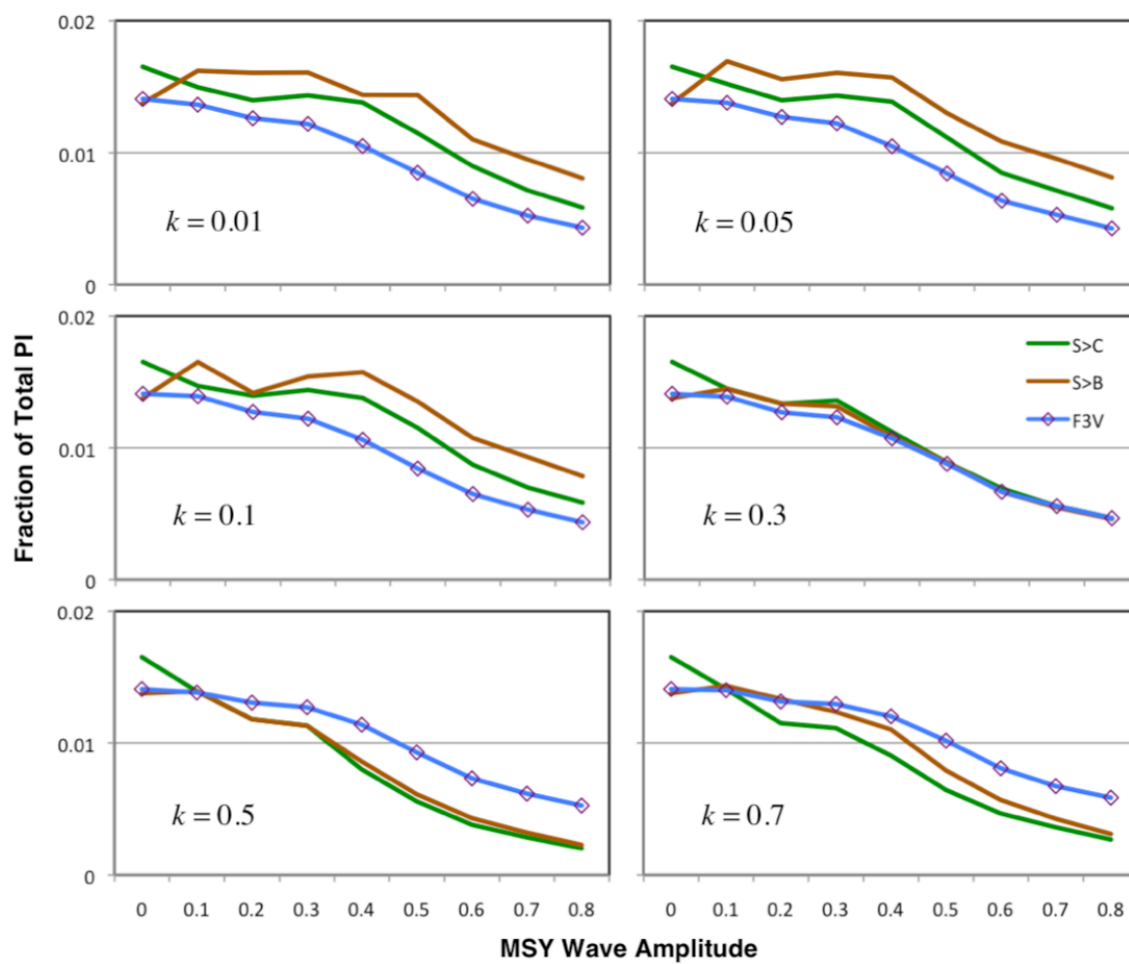


Figure B-20 Magnetosonic $\mathbf{x}$ -Wave effect on Velocity Maxima

Figure B-21 Magnetosonic y -Wave effect on Fluxes

Figure B-22 Magnetosonic y -Wave effect on Efficiency

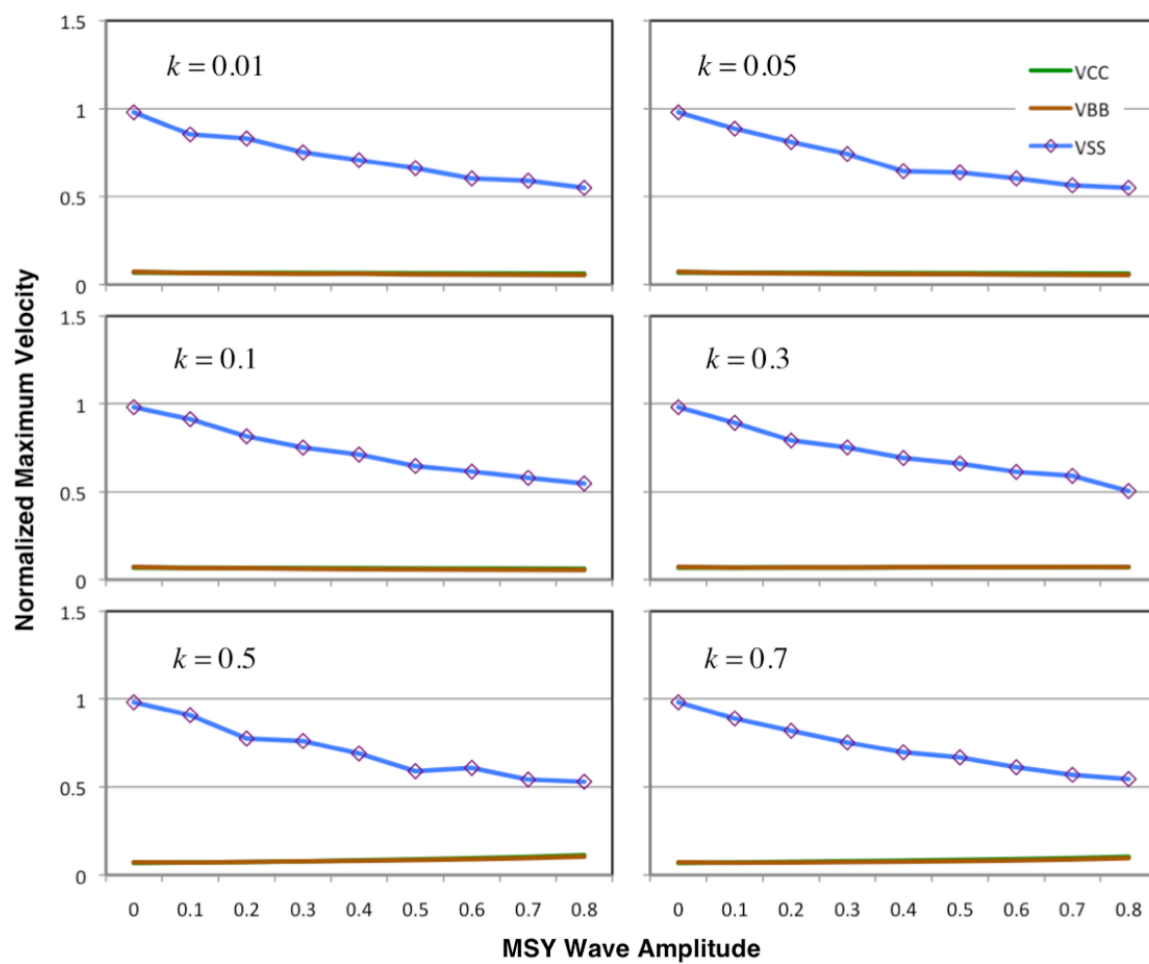


Figure B-23 Magnetosonic y -Wave effect on Velocity Maxima

Appendix C Corrugation Study Supplement

In this supplement we propose two physical models to explain the disruption of the Shock Surfing (SS) mechanism at corrugated fronts as evident from the trajectory plots shown in the text. This appendix is divided into three sections. We first present physical models, mathematical formulations, and derive functions for the surfing velocity ceiling, which is suppressed by rippled fronts below that due to finite ramp widths alone. For each model, we briefly exhibit data that graphically demonstrate the accuracy or failure of each mechanism and follow with supportive analytical discussion. In the second section, we justify in more detail some approximations made during the study, as well as the range of parameters explored, including the corrugation wavenumber and shock Mach number. We then conclude with a comprehensive set of supporting graphical data. For clarity and comparison of the graphical data, the force-balance function we derive first is subsequently labeled ‘Fcn1’ and the second function based upon a curvature-model is referred to as ‘Fcn2’.

Force Balance Solution – Function 1

The first model is based on simple static force balance. The ramp is assumed to be sufficiently thin that no account is made of normal trajectory oscillations inside. Further, velocity mixing outside the ramp is likewise ignored since the very fastest ions must closely follow the front inside the ramp. The accelerating force from the convective electric field is always parallel (locally tangent) to the surface. Forces on the ion normal to the surface are due to the perpendicular electric field, which is everywhere negative

$F_{E\perp} = -qE_{\perp}$ and the Lorentz force, which is always positive $F_B = q\frac{v}{c}B$. Finally, we have

a centripetal force required for the ion trajectory to follow the front. The direction and strength $F_C = \kappa mv^2$ depend on the local curvature where $\kappa = \frac{1}{R}$ is the geometric curvature and R is the radius of curvature. The maximum curvature modulus for sinusoidal perturbations $|\kappa| = ak^2$ occurs at crests and troughs and is derived in the following section. At a crest, the front is convex, as shown in Figure C-1.

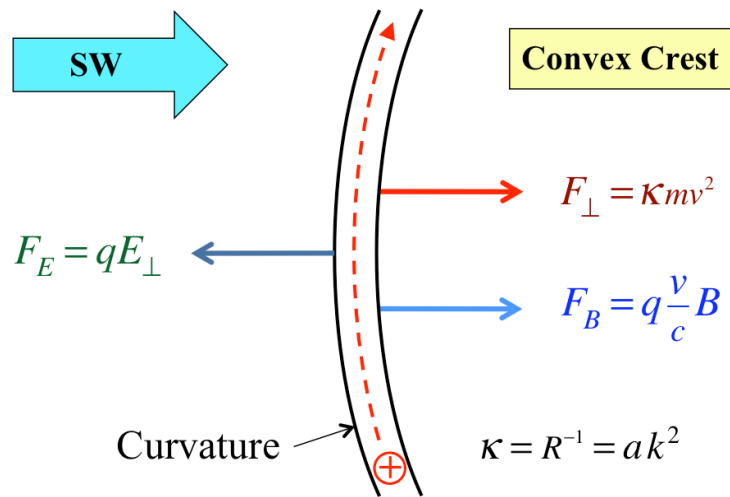


Figure C-1 Force Balance at a Convex Crest

The center of curvature is to the right, Solar Wind flows through the ramp to the right and the forces shown apply inside the ramp (curved lines)

For an ion to follow the convex curvature, the sum of these normal forces must balance

$ak^2mv^2 + q\frac{v}{c}B - qE_{\perp} = 0$. In the limit of vanishing perturbation strength, $v_{\max} \rightarrow \frac{cE_{\perp}}{B}$ so

we define a new variable $\tilde{v} \equiv \frac{B}{cE_{\perp}}v$ which instead asymptotes to unity. The equation in

terms of this variable becomes $\tilde{v}^2 + 2\Gamma\tilde{v} - 2\Gamma = 0$ where $\Gamma \equiv \frac{qB^2}{2c^2 E_{\perp} m a k^2}$. The physical

solution is $\tilde{v}_{>} = \Gamma \left[\sqrt{1 + \frac{2}{\Gamma}} - 1 \right]$ since, when the perturbation vanishes $ak^2 \rightarrow 0$ and

$\Gamma \rightarrow \infty$ so that we recover the smooth-front solution $\tilde{v} \rightarrow \Gamma \left[1 + \frac{1}{\Gamma} - 1 \right] = 1$ or,

equivalently, $v_{\max} = \frac{\Phi}{BL_r}$ in reduced units. The function $\tilde{v}_{>}(\Gamma)$ is defined for all Γ and is

plotted below in Figure C-2.

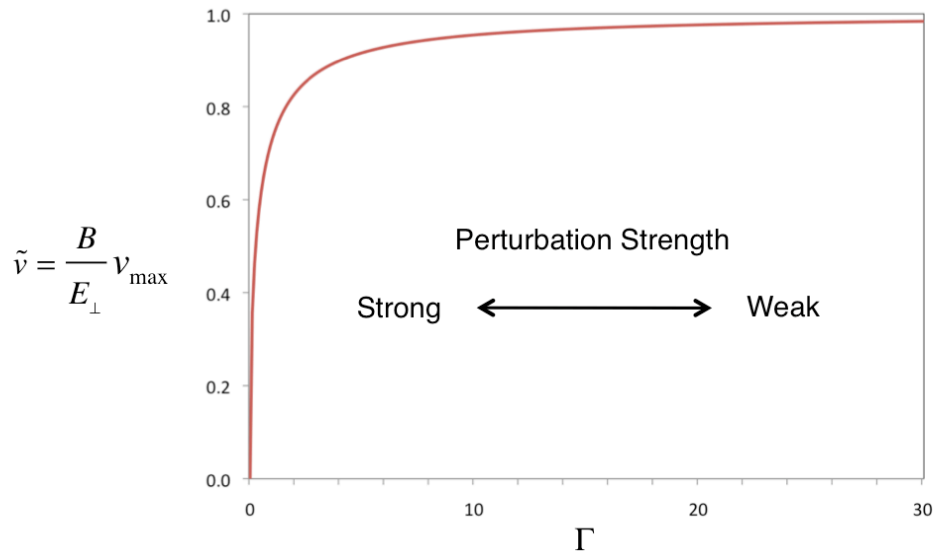


Figure C-2 Perturbation Function at a Crest

Perturbation strength decreases with increasing Γ while the velocity function asymptotes to unity

Conversely, the convex crest shown schematically in Figure C-3 has the force balance equation $\tilde{v}^2 - 2\Gamma\tilde{v} + 2\Gamma = 0$ where we apply the downstream magnetic field strength, since we assume the ion escapes through the back of the ramp. Again, we choose the physical solution $\tilde{v}_< = \Gamma \left[1 - \sqrt{1 - \frac{2}{\Gamma}} \right]$, required in order to satisfy the limit of vanishing perturbation i.e., $\lim_{\Gamma \rightarrow \infty} \tilde{v} = 1$.

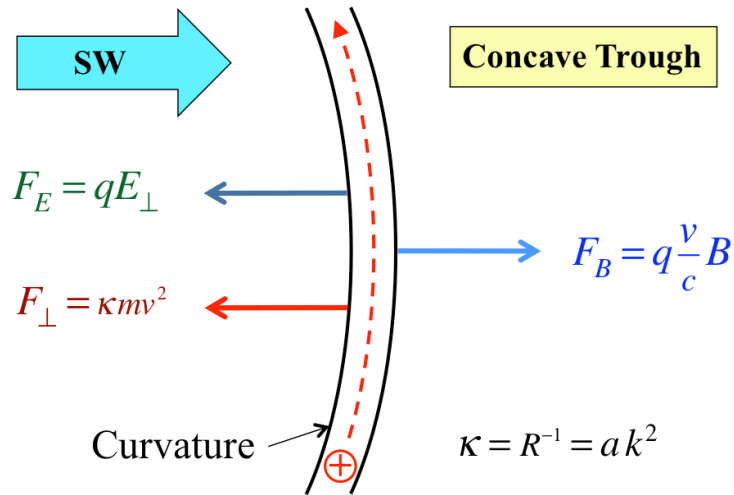


Figure C-3 Force Balance at a Concave Trough
The center of curvature is to the left

Note that the solution is defined only when $\Gamma \geq 2$, limiting the region of application, as shown in the plot of Figure C-4.

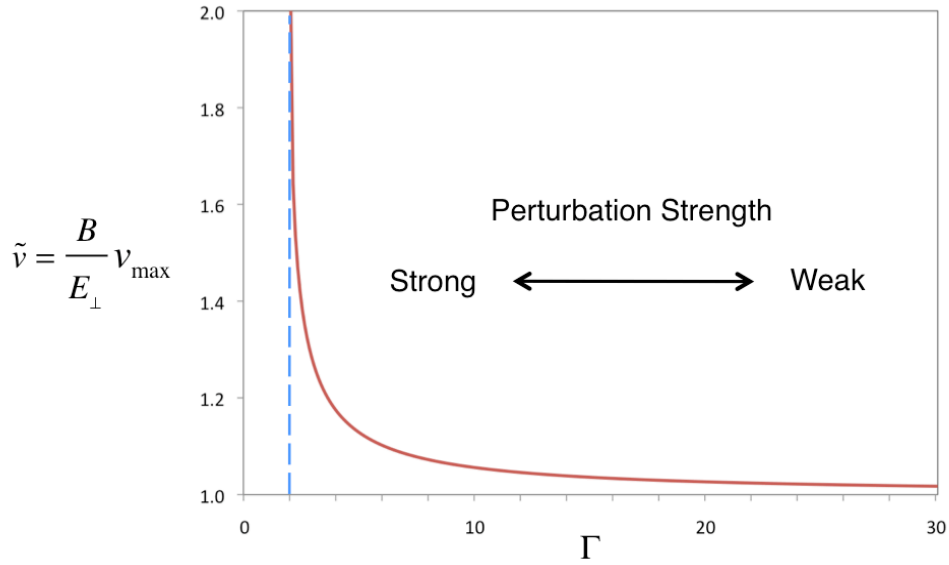


Figure C-4 Perturbation function at a Trough

The function is not defined for $\Gamma < 2$

This is equivalent to restricting the perturbation conditions such that

$$ak^2 < \frac{qB_2^2}{4E_{\perp}mc^2} = \left(\frac{qB_1^2 R_C^2}{4\Phi c^2} \right) \frac{L_r}{m}. \text{ The electrostatic potential step } \Phi \text{ and the shock}$$

compression ratio R_C are determined by the choice of Mach number and plasma beta.

We assume singly ionized species so that, by our choice of normalization, both the charge q and the upstream magnetic field B_1 are unity. Rewriting this constraint as

$$\frac{mak^2}{L_r} < K, \text{ a constant, we see that our predictive function becomes undefined with}$$

increasing mass at narrow shock ramps for all but the weakest perturbations.

Specifically applying the parameters at the Earth's bow shock at Mach 5, we have

$$K = \frac{R_C^2}{4\Phi} = \frac{(2.62)^2}{4(10.6)} \simeq 0.162 \text{ so that, for protons encountering a steep perpendicular}$$

shock of width $L_r = 0.05$, even a very weak perturbation $a = k = 0.1$ would be an order

of magnitude too large, as the function would be undefined. Hence, we abandon this formulation and adopt the ‘crest’ function derived above, which we shall refer to as ‘Fcn1’.

Penetration Depth

One tacit assumption made in deriving the ‘crest’ force balance equation

$$\tilde{v}_> = \Gamma \left[\sqrt{1 + \frac{2}{\Gamma}} - 1 \right]$$

was that the very fastest ions at smooth fronts do not undergo normal oscillation, but remain at the leading edge of the ramp until $qvB_1/c > qE_\perp$.

Should this condition no longer hold at strongly perturbed shocks, we must allow for variation in the magnetic field strength as a function of penetration depth, since the field increases through ramp to its downstream value $B_2 = R_C B_1$. We do so by explicitly replacing the magnetic field strength $B \rightarrow \tilde{B} = \xi(x) B_1$ where $1 < \xi(x) < R_C$ is a measure of the penetration depth as $x:0 \mapsto L_r$, and the upstream field B_1 is normalized to unity.

This may be introduced into Fcn1 $\tilde{v}_{\max} = \Gamma_1 \left[\sqrt{1 + \frac{2}{\Gamma_1}} - 1 \right]$ via the variable $\Gamma \equiv \frac{\xi^2 L_r}{2\Phi \tilde{m}_i a k^2}$,

which we have here explicitly defined in terms of the potential, and where the ‘fitting parameter’ ξ represents the penetration depth. After several trials, we found that the fitting function $\xi = 1 + L_r$ gave the best global results and, while this seems somewhat abhorrent as an ‘ad hoc’ fitting parameter, it at least has the virtue of being a single, ramp dependent function.

Force-balance Function Agreement

We begin by comparing the force-balance function Fcn1 with data at four perturbation wavenumbers $k = 0.1$, $k = 0.3$, $k = 0.5$ and $k = 0.7$ at a Mach 5 shock of thin ramp $L_r = 0.05$, equivalent to two electron inertial lengths $\sim 2c / \omega_{pe}$. This represents the thinnest ramps reasonably possible within the heliosphere, and we immediately see in Figure C-5 below that this function begins to fail almost immediately with increasing perturbation strength.

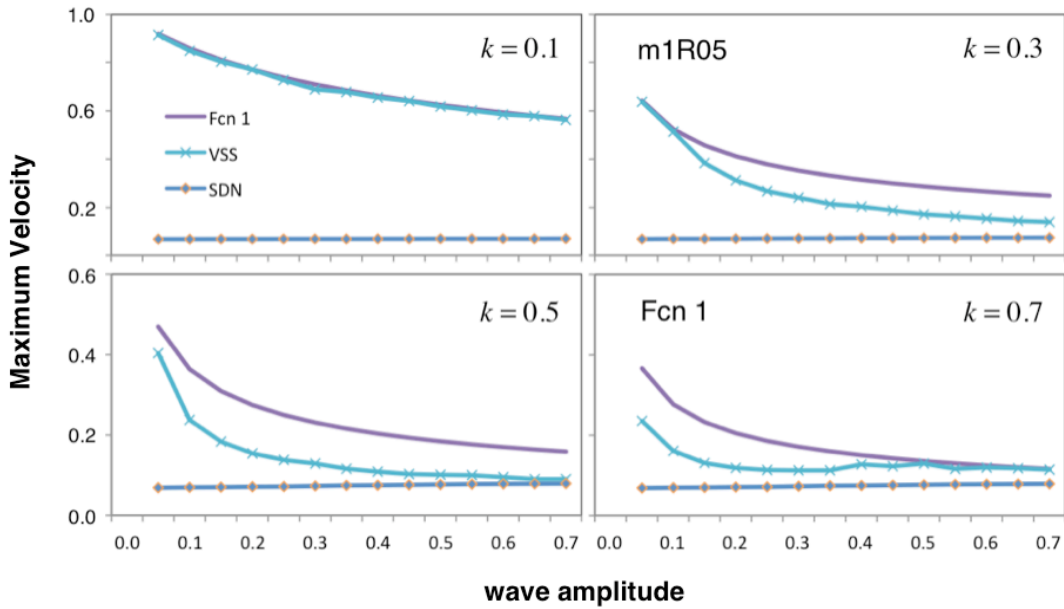


Figure C-5 Fitting Function Fcn1 versus Data

Proton velocity maxima at Mach 5 shock with thin ramp $L_r = 0.05$ are plotted against amplitude at four perturbation wavenumbers. Velocity 'Vmax' is normalized to the unperturbed velocity maximum attainable for this ramp. The vertical scale in the bottom panels is expanded by 40%

Examination of data over the entire parameter range explored confirms that the force-balance solution fails at thin ramps with strong perturbations, more specifically, when $m_i a k^2 > 0.18 L_r$, as indicated by the red arrows in the broadly representative plots as shown in Figure C-6 below. This may indicate that, for strong perturbations at thin ramps, the fastest surfing ions do not traverse the shock front inside the ramp, as initially assumed.

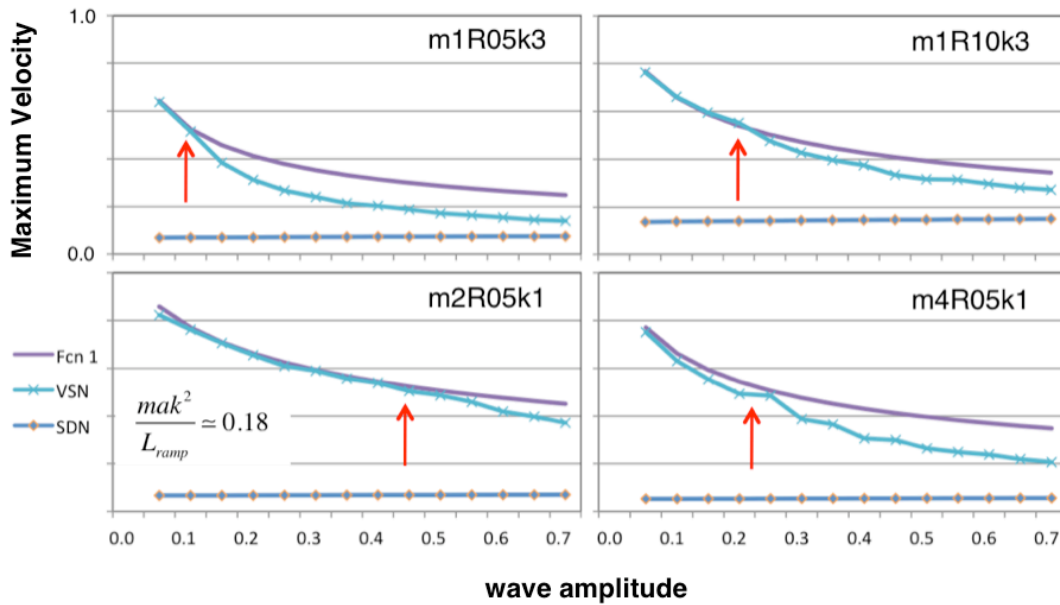


Figure C-6 Force Balance FCN1 Failure

Acronyms indicate ion mass, ramp width and ripple wavenumber e.g.,

m1R05k3 represents $m = 1$ (protons) ramp width $L_r = 0.05$ and

wavenumber $k = 0.3$. Deviation begins when $mak^2 \approx 0.18 L_r$. All panels

have identical scales.

This may possibly be explained by the solution derived for the ‘trough’ force balance

equation, which was undefined when the perturbation parameter $\Gamma = \frac{B^2 L_r}{2\Phi m a k^2} < 2$. Even

though we employed the ‘crest’ formulation, no account was made for deep penetration into the ramp at troughs, aside from the inclusion of $B(x) = \xi(x) B_1$ as a fitting parameter. As we have already shown, at Mach 5 the ‘trough’ function should fail when $mak^2 > 0.162 L_r$, and agreement between this predicted failure and that observed is remarkably close. Since no adjustment of the fitting parameter could overcome this shortcoming, the discrepancy suggests that the force balance model is sensitive and may further indicate that some relevant physical mechanisms are missing from this treatment. We therefore abandon Fcn1 as a predictive function, and seek an alternative model.

Curvature Model – Function 2

We have just treated the case where an ion’s trajectory remains inside the shock ramp, constrained by the shock-normal retarding electric field. As an alternate approach, we now consider what happens when an ion escapes the ramp upstream where it is free to follow an unconstrained guiding center trajectory. The following derivation is based on the assumption that, once the radius of curvature of an accelerated ion’s trajectory exceeds that of the shock front, the normal velocity gain resulting from free-flight velocity mixing during its excursion away from the front will be sufficient to surmount the electrostatic boundary upon its next shock encounter, and thus terminate the surfing process. This is consistent with the analysis of Lee et al, since one cornerstone of their theory is the requirement that the bounce amplitude should decrease with time. We now seek a functional form to which we can fit the data in such cases.

Trajectory curvature

We begin by calculating the radius of curvature of upstream trajectories where the bulk fluid velocity and magnetic field are uniform. Ion trajectories free from the retarding electric field within the shock ramp obey the parametric velocity equations

$$v_x(t) = -v_{\perp} \cos \Omega_{ci} t + u_1 \quad \text{and} \quad v_y(t) = v_{\perp} \sin \Omega_{ci} t$$

C-7.

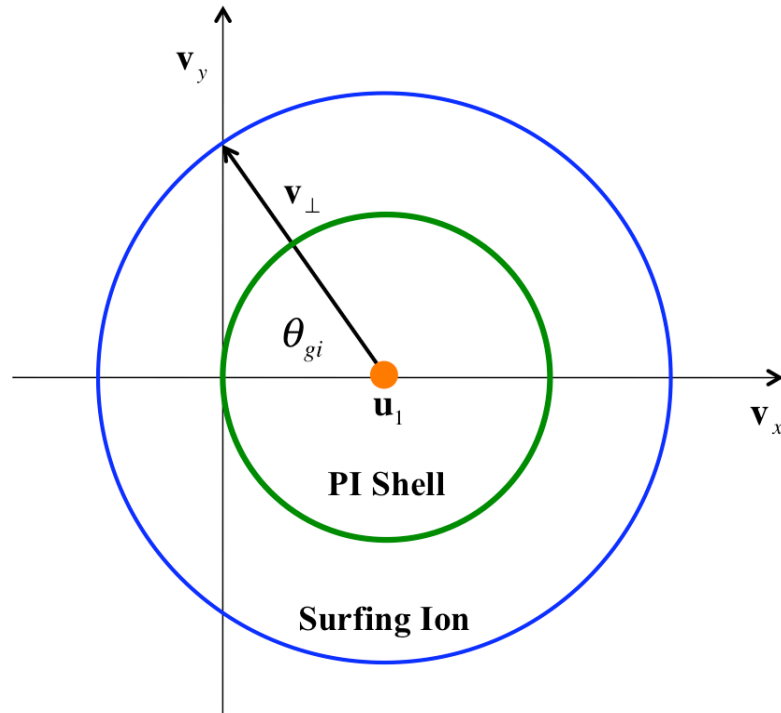


Figure C-7 Velocity space construction

Outer circle represents accelerated upstream ion free of the shock ramp electric field. The inner circle represents fastest incident PI with both species centered at the upstream SW velocity u_1 . Free ion phase angle evolves as $\theta_{gi}(t) = \Omega_{ci} t$. Velocity of front-confined ions is constrained to the v_y -axis. Acceleration by convective electric field expands the free-velocity shell.

We integrate the velocity equations in the shock frame to obtain the parametric equations for the guiding center trajectory $x(t) = -r_{gi} \sin(\Omega_{ci1} t) + u_1 t$ and $y(t) = -r_{gi} \cos(\Omega_{ci1} t)$ where $r_{gi} = v_{i\perp} \Omega_{ci1}^{-1}$ is the upstream cyclotron radius specific to the ion species and its velocity about the guiding center, perpendicular to the ambient magnetic field. For such parametrically defined functions, the Radius of Curvature, that is, the inverse of the curvature function κ , is given by

$$\rho_{ion} = \frac{1}{\kappa} = \frac{\left[(\dot{x})^2 + (\dot{y})^2 \right]^{3/2}}{|\dot{x} \ddot{y} - \dot{y} \ddot{x}|} = \frac{\left[v_x^2 + v_y^2 \right]^{3/2}}{|v_x \dot{v}_y - v_y \dot{v}_x|}$$

where the dots represent differentiation with respect to the argument. Substituting the time derivatives, we have

$$\rho_{ion} = \frac{\left[v_{\perp}^2 - 2u_1 v_{\perp} \cos \Omega_{ci1} t + u_1^2 \right]^{3/2}}{|\Omega_{ci1} v_{\perp}^2 - u_1 v_{\perp} \Omega_{ci1} \cos \Omega_{ci1} t|} = r_{gi} \frac{\left[1 - 2\mu \cos \Omega_{ci1} t + \mu^2 \right]^{3/2}}{|1 - \mu \cos \Omega_{ci1} t|}$$

where $\mu \equiv \frac{u_1}{v_{\perp}}$. In order to obtain a closed algebraic form for a predictive velocity

attenuation function, we may simplify this expression by assuming the particle has only just escaped confinement within the ramp. In that case, even though v_y may have increased substantially, the normal velocity v_x must lie very near the v_y axis, and the cosine of the instantaneous gyro-phase angle θ_{g0} may be approximated by

$\cos \theta_{g0} \simeq \frac{u_1}{v_{\perp}} = \mu$ as seen in construction of Figure C-7. As the ion velocity increases

$\mu \rightarrow 0$ whence $\rho_{ion} \simeq r_{gi} (1 - \mu^2)^{1/2} \simeq r_{gi} \left(1 - \frac{1}{2} \mu^2\right)$ to second order, while at sufficiently

high velocity we simply have $\rho_{ion} \rightarrow r_{gi}$.

Recall that our initial force-balance model was static, and made no provision for ion dynamics inside the ramp. Note also that severe corrugations must necessarily induce normal oscillations as the ion traverses along the front within or without the ramp. Hence wide ramps or large perturbations may limit the predictive scope of our analytic approaches. We next calculate the curvature of the front itself, complete the derivation with a few minor corrections, and see whether this function may prove acceptable in predicting the data.

Front curvature

We now consider a static corrugated front given by $x(y) = a \cos ky$ whose first and second derivatives with respect to the argument are $x'(y) = -ak \sin ky$ and $x''(y) = -ak^2 \cos ky$ respectively. For functions of this form, the Radius of Curvature is

given by $\rho_{front} = \frac{(1 + x'^2)^{3/2}}{|x''|}$ and, substituting the derivatives yields

$$\rho_{front} = \frac{(1 + a^2 k^2 \sin^2 ky)^{3/2}}{|ak^2 \cos ky|} \text{ which has minima when } ky = n\pi, \text{ whence } \rho_{min} \simeq \frac{1}{ak^2}.$$

Recall that the condition for the trajectories to remain at the front is that $\rho_{ion} \leq \rho_{front}$.

Applying our criterion that, to be of any consequence, surfing ions must exceed the velocity ceiling of shock drift, $v_y \geq 3u_1$, we see that the variable μ must satisfy

$\mu = \frac{u_{\perp}}{v_{\perp}} < \frac{1}{3}$. As the reader can easily verify, the substitution $v_{\perp} \approx v_y$ affects the solution

by at most six percent, and we may obtain a reasonable estimate for the velocity ceiling

– even under adverse conditions – by setting $\frac{v_y}{\Omega_{ci1}} \leq \frac{1}{ak^2}$. In our normalized units

$c = q = 1$, and we find that $v_{pert} \leq \frac{B}{mak^2}$.

Smooth front divergence correction

The velocity ceiling function just derived takes no account of the electric field within the ramp, and obviously diverges in the limit of vanishing perturbation. As we have shown, surfing velocity at a smooth front is always limited by the ramp width

$v_{ramp} \leq \frac{E_{\perp}}{B} = \frac{\Phi}{BL_r}$ when the Lorentz force exceeds the retarding electrostatic field, which

we may incorporate by employing the functional form $\frac{1}{v_{max}} = \frac{1}{v_{ramp}} + \frac{1}{v_{pert}}$. Note that while

large perturbations dominate over the ramp limit, the expression correctly returns the correct ramp-dependence when ripples vanish. As we did with Fcn1, to better compare our results we normalize with respect to the maximum attainable velocity for a given

smooth ramp width, $\tilde{v}_{max} = \frac{BL}{\Phi} v_{max}$, which gives us

$$\tilde{v}_{max} = \left[1 + \frac{\Phi mak^2}{B^2 L_r} \right]^{-1} = \left[1 + \frac{2}{\Gamma} \right]^{-1}.$$

Free flight Energy gain correction

As we showed in the baseline studies of the smooth front, surfing ions gain roughly double the piston energy $\Delta\epsilon \simeq m_p u_1^2$ with each reflection. Since our convex function models the precursor to upstream escape from the ramp, we expect that some additional energy may be gained during the ion's free flight before it subsequently encounters the front. There are several ways to predict the duration of this last flight, which would in turn give us a velocity gain under the influence of the convective electric field. However, we shall assume for simplicity that the shock drift energy ceiling will suffice as a rough estimate. The final surfing ion energy should then be

$$\epsilon_{fin} = \epsilon_{max} + \Delta\epsilon_{flight} \text{ or } \frac{1}{2}m_i v_{fin}^2 = \frac{1}{2}m_i v_{max}^2 + \frac{1}{2}m_p v_{SD}^2 \text{ which gives an adjusted final velocity}$$

$$\tilde{v}_{fin} = \tilde{v}_{max} \sqrt{1 + m_i^{-1} \left(\frac{\tilde{v}_{SD}}{\tilde{v}_{max}} \right)^2} \text{ where the proton mass is normalized to unity, } \tilde{v}_{max} \text{ is the}$$

$$\text{curvature-limited velocity derived above } \tilde{v}_{max} = \left[1 + \frac{2}{\Gamma} \right]^{-1}, \text{ and } \tilde{v}_{SD} \simeq \frac{BL_r}{\Phi} v_{SD} \text{ is the ramp-}$$

normalized velocity ceiling. The absolute SD ceiling $v_{SD} \simeq (4 - 3L_r)u_1$ was found

empirically from the baseline studies, but did vary somewhat with ion mass down to

$v_{SD} \sim 3u_1$. However, the precise form of this surfing velocity 'floor' function is of little

moment since, once the velocity has been attenuated that far, the surfing mechanism

ceases to be effective at all. Indeed, any estimate in the range $3u_1 \leq v_{SD} \leq 4u_1$ should

suffice for all cases of interest since this free-flight velocity correction is small for narrow

unperturbed ramps when the surfing velocity is high, and only becomes prominent when

the velocity has already been substantially suppressed. Note also that there was no

need to extend this correction to force-balance solution Fcn1 since it was derived under the explicit assumption that surfing ions remain ramp-bound as they approach terminal velocity prior to downstream escape.

Penetration depth correction

While the ramp penetration ‘fitting function’ may have been appropriate for Fcn1, which was explicitly derived for trajectories inside the ramp, we now argue in favor of extending this to Fcn2 even though this was based on expulsion from the front into free-flight. It is possible that an ion trajectory prior to ejection the front may experience transient penetration into the ramp where the magnetic field is stronger – especially at higher perturbation strengths that may induce normal oscillations – and we expect that this may possibly be more pronounced at wider ramps. Further, there is no reason to assume that ions reaching terminal velocity necessarily do so at the leading edge of the ramp, but may in fact be anywhere inside. We therefore justify extending this fitting parameter to Fcn2 noting that, once ions escape the ramp upstream, our free gyrophase argument still holds.

Final Functional Form

With inclusion of these corrections, we may now write the curvature-limited

velocity ceiling function as $\hat{v}_{\max} = \left[1 + \frac{2}{\Gamma_2}\right]^{-1} \sqrt{1 + m_i^{-1} v_{SD}^2 \left[1 + \frac{2}{\Gamma_2}\right]^2}$ in the reduced

variable $\Gamma_2 \equiv \frac{\xi_2^2 L_r}{2\Phi \tilde{m}_i a k^2}$ with ‘fitting’ parameter ξ_2 , which represent the magnetic field

strength correction as a function of ramp penetration depth.

Finally, since surfing ions by definition undergo multiple reflections, we expect a minimum velocity ‘floor’ corresponding to a kinetic energy of order twice the piston energy. This velocity floor was seen in the baseline studies of Chapter 2 at very wide ramp widths, and should be taken into account when composing our predictive functions. Even though we have taken into account a ‘free-flight’ energy gain correction, for completeness we will include this hard floor in our curvature function

$$v_{\max} = \max(\text{Fcn2}, v_{SD}), \text{ where } v_{SD} \text{ is again defined as above.}$$

Agreement with Data

We see from the parameter set shown in Figure C-8 that the curvature-based function Fcn2 tracks the data well at the thin ramp width and large perturbation strengths precisely where Fcn1 failed. This should be expected, since the postulates applied in deriving Fcn2 were based on ions following trajectories outside the ramp, where large ripples were seen to propel ions.

As may be verified from the data exhibited at the end of this appendix, agreement with data is remarkably good over the entire range of shock, ion and perturbation parameters. This is especially noteworthy considering that only one ‘fitting’ constant $\xi_2 = 1.2$ for this function as been used throughout.

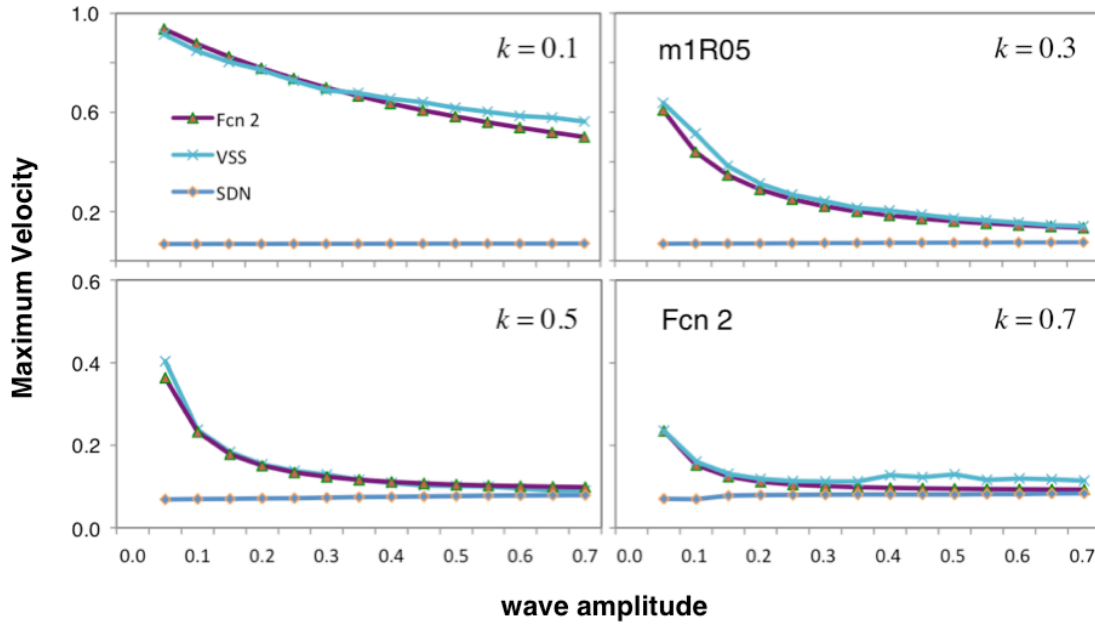


Figure C-8 Curvature Velocity Function Fcn2 vs. Data

Data for protons is shown for Mach 5 shock and thin ramp $L_r = 0.05$

showing agreement with 'Fcn2' which improves with increasing perturbation strength. Vertical scale in bottom panels is expanded by 40%

We also offer exemplary plots comparing both predictive functions with simulation data over broader parameter samples and stronger perturbations as shown in Figure C-9, and see that both appear to work reasonably well over the entire exhibited range until the perturbation becomes strong enough to suppress the velocity to values near the shock drift ceiling $v_{SD} \sim 3u_1$.

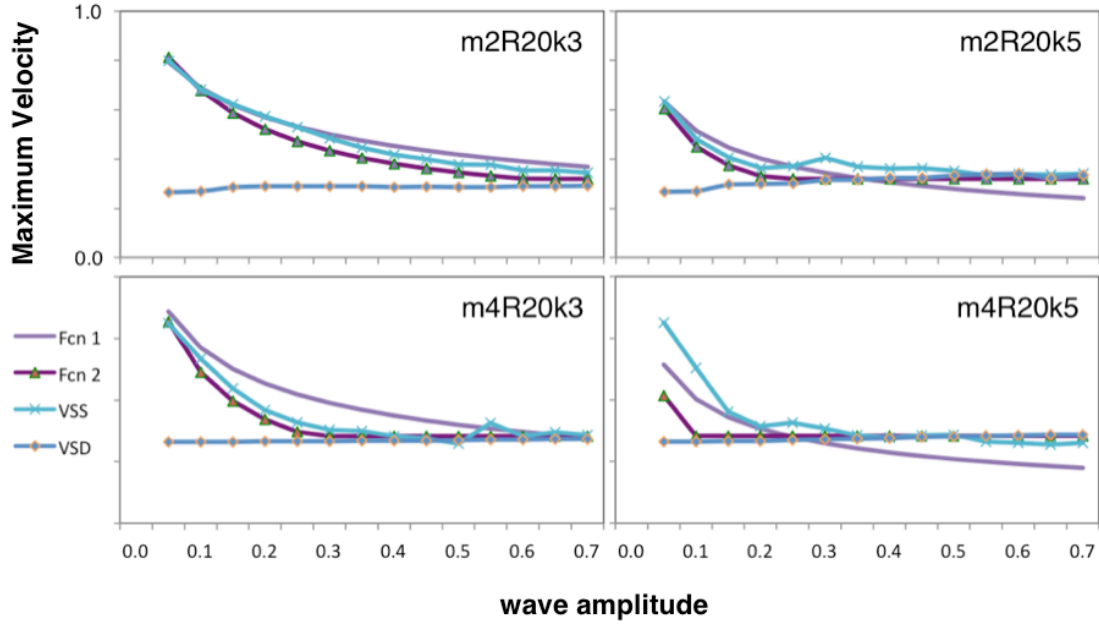


Figure C-9 Comparison of fitting functions Fcn1 and Fcn2.

The top panels represent deuterons $m = 2$ with singly ionized helium $m = 4$ below. Wavenumbers $k = 0.3$ are on the left panels and $k = 0.5$ on the right. The ‘fitting’ constant $\xi_2 = 1.2$ for Fcn2 was used throughout for all variation of parameters. The four panels have identical vertical scales since they represent the same ramp width $L_r = 0.20$.

It should be noted parenthetically that the deviation from unity in the constant ξ_2 applied to Fcn2 is smaller than might appear, since the function range is $1 < \xi(x) < R_c$ where $\xi(x)$ represents the magnetic field strength at a ramp penetration depth $0 < x < L_r$ and where $R_c \approx 2.62$ is the compression ratio at Mach 5 equivalent to the field strength at the trailing edge of the shock ramp. Rewriting the function

$B(x) = \xi(x)B_1 = 1 + (R_c - 1)\frac{x}{L_r}$ with $\xi(x) = 1 + \Delta(x)$, we find the penetration depth

$$\Delta(x) = 0.20 \Rightarrow \frac{x}{L_r} = \frac{0.20}{1.62} \approx 0.12, \text{ a little more than 10\%}.$$

Finally, we validate in Figure C-10 our inclusion of the velocity floor term in Fcn2 when the surfing velocity approaches the shock drift ceiling $v_{\max} \rightarrow v_{\text{floor}} \sim v_{SD} \sim 3u_1$. We remind the reader that this is consistent with the minimum energy gain from a single excursion away from the front, as discussed in the analysis and derivations above.

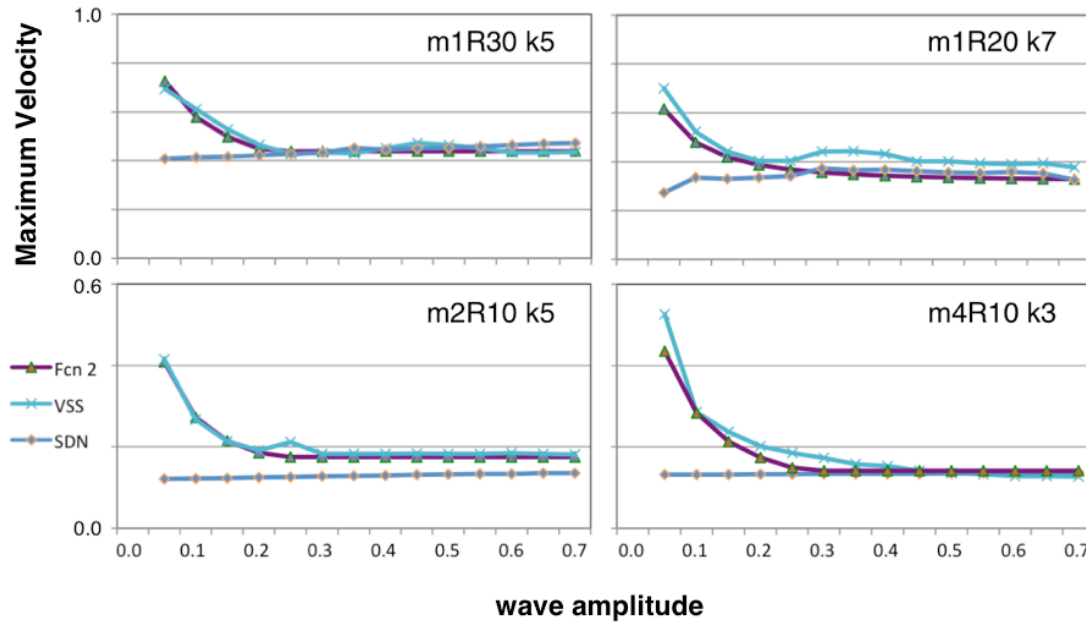


Figure C-10 Surfing Velocity Floor

These four plots demonstrate agreement between the velocity floor term and data over a range of strong perturbations.

Fitting Function Comments

The primary objectives of this exercise are to demonstrate the disruptive effect of ripples on the shock surfing mechanism and verify to what extent our physical model explains the results. We also seek to produce a reasonably accurate predictive function for the velocity ceiling, with which to evaluate shock surfing's ability to attain the DSA pre-injection threshold from observational data.

The curvature function Fcn2 using the constant 'fitting' parameter $\xi_2 = 1.2$ appears to work globally, and thus provides us with a reliable tool with which to evaluate shock surfing as an effective DSA pre-injection mechanism, assuming that such an injection velocity threshold actually exists. We have also seen that Fcn1 works extremely well at wider ramps. However, good prediction is essential for strong 'turbulence' at thin ramps where ions have the best opportunity to achieve high velocities but, at the same time, are likely to suffer the greatest disruption due to front ripples. From our baseline studies, we have seen that wide ramps produce low surfing velocities and hence are of little interest for DSA pre-injection, and therefore acknowledge Fcn1 as a model-validating tool, but dismiss it as a predictive tool for this purpose.

Comments on Mach Number

Recall that the ramp limit on surfing ceiling given in terms of absolute velocity

$v_{\max} = \frac{\Phi}{BL_r}$, while the potential step calculated from $q\Phi = \frac{1}{2}m_p u_1^2 (1 - R_C^{-2})$ may be given

in reduced units as $\Phi = K M_A^2$ where $K = \frac{1}{2}(1 - R_C^{-2})$ becomes nearly constant since the compression ratio R_C rapidly asymptotes with increasing Mach number. Fixing the

magnetic field to unity then gives $v_{\max} \sim \frac{M_A^2}{L_r}$, which at first appears to be a strong Mach

dependence indeed. However, for the purposes of DSA injection, it is the threshold

velocity relative to the upstream wind velocity $\tilde{v}_{inj} = \frac{v_{inj}}{u_1} \gg 3\sqrt{1+\eta^2}$ that matters.

Expressing the surfing velocity in reduced units, we have $\tilde{v}_{\max} = \frac{v_{\max}}{M_A} \sim \frac{M_A}{L_r}$ so that the

effective ceiling is proportional to the Mach number, not the Mach squared. Hence the Alfvén Mach 5 results are representative of the majority of shocks likely encountered in the heliosphere within a range of less than 10%, since the magnetosonic Mach number for unity plasma beta at Alfvén Mach 5 would be around seven.

We now consider acceleration time, and limits thereupon placed by periodic shock reformation. This limit was derived as $v_{\max}(\tau) = \frac{q}{m} E_{\parallel} \tau = u_1 \frac{qB_1}{m} \tau$ which we again express in reduced units as $\tilde{v}_{\max}(\tau) = \frac{\tilde{v}_{\max}(\tau)}{u_1} = K\tau$ where the constant $K = \frac{qB_1}{m}$.

Hence, in terms of units reduced appropriately for DSA injection, the velocity with respect to the injection threshold is independent of the shock Mach number, and the limits we have demonstrated must apply to higher Mach number shocks as well.

Unlike the rigid constraints we imposed upon our model shocks to freely probe parameter space, real astrophysical shock parameters are interconnected by a host of local phenomena often generated by accelerated ions themselves. Indeed, such feedback mechanisms contribute to the long-term stability of standing shocks, as noted by Wu et al (2007) and Winske et al (2007) who modeled the Termination Shock. But the mechanism is cyclic: the shock front is broadened by the population of reflected ions

until the potential ramp is too weak to effect reflection, then narrows until reflection begins anew. Whether this reformation periodicity persists at extremely high Mach number shocks such as those associated with supernovae is at this moment unknown. However, for the present consideration of shocks within the heliosphere, we may safely conclude that our results should be generally representative.

Comments on Wavenumber

We also claim that we need not continue our investigation to higher wavenumbers, since $k = 0.7$ corresponds to a corrugation wavelength $\lambda \simeq 9.0$, slightly shorter than the ion cyclotron radius of the fastest pickup protons. In other words, we do not expect wavelengths substantially shorter than this to be supportable especially since the ‘gyrating’ upstream pickup ion population has often been cited in the literature as a source when discussing shock features such as the Earth’s bow shock ‘foot’. Indeed, the wavenumber ceiling constraint may be generalized by the condition $\lambda(k) \geq r_{gi}(2u_1)$, which in normalized units gives $k \leq \frac{\pi}{M_A}$ or $k_{\max} \simeq 0.63$ for a Mach 5 shock. As we discuss later in detail, this criterion is met precisely at the Bow Shock, as was observed *in situ* by Cluster.

Another point of note when considering a wavenumber ceiling is that both Kolmogorov turbulence theory and *in situ* observations by the Cluster Mission (where the turbulence is far from fully-developed) show that, due to cascading, mechanical energy is primarily concentrated at longer wavelengths. Assuming that some form of coupling mechanism is responsible for front corrugations, we expect the amplitude at higher wavenumbers to be fairly weak in any case, even if no resonant physical limiting

mechanism were to exist. We can easily quantify this, and do so in depth in the next Chapter. The spectral index for magnetic turbulence at the Bow Shock observed by Cluster is approximately an inverse square, more specifically $\varepsilon_B \sim k^{-1.9}$. Fully developed Kolmogorov turbulence is nearly the same at $\varepsilon_B \sim k^{-5/3}$. Coupling by equipartition between turbulent magnetic energy and kinetic energy of oscillations along the shock front would lead to the dependence $\varepsilon_K \sim \varepsilon_B \Rightarrow \frac{1}{2}\rho\Omega^2 a^2 \sim \frac{\delta B^2(k)}{8\pi}$ or $a^2 \sim \delta B^2$

where a is the amplitude of the corrugation. We now have $a^2 \sim k^{-\gamma}$ and are free to choose the spectral index appropriate to the problem. Since the difference is small, we choose $\gamma = 2$ for convenience so that our surfing velocity suppression factor in terms of wavenumber alone becomes $(mk)^{-1}$ making the velocity ceiling dependence inversely proportional to both mass and wavenumber.

Conclusions and Discussion

The fact that we have modeled a perturbed shock with quasi-static corrugations should not materially affect our results or predictions. We implicitly assumed dynamic perturbations, but with times scales which appear long in the surfing ion's frame. This may be understood by considering both the limiting Alfven wave speed along the front as well as the frequency regime of possible front oscillations. Once an ion has considerably exceeded the shock velocity, any such time dependence or wave speed will appear at worst quasi-stationary in the particle's frame of reference.

Whether such corrugations are caused by resonant coupling to turbulence in the background plasma, by currents via the modified two-stream instability or by manifold other mechanisms occurring at the front is of little consequence to this investigation. As

we show in the next chapter, both recent simulations and in situ observations at several shocks have shown that these ripples do in fact exist, may be ubiquitous and, as we have just seen, impose serious constraints upon the shock surfing mechanism.

It is unfortunately impossible to verify the predictive power of this function *in situ* since, to date, there has been no observational data of streaming ions directly attributable to shock surfing. Although we have failed to achieve a high degree of precision in predicting the simulation results, we feel confident that our functions are qualitatively accurate enough for us to conclude that perturbations disrupt shock surfing to the extent shown, and that no further post-mortem is required. Since our model is highly idealized, missing many real shock features including waves and temporal fluctuations, the depth of our mathematical treatment need not exceed the physical details of our model nor the limitations imposed by our approximations. As we will show, the very simplest estimates coupled with observational evidence will prove more than ample for dismissing shock surfing as a viable pre-injection mechanism. Hence there is no motive in further seeking to produce a practical application out of our predictive function.

Finally, we claim that the details of our model have been kinder to surfing ions than a real shock corrugation would be. The consequences of non-normal flow through the ramp, where the convective electric field would not follow the front surface at the leading edges of crests, make the force balance much worse. Hence our model represents the most optimistic case for shock surfing so that, if shock surfing fails under these idealized assumptions, it must certainly fail under the conditions found at real astrophysical shocks.

Data Section

In the following pages, we present a comprehensive set of data with both predictive functions (force-balance function Fcn1 and curvature function Fcn2) superimposed for comparison. All vertical velocity scales are normalized to the maximum attainable at the given ramp width without front perturbations, and all examples given for a perpendicular shock at Mach 5. Ripple amplitudes are given in units of the ion inertial length $\lambda_{pi} = c / \omega_{pi}$ and wavenumbers to $2\pi / \lambda_{pi}$.

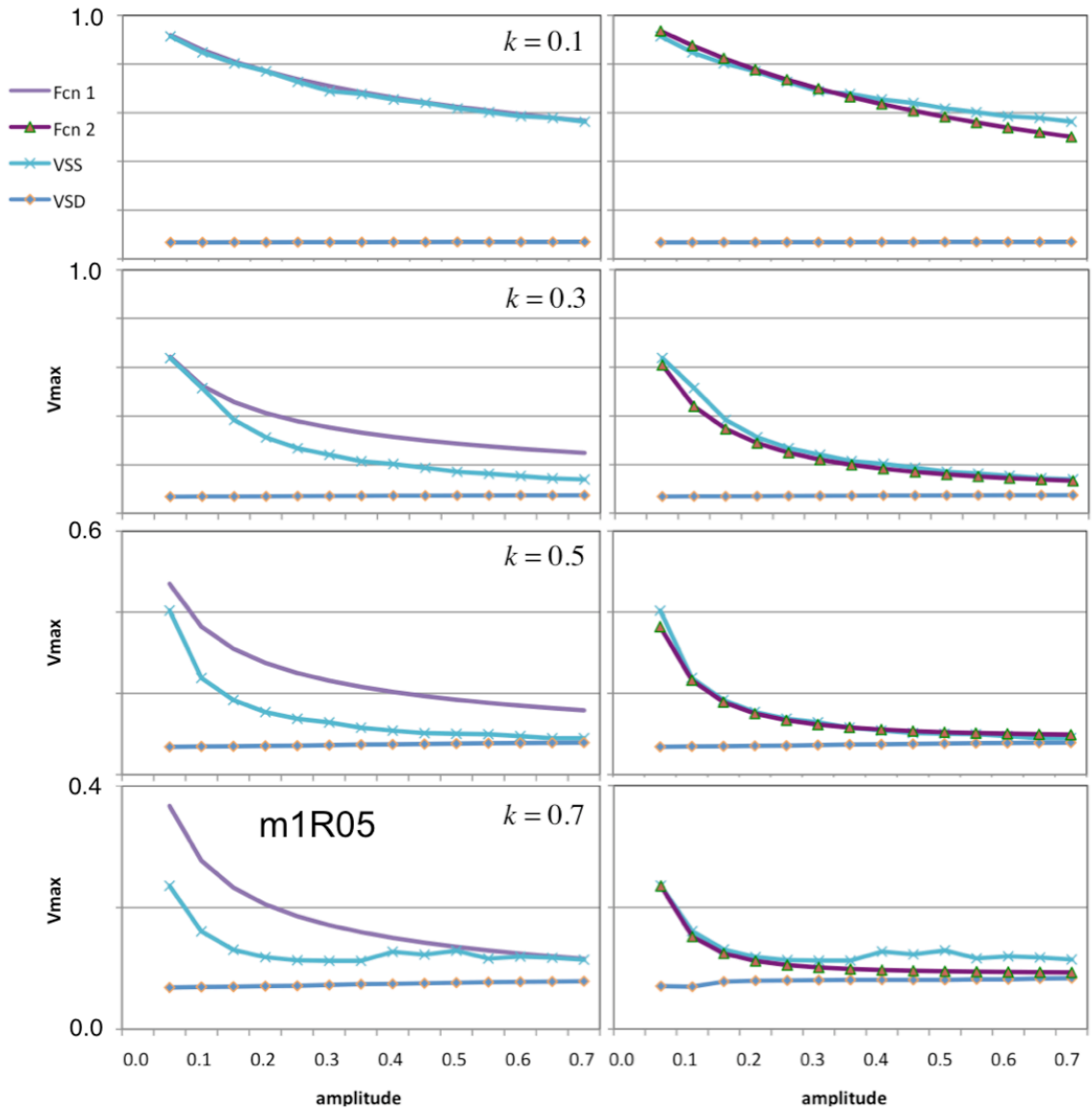


Figure C-11 Velocity Data and predictive functions

Protons H^+ at ramp width $L_r = 0.05 \lambda_{pi}$

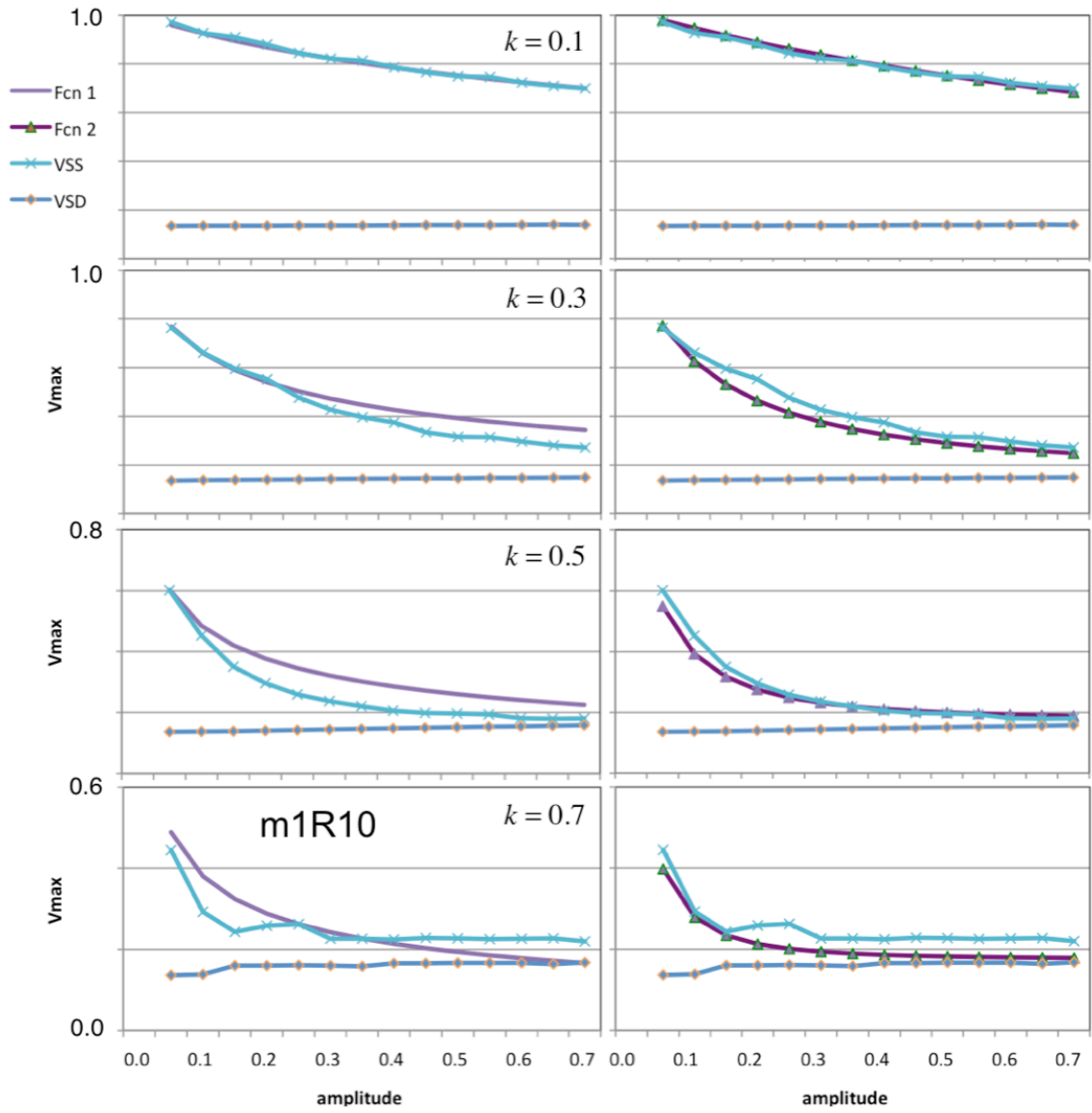


Figure C-12 Velocity Data and predictive functions

Protons H^+ at ramp width $L_r = 0.10 \lambda_{pi}$

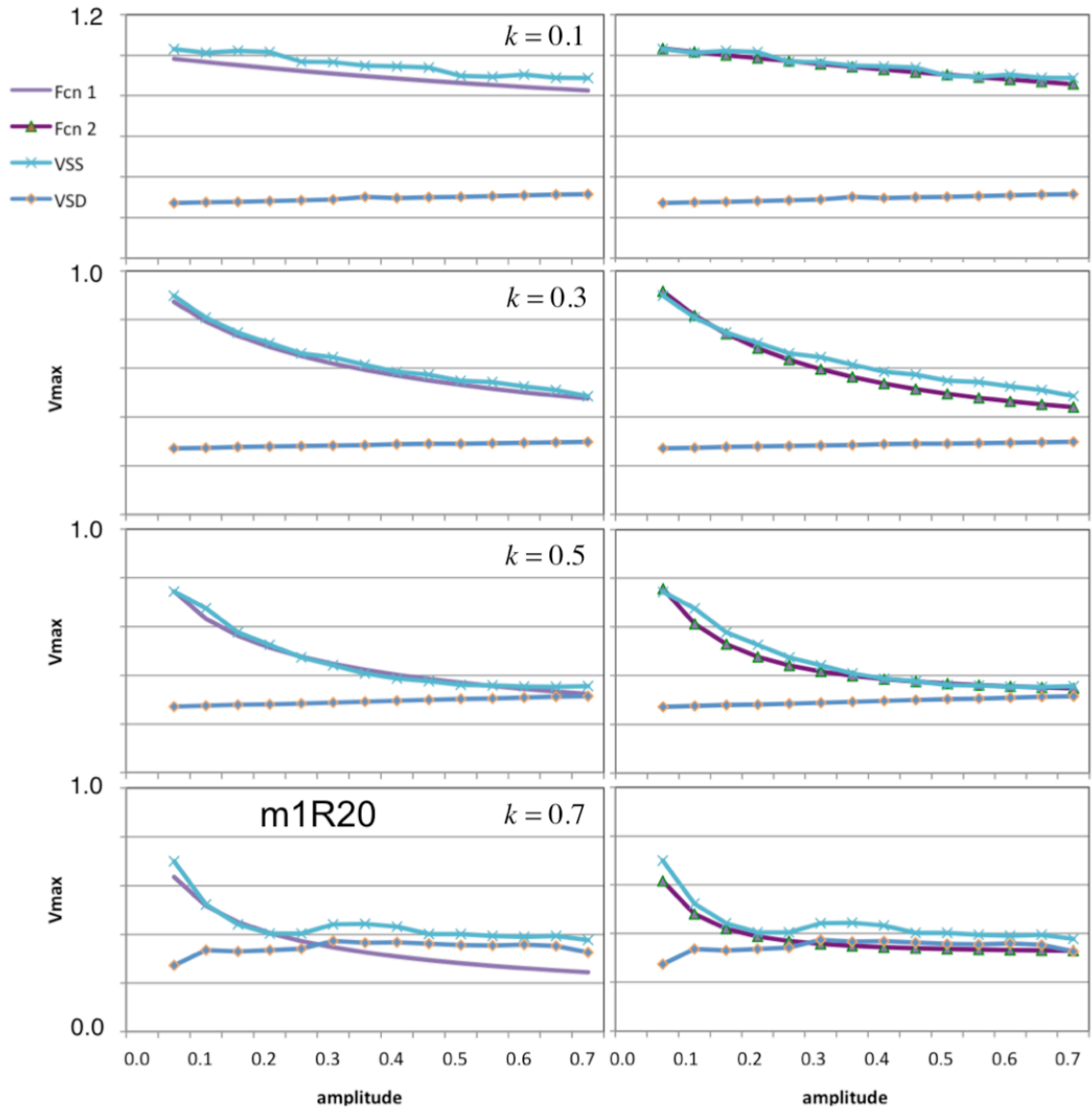


Figure C-13 Velocity Data and predictive functions

Protons H^+ at ramp width $L_r = 0.20 \lambda_{pi}$

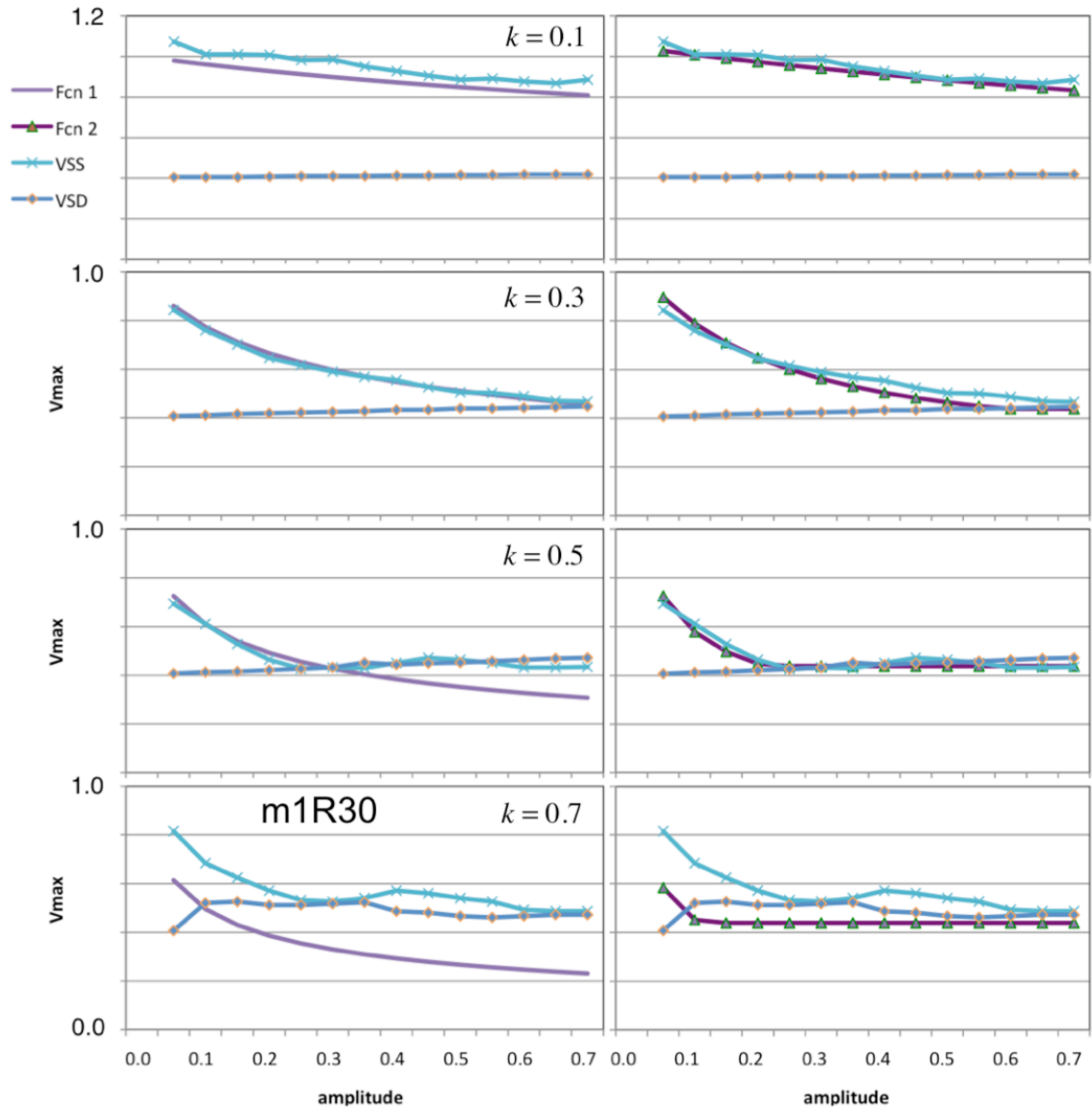


Figure C-14 Velocity Data and predictive functions

Protons H^+ at ramp width $L_r = 0.30 \lambda_{pi}$

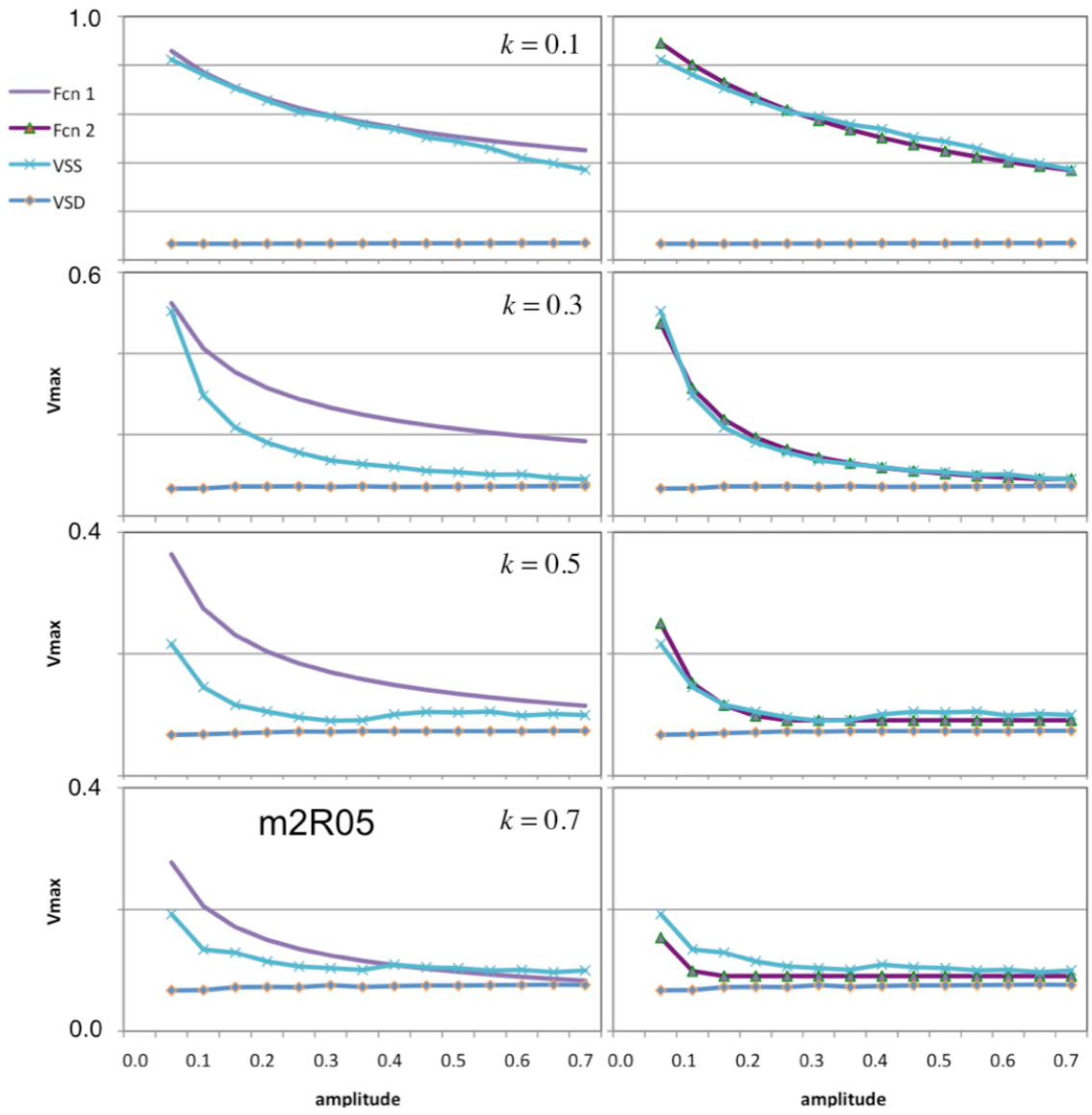


Figure C-15 Velocity Data and predictive functions

Deuterons D^+ at ramp width $L_r = 0.05 \lambda_{pi}$

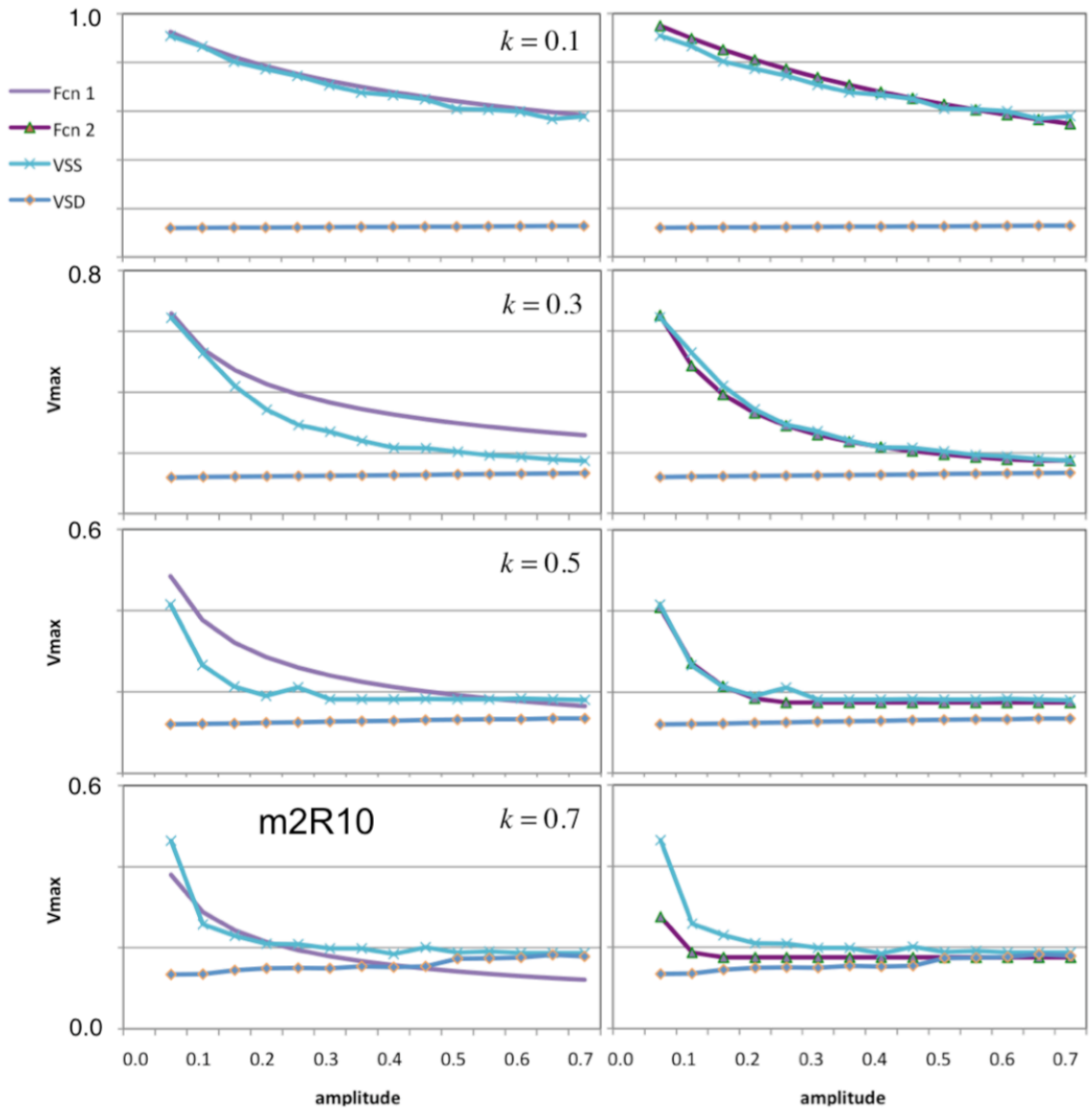


Figure C-16 Velocity Data and predictive functions

Deuterons D^+ at ramp width $L_r = 0.10 \lambda_{pi}$

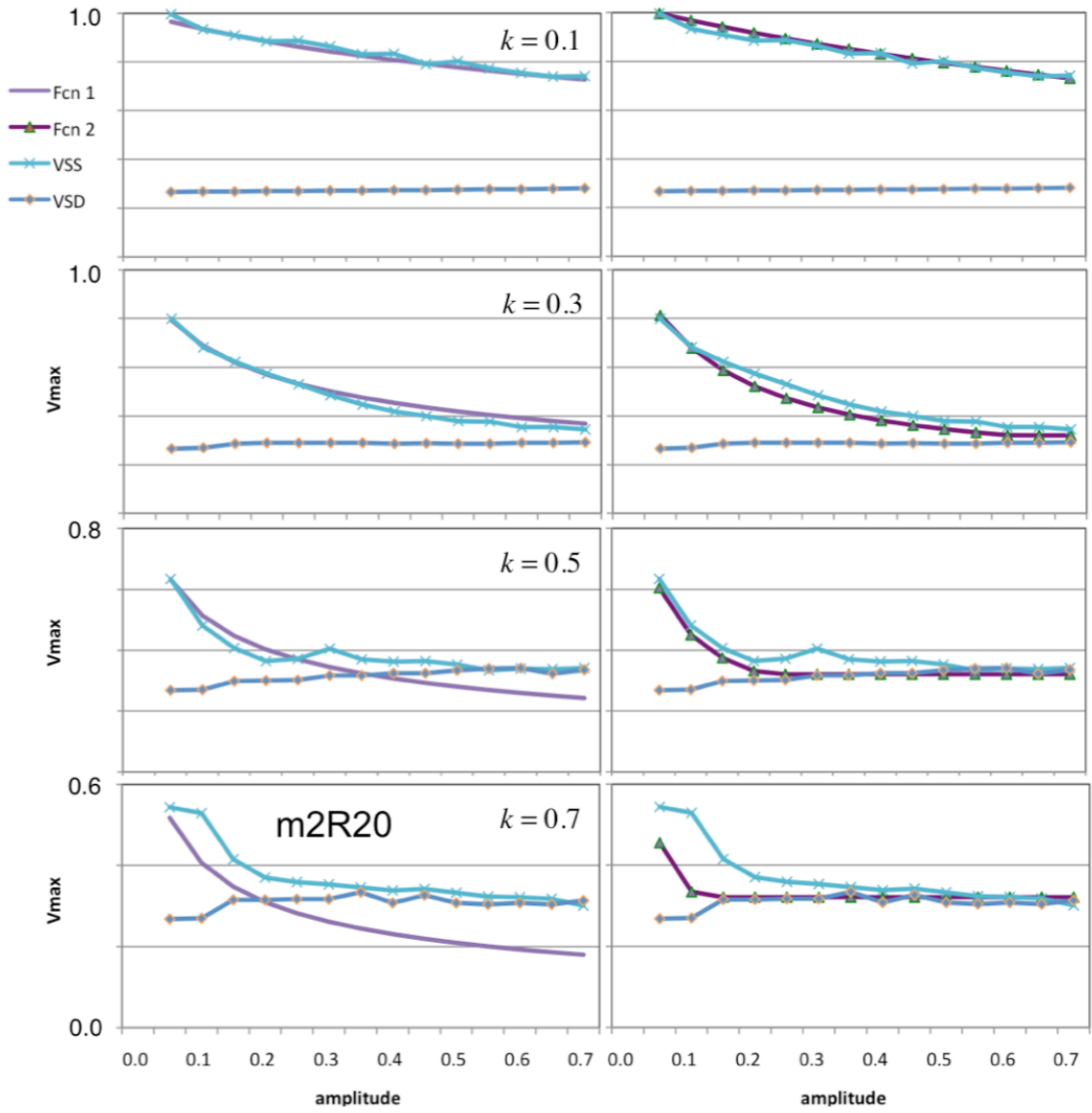


Figure C-17 Velocity Data and predictive functions

Deuterons D^+ at ramp width $L_r = 0.20 \lambda_{pi}$

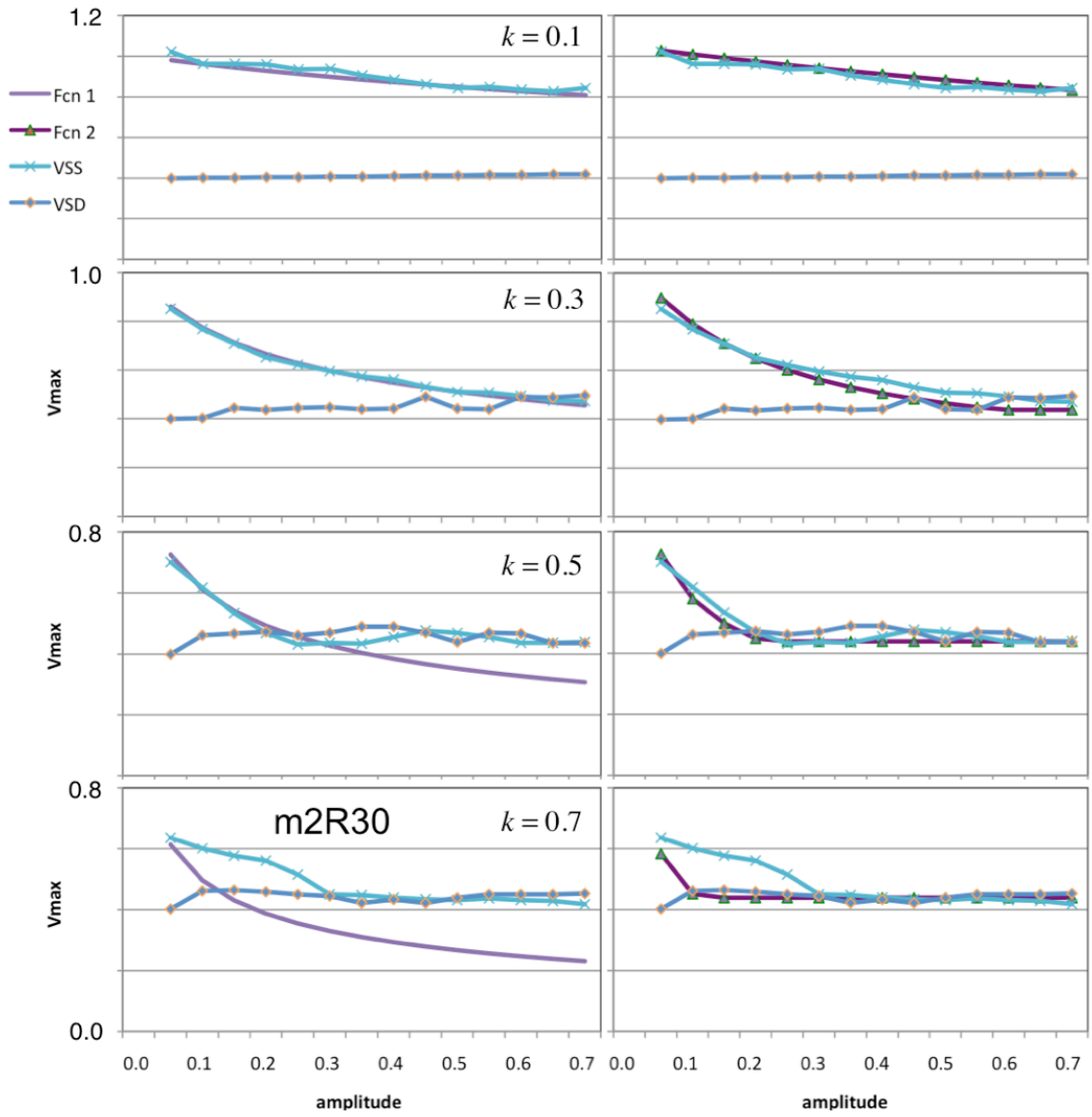


Figure C-18 Velocity Data and predictive functions

Deuterons D^+ at ramp width $L_r = 0.30 \lambda_{pi}$

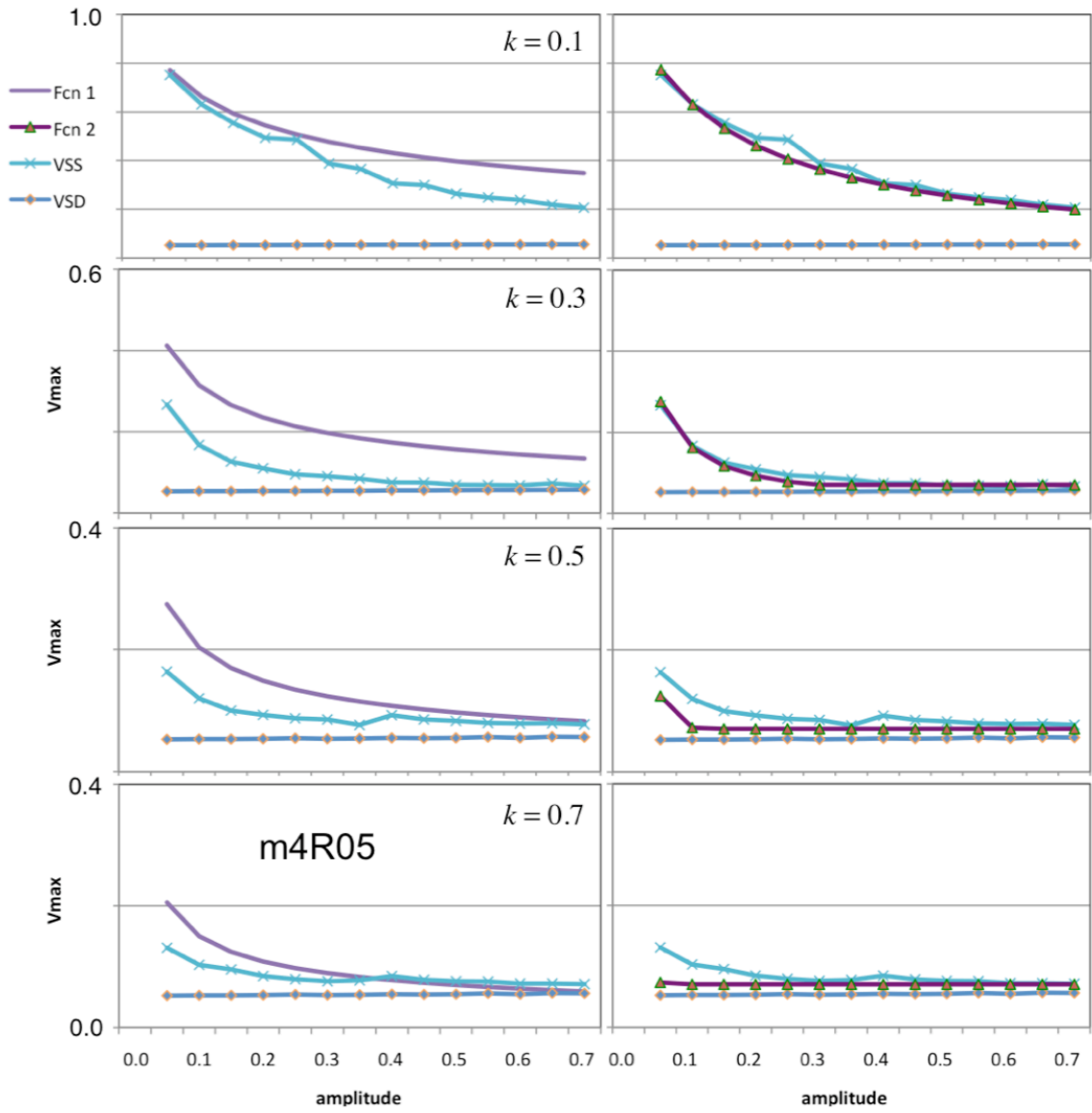


Figure C-19 Velocity Data and predictive functions

Singly ionized helium He^+ at ramp width $L_r = 0.05 \lambda_{pi}$

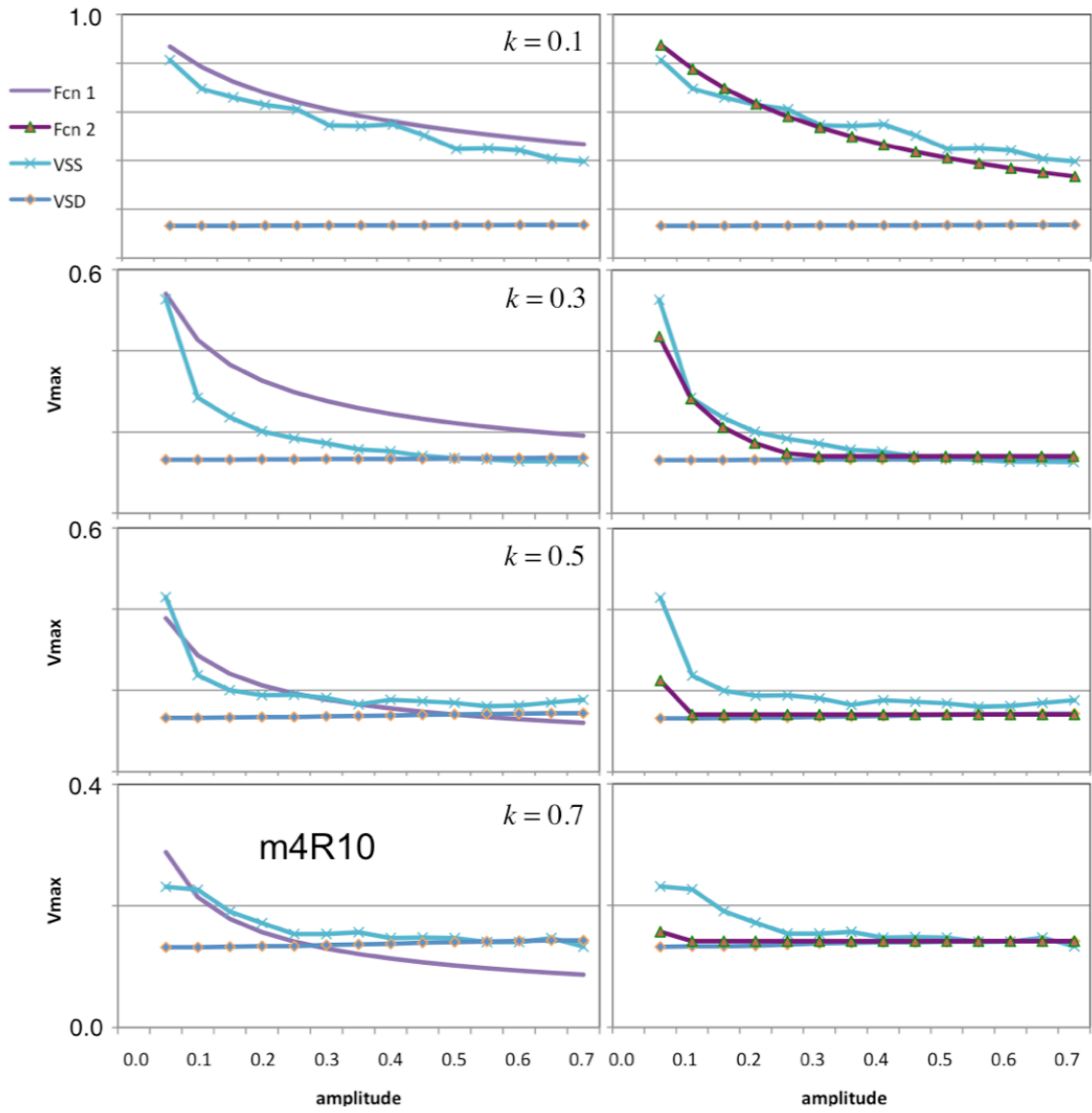


Figure C-20 Velocity Data and predictive functions

Singly ionized helium He^+ at ramp width $L_r = 0.10 \lambda_{pi}$

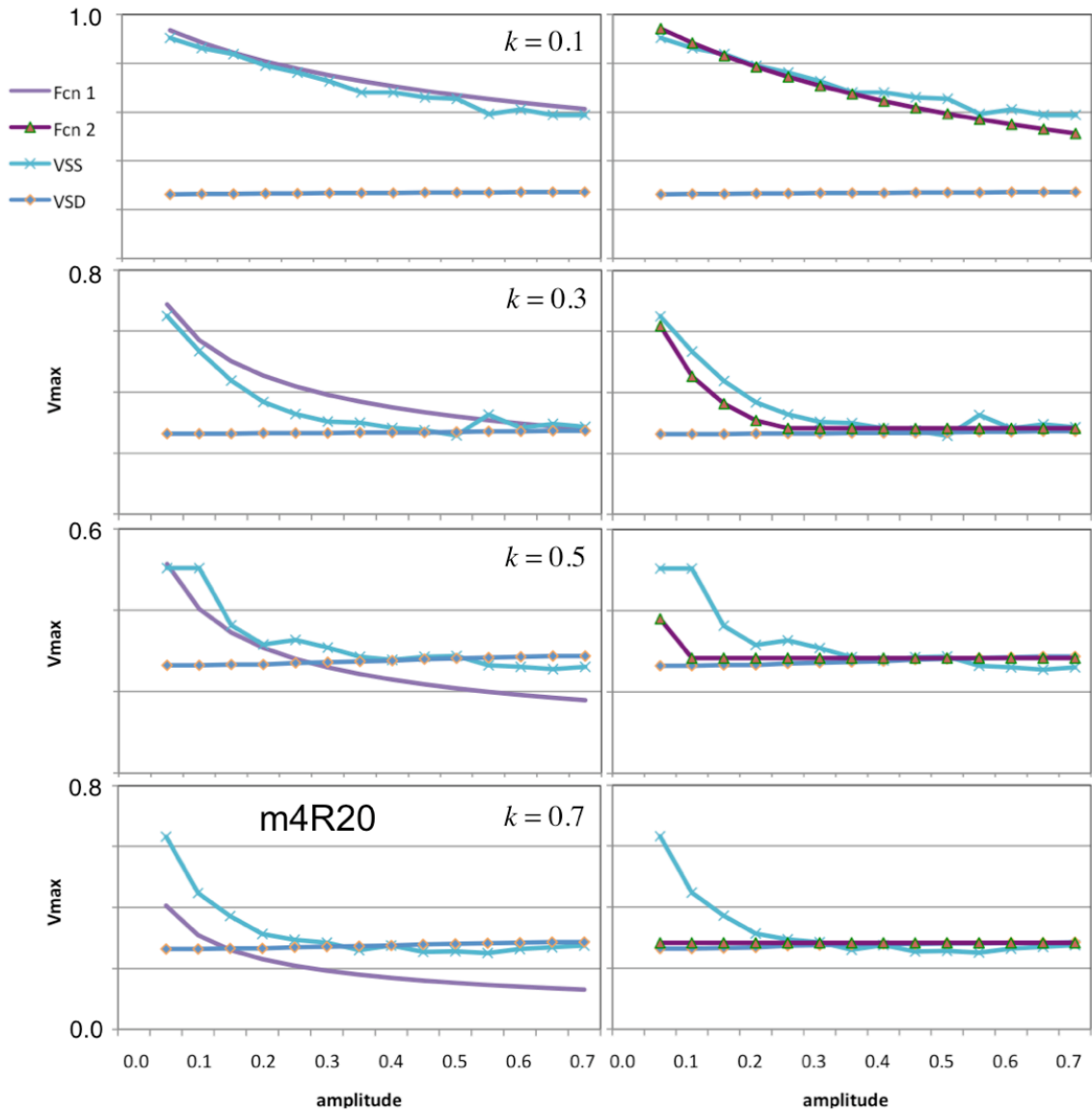


Figure C-21 Velocity Data and predictive functions

Singly ionized helium He^+ at ramp width $L_r = 0.20 \lambda_{pi}$

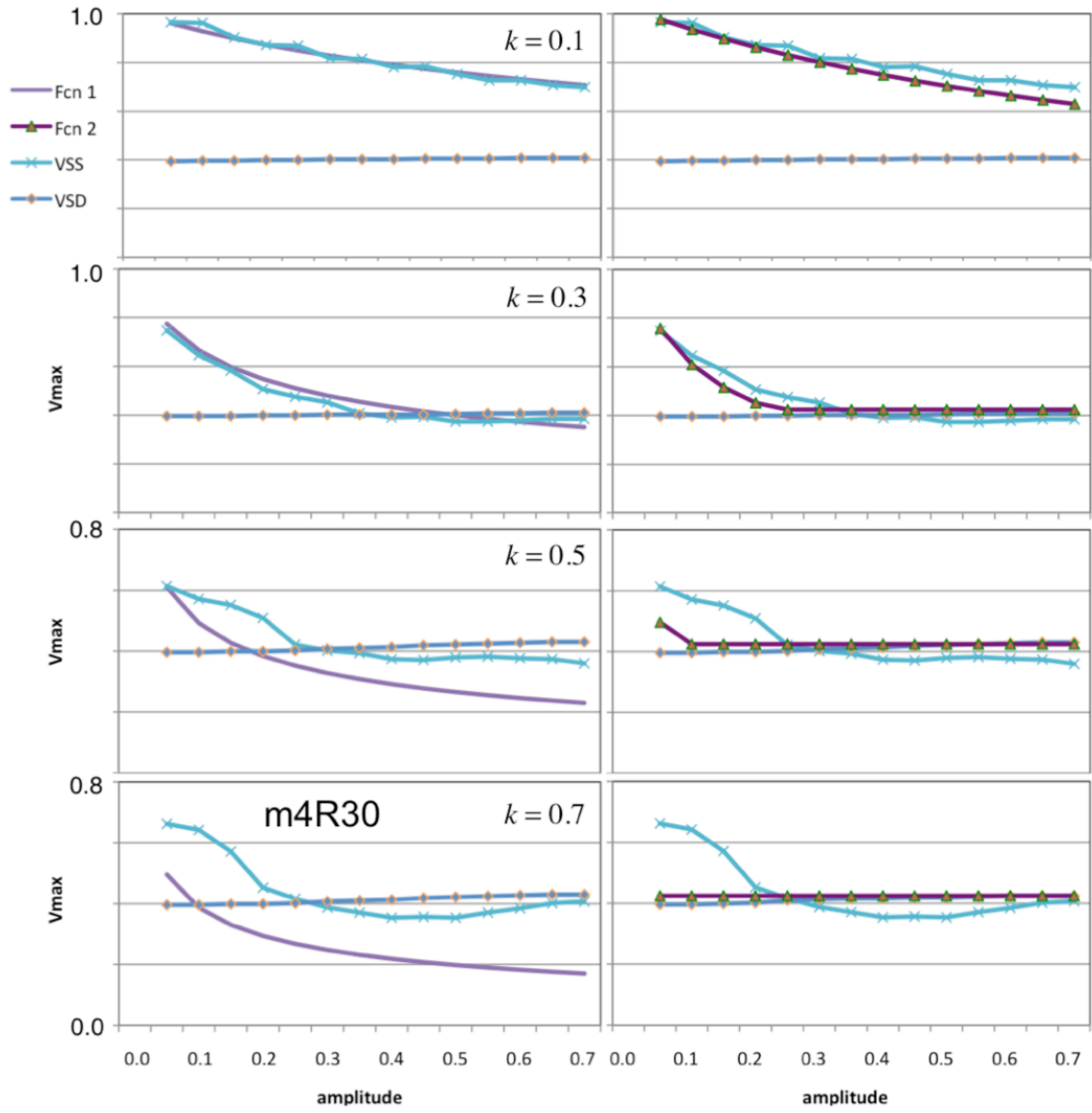


Figure C-22 Velocity Data and predictive functions

Singly ionized helium He^+ at ramp width $L_r = 0.30 \lambda_{pi}$

Appendix D Data Supplement

In this supplement we present more complete sets of data from our numerical simulations at corrugated shock fronts with detailed analysis and commentary. We begin with examples of ripple-induced velocity ceiling attenuation on shock surfing (VSS) along with data for the two shock drift mechanisms, ‘single-bounce’ (VBB) and ‘crossing’ (VCC) for comparison. We also recapitulate abbreviated examples and graphics contained in the main text our paper for easy reference. We then follow the same format for the total fluxes accelerated by each mechanism, and conclude with examples and discussion of the mechanism efficiency, defined as the flux of surfing ions promoted above the shock drift velocity ceiling.

Velocity Ceiling Data

In Figure D-1 below, we show raw data for protons of mass $m = 1$ at a Mach 5 perpendicular shock, in abbreviated parameter sets superimposed by wavenumber $k = 0.1, 0.3, 0.5, 0.7$ and grouped for comparison at four ramp widths

$L_r = 0.05, 0.1, 0.2, 0.3$. The velocities are normalized $\tilde{v} = v_{\max} \frac{B}{E_{\perp}}$, so that the maximum unperturbed velocity for each ramp width at zero amplitude is unity. These unity maxima correspond respectively to the absolute velocities $v_{\max} = 212, 106, 53, 38$ given in units normalized to the upstream Alfven speed. We also display for comparison the normalized shock drift velocity ceiling at $k = 0.5$ as tabulated at each amplitude increment in the simulations, and shown by the dotted curve labeled ‘SDN’. This not only emphasizes the scale difference between the panels, but also highlights the degree of

surfing suppression by comparison with shock drift. Note that the SD ceiling velocity is not strictly invariant, but increases slightly with perturbation strength.

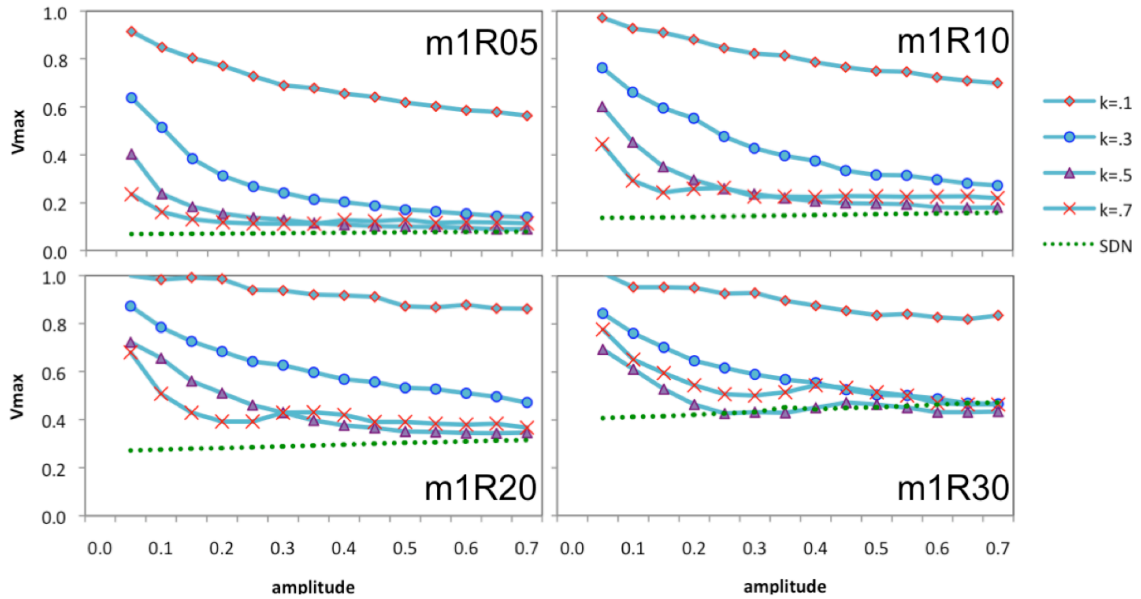


Figure D-1 Velocity Suppression of Surfing Protons

Results for a perpendicular Mach 5 shock are shown at four ramp widths indicated by the legends e.g., m1R10 means mass $m = 1$ and ramp width $L_r = 0.10$. Four wavenumber curves are superimposed in each panel, normalized to the maximum velocity attainable for the given ramp width $v_{\max} \sim L_r^{-1}$. The scale expands accordingly, as evident by the ‘SDN’ or normalized shock drift velocity, which is nearly insensitive to ramp width.

Suppression of the shock surfing velocity ceiling is immediately apparent. Even at the very thinnest ramp, the smallest wavenumber $k = 0.1$ (longest wavelength) attenuates the velocity by over 40% at the amplitude $a = 0.7$, while the shortest wavelength corresponding to $k = 0.7$ reduces the maximum surfing velocity all the way down to the shock drift ceiling at half that amplitude. Bearing in mind that the displayed velocity scale contracts with increasing ramp width, at the widest ramp shown where $L_r = 0.3$, the

$k = 0.7$ curve quickly reaches its velocity floor, while the $k = 0.5$ curve is not far above it in absolute velocity, and suppression is still substantial at the modest wavenumber $k = 0.3$. Since we expect the maximum velocity for strong perturbations to drop roughly as $v_{\max} \sim (mak^2)^{-1}$, it is not surprising that wavenumber dependence is more pronounced than that of amplitude. Note that the difference between the wavenumber curves at the widest ramp appears less drastic, but this is simply a result of the maximum unperturbed velocity being already substantially damped by the ramp.

In Figure D-2, we show sets of data for deuterons of mass $m = 2$ at the same parameters for comparison. For each wavenumber, we see that the minimum shock surfing velocity floor is reached far more quickly, that is, at weaker amplitude than that required for similar suppression of Protons. Here, the mass dependence is abundantly clear, since even the $k = 0.3$ curves reach the SD ceiling at much lower amplitudes than those for protons, which reaffirms the mass dependence in the relation $v_{\max} \sim (mak^2)^{-1}$.

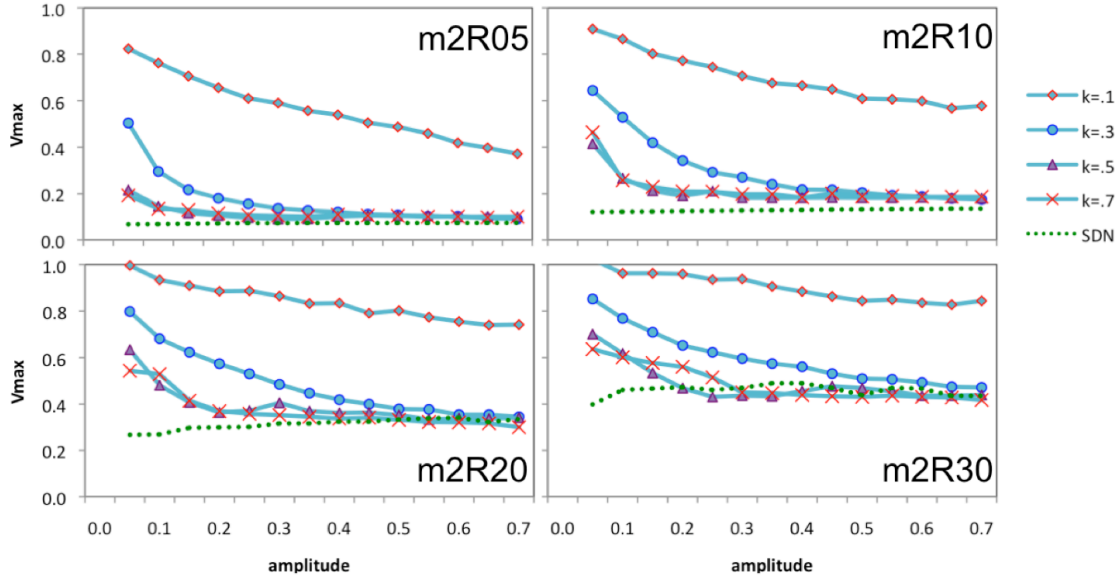


Figure D-2 Velocity Suppression for Deuterons

Panels represent the same parameters as shown in the previous figure.

The perturbation effects are more pronounced than for protons, the shortest wavelength at $k = 0.1$ suppressing the $L_r = 0.05$ surfing velocity ceiling by 60% at amplitude $a = 0.7$ with the suppression being nearly identical for wavenumbers above $k = 0.3$ at all amplitudes. At the widest ramp $L_r = 0.30$, the weakest perturbation amplitude $a = 0.05$ already drops the velocity by 40%. The ramp-dependent velocity maxima are from left to right and top to bottom are respectively, $\sim 40u_1$, $\sim 20u_1$, $\sim 10u_1$ and $\sim 6.7u_1$. Hence even the weakest ripples above $k = 0.3$ bring the maximum attainable surfing velocity respectively to $2.7v_{SD}$, $2.7v_{SD}$, $2.0v_{SD}$ and $1.5v_{SD}$ which is severe.

Our prediction regarding the velocity ‘floor’ also appears to be vindicated, that is, $v_{\max} \rightarrow v_{\text{floor}} \sim v_{SD} \sim 3u_1$. This is consistent with the minimum energy gain from a single excursion away from the front, approximately twice the piston energy. In other words,

even when the surfing mechanism fails after a single reflection, there is still additional energy gain due to the guiding center displacement outside the ramp. This is most noticeable at wide ramps where the velocity ceiling is already suppressed. Then for large perturbations, $\Delta\epsilon / \epsilon_{\max}$ due to the free flight outside the ramp is no longer small, which explains the surfing velocity ‘floor’ shown in the data. Recall that our baseline studies show that $\Delta\epsilon \sim 2\epsilon_0$ so that we should have $v_{\min} \sim 4u_1 \simeq 20$ for protons at a Mach 5 shock.

As expected, the Shock Drift “Crossing” velocity ceiling remains essentially invariant under all variation of parameters, including ramp width, perturbation amplitude and wavelength. However, for the “Single-Bounce” PI subset the velocity ceiling increases slightly with perturbation strength, which as we shall see, affects the relative surfing efficiency when this species is included in the aggregate shock drift mechanism.

We note additional features at the wider ramps, $L_r = 0.20$ and $L_r = 0.30$, that appear to be resonances induced by strong perturbations. We have not modeled processes such as normal oscillations within the ramp nor do we feel it necessary to do so. While these features are of interest and would warrant further investigation if they were plots of actual *in situ* observational data, their study will not further our analysis of the extent to which ripples suppress the shock surfing velocity ceiling.

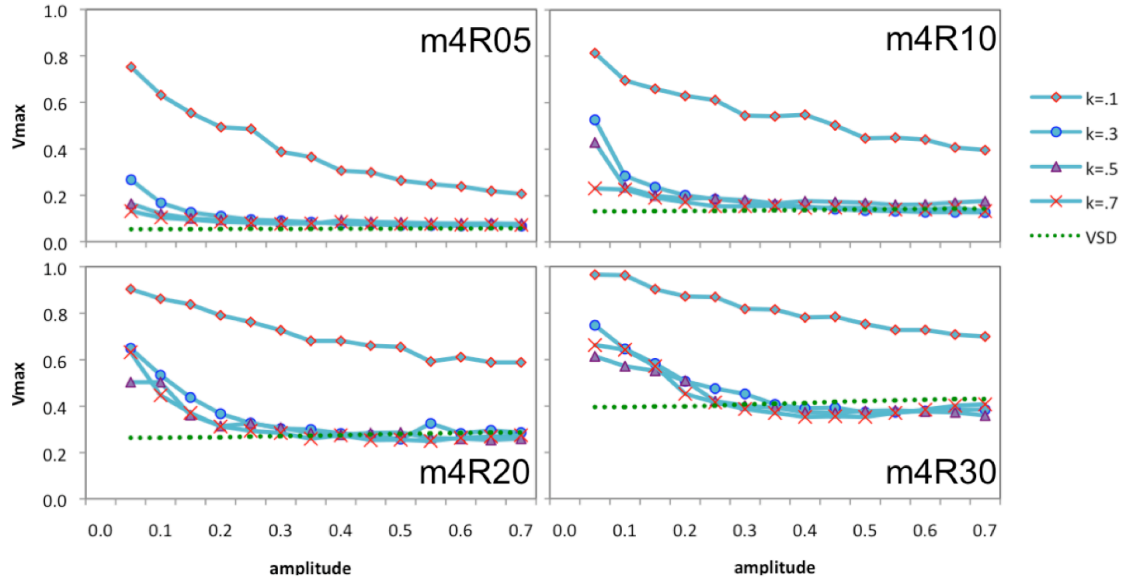


Figure D-3 Velocity Suppression Singly Ionized Helium He^+

Here the suppression is severe even for long wavelengths corresponding to $k = 0.3$ at the very weakest amplitude.

Finally, we show in Figure D-3 above the corresponding set of velocity plots for singly ionized Helium. As can be seen from a comparison of these three species from H^+ to He^+ , the deleterious effect of front ripples on shock surfing becomes more and more severe for heavier ions. That is, the surfing velocity ceiling is more readily suppressed with ripple amplitude at longer wavelengths as the ion mass increases.

Data Summary

These three sets of plots give us a convincing and comprehensive overview of the disruptive effects of front corrugations upon the shock surfing mechanism. We have seen that, even at very modest amplitudes and long wavelengths, velocity suppression can be significant and, for physical parameters still reasonable at real shocks, this suppression is sufficient to render the surfing mechanism inconsequential when

compared to shock drift. Further, the disruption of the surfing mechanism is more pronounced for heavier ions, boding ill for the application of shock surfing as a DSA pre-injection mechanism to explain the observed oxygen fluxes of anomalous cosmic rays.

Integrated Fluxes

In the following three sets of figures, we show these fluxes for three different masses at selected ramp widths, in groups of four perturbation wavelengths, each plotted against the full range of perturbation amplitudes. For comparison, we superimpose the flux curves of shock surfing ions (FSS), crossing ions (FCC), and single-bounce (FBB) mechanisms. In addition, we plot the sum of the latter two, which we have shown are essentially the same mechanism and refer to as aggregate shock drift (FSD). The flux function values plotted on the ordinate represent the decimal fraction of the total incident PI population.

We show in Figure D-4 the total fluxes of protons accelerated by each of the ‘tagged’ mechanisms. The acronym m1R10 stands for mass $m = 1$ and ramp width $L_r = 0.10$, which we will use throughout this paper. We note from this first set of figures that the aggregate shock drift flux (FSD) remains invariant with perturbation strength both in amplitude and wavenumber, while. Another feature is that the ‘crossing’ flux (FCC) and ‘single-bounce’ flux (FBB) diverge with increasing amplitude, the latter decreasing. This separation is divergence dependence is nearly linear in both amplitude and wavenumber. Although the surfing flux (FSS) increases slightly at the strongest perturbation strength shown, the difference is small in absolute terms. Finally, we see that the shock drift flux is roughly an order of magnitude below that of the entire PI population, and the shock surfing flux an order of magnitude below that.

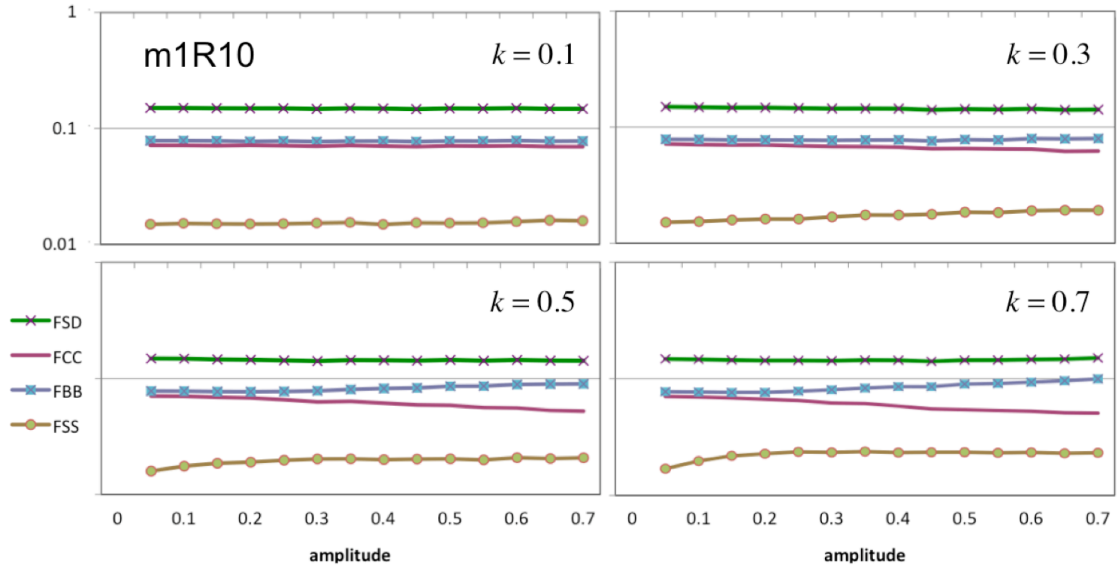


Figure D-4 Total Accelerated Proton Fluxes

The panels represent four wavenumbers as shown. The aggregate Shock Drift flux (FSD) composed of the ‘Crossing’ (FCC) and ‘single-bounce’ (FBB) components is insensitive to perturbation strength. Note that for the ramp shown, $L_r = 0.10$ FSS is an order of magnitude below FSD

In Figure D-5, we show a representative plot sample for Deuterons of mass $m = 2$ at a wider ramp $L_r = 0.20$. The aggregate shock drift flux is nearly the same, but the two components have diverged at zero amplitude, with the ‘single bounce’ having dropped by a factor of two, as has the total shock surfing flux (FSS), both as predicted from our *a priori* analyses in Chapter 2. The deviation of the two shock drift components FBB and FCC as a function of perturbation amplitude and wavenumber is still apparent but the dependence decreases with mass for FCC and increases for FBB. Further, the separation is now inverted, since the unperturbed ‘bounce’ flux now lies below the ‘crossing’.

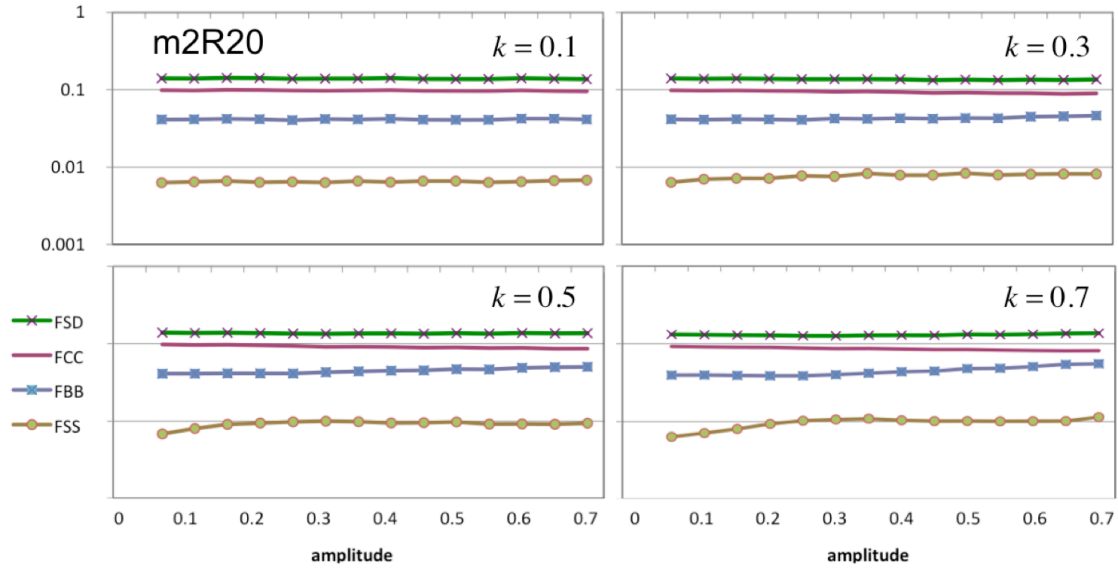


Figure D-5 Accelerated Deuteron Fluxes

For this mass and ramp width $L_r = 0.20$ the surfing flux FSS is more than an order of magnitude below the shock drift FSD

Finally, we show in Figure D-6 plots for singly ionized helium, and see that the trend in mass dependent flux characteristics has continued. Again we see a drop in the FBB and FSS fluxes by another factor of two below those of Deuterons. As was seen from the baseline studies of Chapter 3, both these mechanisms exhibit inverse mass dependence $flux \sim m^{-1}$. This follows from the fact that both these mechanisms consist of reflected ions that must satisfy the mass dependent critical incident velocity criterion $v_{x0} \leq v_{crit}$, which then determines the reflection coefficient and ultimately the total reflected flux.

In addition to the behavior of the shock drift components, we also see a pronounced increase in the total shock surfing flux with increasing perturbation strength. However, these plots merely indicate the ensemble fractions of the pickup ion population

accelerated by each mechanism, and do not in any way give us information about pre-injection efficiency as determined by the distribution tails.

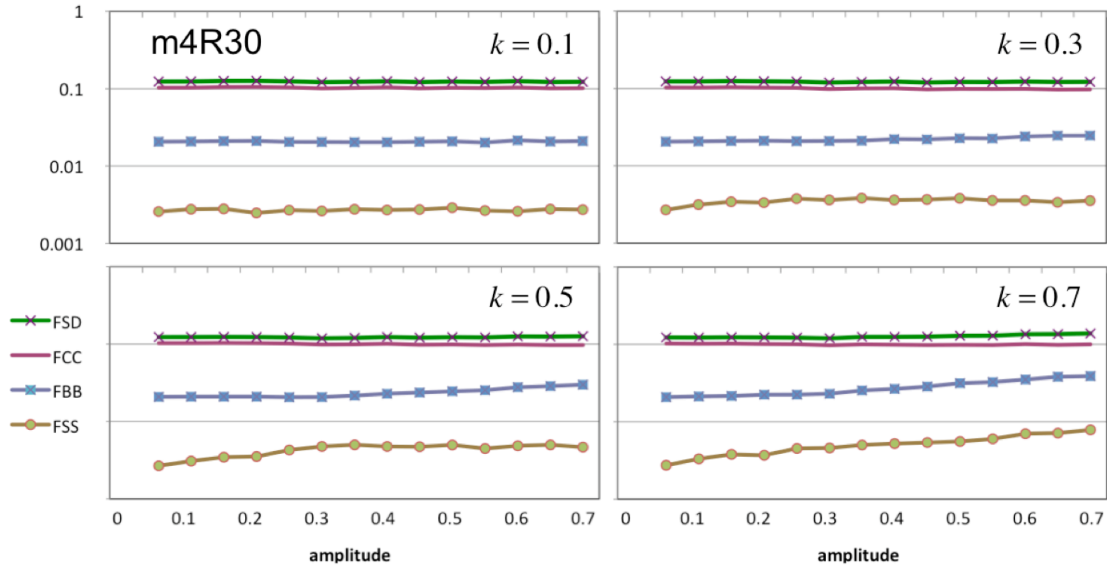


Figure D-6 Accelerated He+ Fluxes

The shock ramp illustrated is $L_r = 0.30$, wider than the sharpest profiles required for acceleration to high velocities. Note the increase in surfing flux FSS with increasing amplitude at the largest wavenumber $k = 0.7$

All of these features may alternatively be seen at a glance from the following three sample plot ensembles which group the data by mechanism for mass dependence comparison. Each figure represents a different mechanism, but with identical parameters in corresponding panels. Figure D-7 shows fluxes for the ‘crossing’ (FCC) component of the shock drift mechanism and Figure D-8 the ‘single bounce’ (FBB). In Figure D-9 we show the ‘surfing’ ion ensembles. The abscissa indicates the perturbation amplitude in reduced units $\lambda_{pi} = c / \omega_{pi}$ and the ordinate the decimal fraction of the total incident pickup ion population.

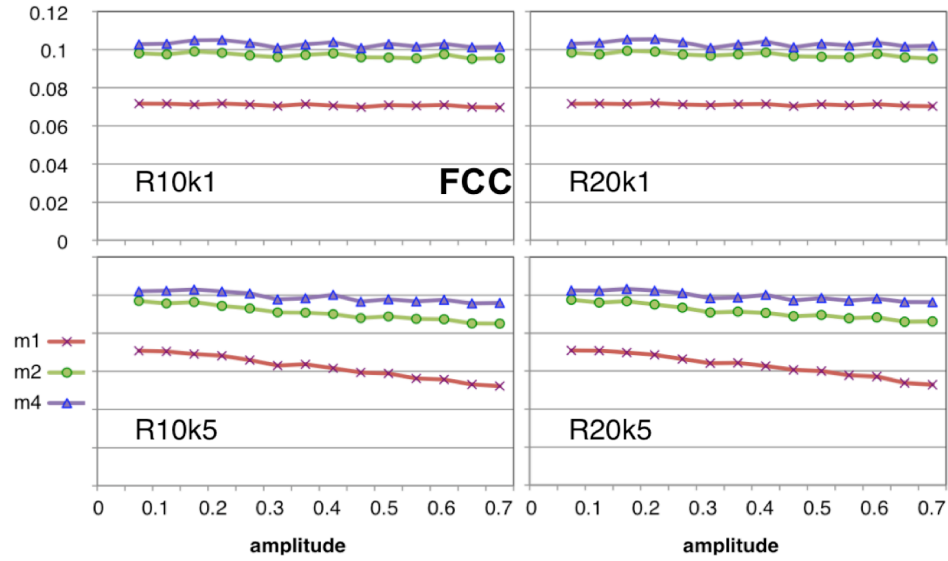


Figure D-7 'Crossing' Flux Mass Comparison

Two ramp widths $L_r = 0.10$ (left panels) and $L_r = 0.20$ (right panels) are shown. Weak perturbation wavenumber $k = 0.1$ are shown in the top panels and strong $k = 0.5$ below. All panels have identical vertical scales.

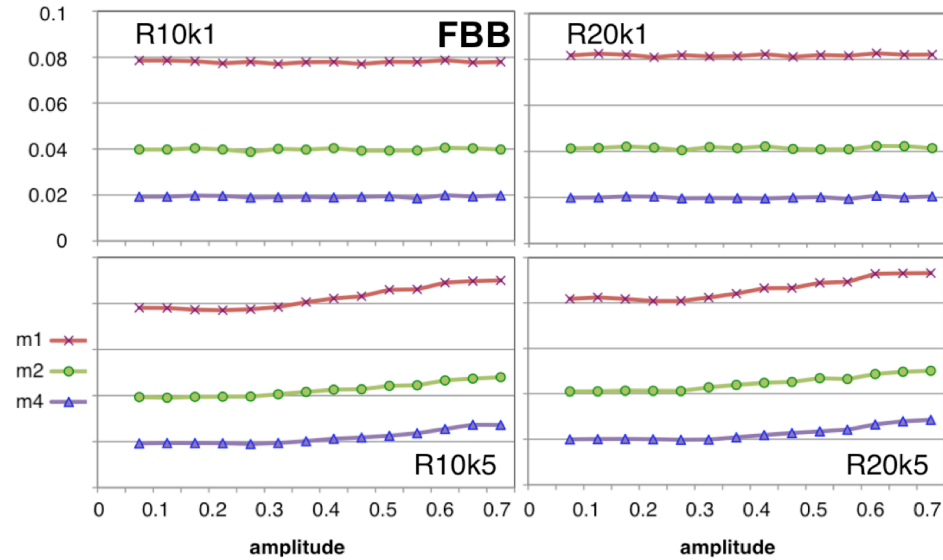


Figure D-8 'Single Bounce' Flux Mass Comparison

Using the same samples from parameter space, the perturbed FBB fluxes remain inversely proportional to mass and vary only marginally with increasing perturbation strength and ramp width.

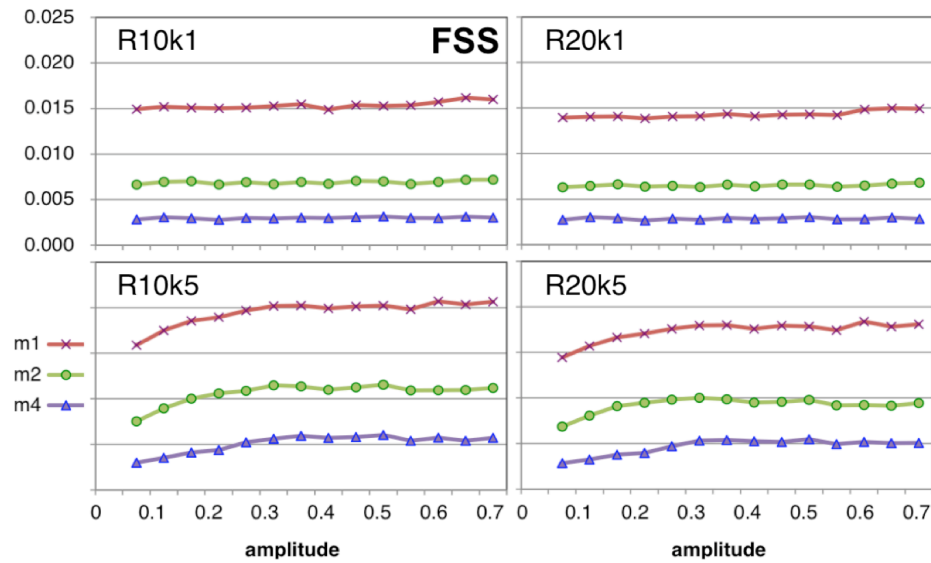


Figure D-9 'Shock Surfing' Flux Mass Comparison

As with the 'single bounce' species, fluxes are inversely proportional to mass, consistent with their dependence in the reflection coefficient.

Surfing Efficiency Section

We now offer a set of representative ensemble figures, one for each efficiency measure, using the same scheme as we did for the total fluxes above. These illustrate at a glance some of the parametric dependence of the flux tail attenuation. We begin with the surfing flux of ions accelerated above three times the upstream solar wind velocity (F3V) in Figure D-10 followed by the flux above the 'crossing' ceiling ($S > C$) in Figure D-11 and finish in Figure D-12 with the flux above the 'single bounce' ceiling ($S > B$). To be more precise, the efficiency measure may be a function of ion mass, ramp width, ripple amplitude and wavenumber.

For example, we may define $F3V(L_r, m_i, a, k) = \int_{3u_i}^{\infty} f(v_{ss}) dv$ where $f(v_{ss})$ is

the surfing distribution density function. In each of the following three figures, we superimpose curves for three different masses $m = 1, 2, 4$ representing respectively the masses of protons H^+ , deuterons D^+ and singly ionized helium He^+ .

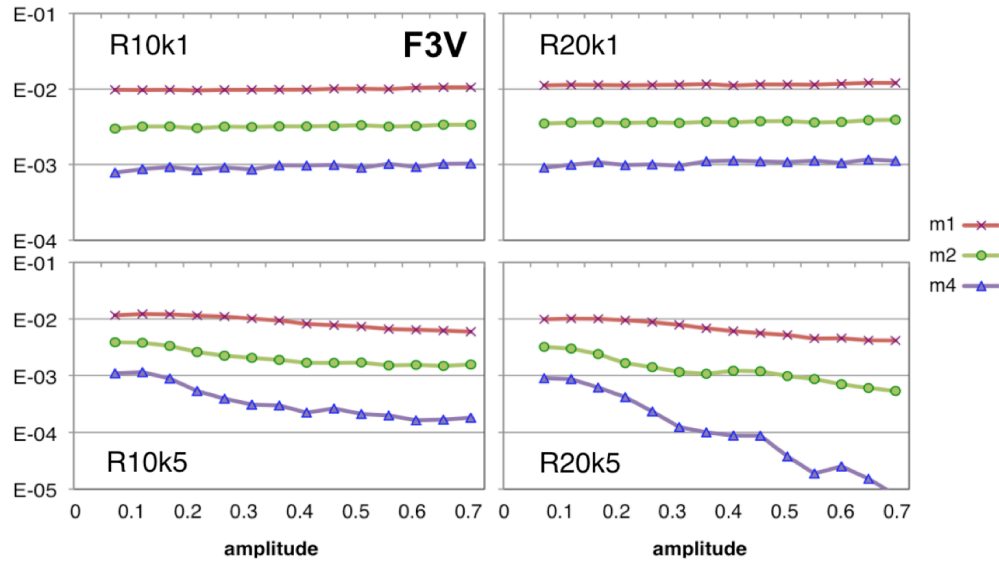
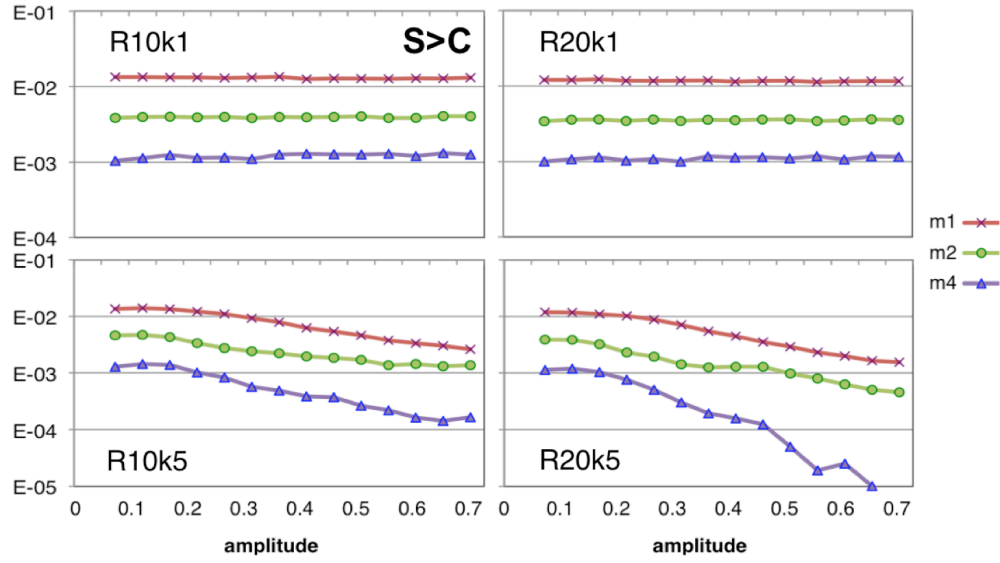
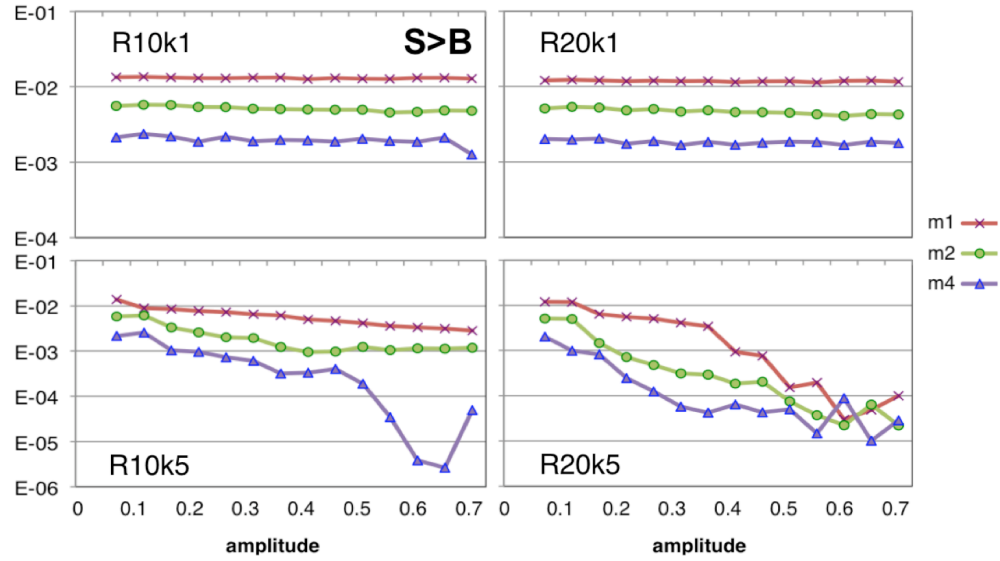


Figure D-10 Surfing Efficiency F3V

Fractions of PI promoted above $3u_{sw}$. Unperturbed values are inversely proportional to mass, with suppression at wider ramp $L_r = 0.2$ (right) compared with $L_r = 0.1$ (left) and more pronounced for He^+ ions 'm4'.

Figure D-11 Efficiency $S>C$

Fluxes of pickup ions promoted above the 'crossing' ion velocity ceiling $\hat{v}_{cc}$ are nearly identical with those of F3V, surfing fluxes above $3u_{sw}$.

Figure D-12 Efficiency $S>B$

Fluxes promoted above the single-bounce velocity ceiling $\hat{v}_{BB}$ show similar mass dependence as other measures, but are more profoundly suppressed at high perturbation strength. Singly-reflected ions attain marginally higher velocity as they are scattered by front corrugations.

Some general observations can be made from these initial examples. The flux tail suppression by all three measures is minimal at the wavelength $k = 0.1$ although, as we have already seen both analytically and computationally, the total unperturbed surfing flux declines with increasing ramp width. Note that, at the Earth's bow shock, this is equivalent to a wavelength $63\lambda_{pi} \sim 6300km$, which is extremely gentle. The efficiency by all three measures exhibits dependence on every parameter – mass, ramp width and perturbation strength, and appears very nearly linear in amplitude.

One feature of particular note is that the single-bounce measure $S>B$ drops catastrophically for He^+ when $L_r = 0.10$, $k = 0.5$ and $a = 0.45$. We assert this because the velocity ceilings v_{CC} and v_{BB} are quite close, and so drastic a flux difference shows that the shock surfing distribution function is waning quite rapidly at that point. Indeed, even if one takes the propriety of the single-bounce data with some reservation, the curve cannot be overlooked since it represents a marker only slightly higher than the others. We emphasize that, even though the velocity ceiling of these particles is extended by strong shock front perturbations, the maximum velocity is still no more than four times the piston velocity $v_{BB} < 4u_1$.

We display the single-bounce efficiency function $S>B$ for protons below in Figure D-13. The four panels each represent a different ramp width as shown, with four overlaid plots of wavenumber from $k = 0.1$ through $k = 0.7$. As we have already demonstrated, the latter represents the lower limit of supportable corrugation wavelengths at Mach 5 and therefore the steepest curves for surfing ions to follow. Even so, such wavelengths have been observed and may indeed be common shock features. These plots show that

at modest perturbations strength the surfing tail flux drops by an order of magnitude and, at the strongest perturbation, by more than two orders.

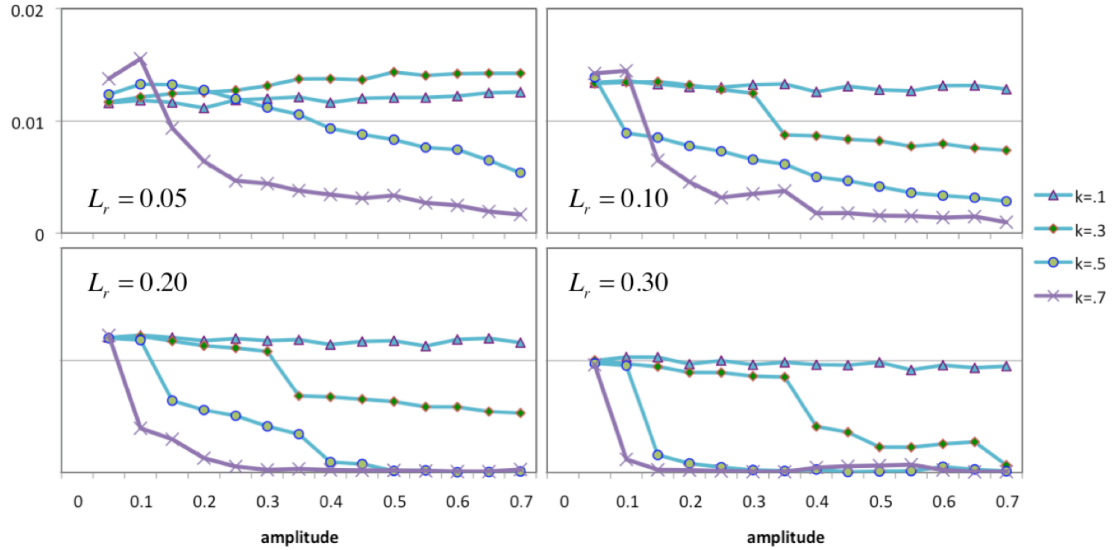


Figure D-13 Proton Surfing Efficiency $S > B$

The four panels for ramp widths shown have the same linear vertical scale and demonstrate the flux dependence on ramp width, wavenumber and amplitude.

Note that the suppression becomes stronger at wider ramps. Although we have asserted that very wide ramps are of little consequence as they severely limit the shock surfing velocity ceiling, we draw attention here to two points of interest. The first is that, should there be cases where the DSA threshold may be low enough to allow wider ramps to be admissible, thinning of the surfing tail due to corrugations in the front become of paramount importance. The second point is that shock reformation has been observed, and may perhaps be common. If this is so, the enhanced flux suppression

during shock ‘breathing’ periods where the ramp widens, necessarily restricts the surfing mechanism further in time.

While there are additional features of interest, the linear scale we have shown is of little consequence when considering that the cosmic ray spectrum spans many orders of magnitude. We now offer a more comprehensive visualization of the data, which shall make the effects of corrugation on surfing efficiency abundantly clear. As can be verified in the appendices, long wavelengths produce negligible suppression of the surfing distribution tails. We will therefore display only abbreviated sets of examples here.

In Figure D-14 below, we show the three efficiency measures $F3V$, $S>C$ and $S>B$ for protons with perturbation wavenumber applied to the shock front. The abscissa is the ripple amplitude, the ordinate in powers of ten represent the fraction of the total incident PI population, and the four panels each represent a different ramp width as shown in the acronym, e.g., ‘m1R20k5’ means mass $m = 1$, ramp width $L_r = 0.20$ and wavenumber $k = 0.5$. Here we hold the mass and wavenumber constant for ramp width comparison. For the moderate wavenumber $k = 0.5$ shown, protons are only modestly affected at narrow ramps, and the flux is reduced by less than an order of magnitude at the largest amplitude shown. The suppression becomes more substantial at wider ramps, especially as defined by the function $S>B$, as discussed above.

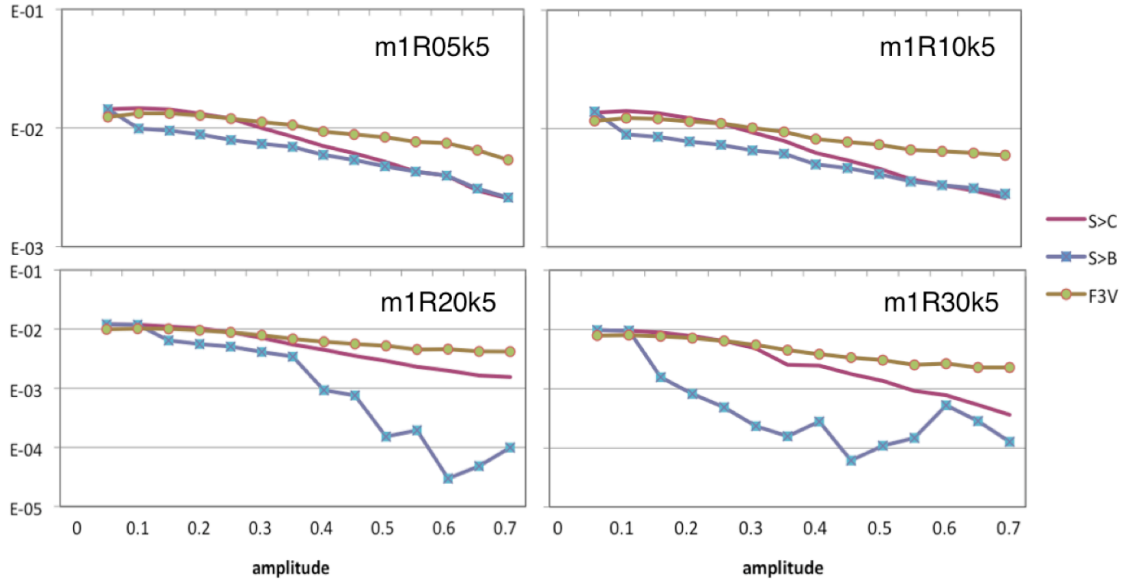


Figure D-14 Efficiency Sample for Protons

Four ramp width are shown from $L_r = 0.05$ to $L_r = 0.30$

In Figure D-15, however, we clearly see the ramp, mass, wavenumber and amplitude dependencies. Again, we stress that the dramatic appearance of the measure $S>B$ should not be entirely dismissed, since it strongly indicates the rapid drop in the surfing distribution function as it approaches $v \sim 4u_1$, only marginally greater than the crossing ceiling VCC. As evidence that the VBB velocity increase with increasing perturbation strength seen for the 'single bounce' ions is by no means extreme we offer a single representative example below in Figure D-16.

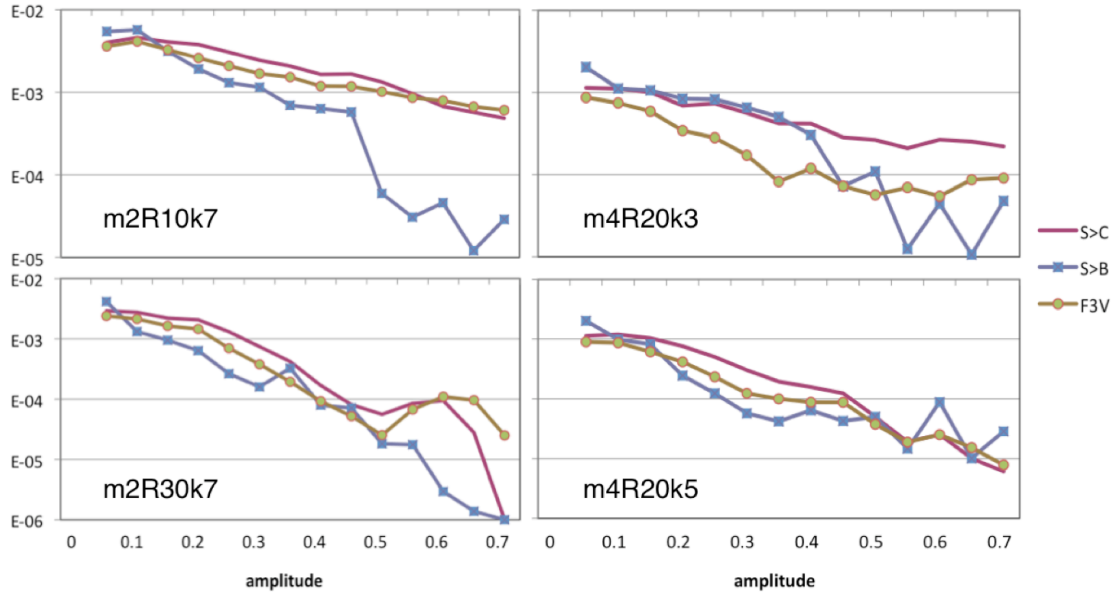


Figure D-15 Representative Efficiency Ensembles

Three measures are plotted in each of the four panels, labeled with acronyms denoting mass, ramp width and wavenumber. The upper left panel m2R10k7 means mass $m = 1$, ramp $L_r = 0.10$ and wavenumber $k = 0.7$. The extreme behavior of the function S>B may be due to paucity of data points, limitations in the model, or real features such as ripple resonant velocity enhancement.

The velocity VBB is plotted in units of the upstream solar wind speed, making clear that the deviation above $3u_1$ is small – approximately $0.5u_1$ or about 16%. But, in any case, one cannot disregard the other functions, especially F3V which is based upon a fixed marker. By this measure alone, we see attenuation by orders of magnitude just for the examples shown and greater still as seen in the further examples exhibited in the supplement at the end of this appendix.

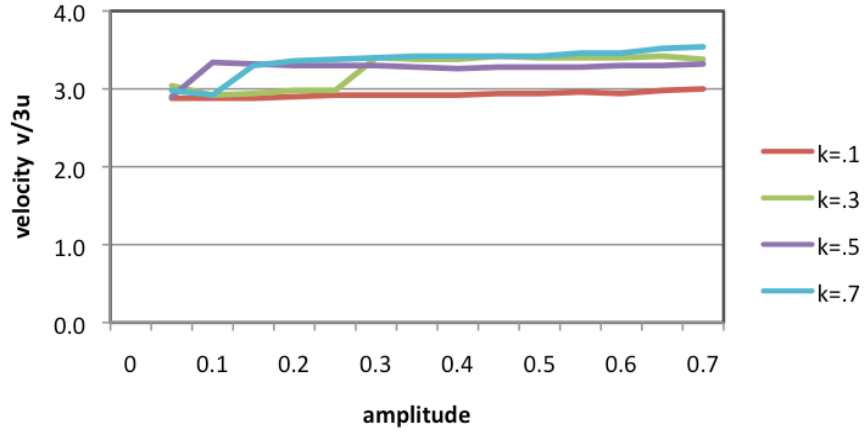


Figure D-16 Single Bounce Velocity Ceiling

Figure represents protons at Mach 5 shock of ramp width $L_r = 0.05 \lambda_{pi}$ at four corrugation wavenumbers. Velocity is given in units of the upstream wind speed.

Parameter Range Discussion

We need not extend our investigation to higher wavenumbers, since $k = 0.7$ corresponds to a corrugation wavelength $\lambda \approx 9.0$, slightly shorter than the ion cyclotron radius of the fastest pickup protons. In other words, we do not expect wavelengths substantially shorter than this to be supportable especially since the ‘gyrating’ upstream pickup ion population has often been cited in the literature as a source when discussing shock features such as the Earth’s bow shock ‘foot’. Indeed, the wavenumber ceiling constraint may be generalized by the condition $\lambda(k) \geq r_{gi}(2u_1)$, which in normalized units gives $k \leq \pi / M_A$ or $k_{\max} \approx 0.63$ for a Mach 5 shock. As we show in Chapter 6, this criterion is met precisely at the Bow Shock, as was observed *in situ* by Cluster.

Another point of note when considering a wavenumber ceiling is that both Kolmogorov turbulence theory and *in situ* observations by the Cluster Mission (where the

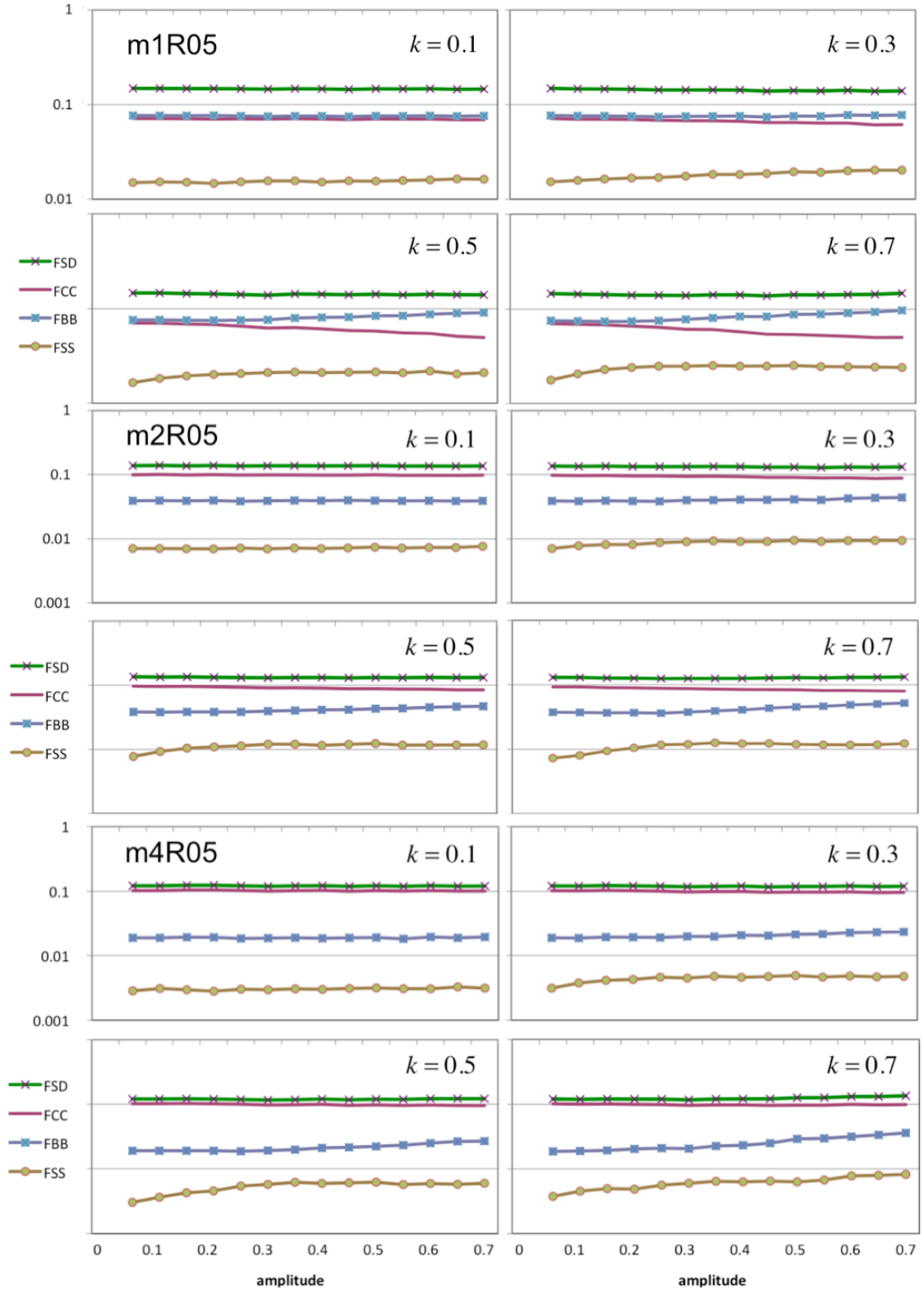
turbulence is far from fully-developed) show that, due to cascading, mechanical energy is primarily concentrated at longer wavelengths. Assuming that some form of coupling mechanism is responsible for front corrugations, we expect the amplitude at higher wavenumbers to be fairly weak in any case, even if no resonant physical limiting mechanism were to exist. We can easily quantify this, and do so in depth in the next Chapter. The spectral index for magnetic turbulence at the Bow Shock observed by Cluster is approximately an inverse square, more specifically $\varepsilon_B \sim k^{-1.9}$. Fully developed Kolmogorov turbulence is nearly the same at $\varepsilon_B \sim k^{-5/3}$. Coupling by equipartition between turbulent magnetic energy and kinetic energy of oscillations along the shock front would lead to the dependence $\varepsilon_K \sim \varepsilon_B \Rightarrow \frac{1}{2}\rho\Omega^2 a^2 \sim \frac{\delta B^2(k)}{8\pi}$ or $a^2 \sim \delta B^2$ where a is the amplitude of the corrugation. We now have $a^2 \sim k^{-\gamma}$ and are free to choose the spectral index appropriate to the problem. Since the difference is small, we choose $\gamma = 2$ for convenience so that our surfing velocity suppression factor in terms of wavenumber alone becomes $(mk)^{-1}$ making the velocity ceiling dependence inversely proportional to both mass and wavenumber.

Finally, although we have modeled a perturbed shock with quasi-static corrugations, this should not materially affect our results. We implicitly assumed dynamic perturbations, but with times scales which appear long in the surfing ion's frame. This may be understood by considering both the limiting Alfven wave speed along the front as well as the frequency regime of possible front oscillations. Once an ion's velocity has considerably exceeded the ram speed, any such time or wave velocity dependence will appear at worst quasi-stationary in the particle's frame of reference.

Total Flux Data

This section begins with plots of the total fluxes of ions accelerated by Shock Surfing (FSS), Single-Bounce (FBB), and Crossing (FCC) mechanisms, and aggregate Shock Drift (FSD) defined as the sum of FBB and FCC. The first set consists of comparative plots of all four mechanisms, in groups of four panels each with a different wavenumber. Each of the quad sets represents a single ion mass and ramp width, as denoted by the anagram label. For example, the first quad is labeled 'm1R05' which represents mass $m = 1$ and ramp width $L_r = 0.05$. Consistent with the convention adopted throughout this paper, amplitudes are given in units normalized to the ion inertial length $\lambda_{pi} = c / \omega_{pi}$, wavelength to $2\pi / \lambda_{pi}$, and flux values on the ordinate represent the portion of the pickup ion population accelerated by the particular mechanism, given as a decimal fraction.

Beginning with Figure D-17 through Figure D-20, the groups on each page all represent the same ramp width, and illustrate the comparative insensitivity to variation in parameters of the Shock Drift mechanism. The entire ensemble shown represents results from simulations of Mach 5 perpendicular shocks.

Figure D-17 Accelerated Fluxes at ramp $L_r = 0.05$ grouped by ion mass

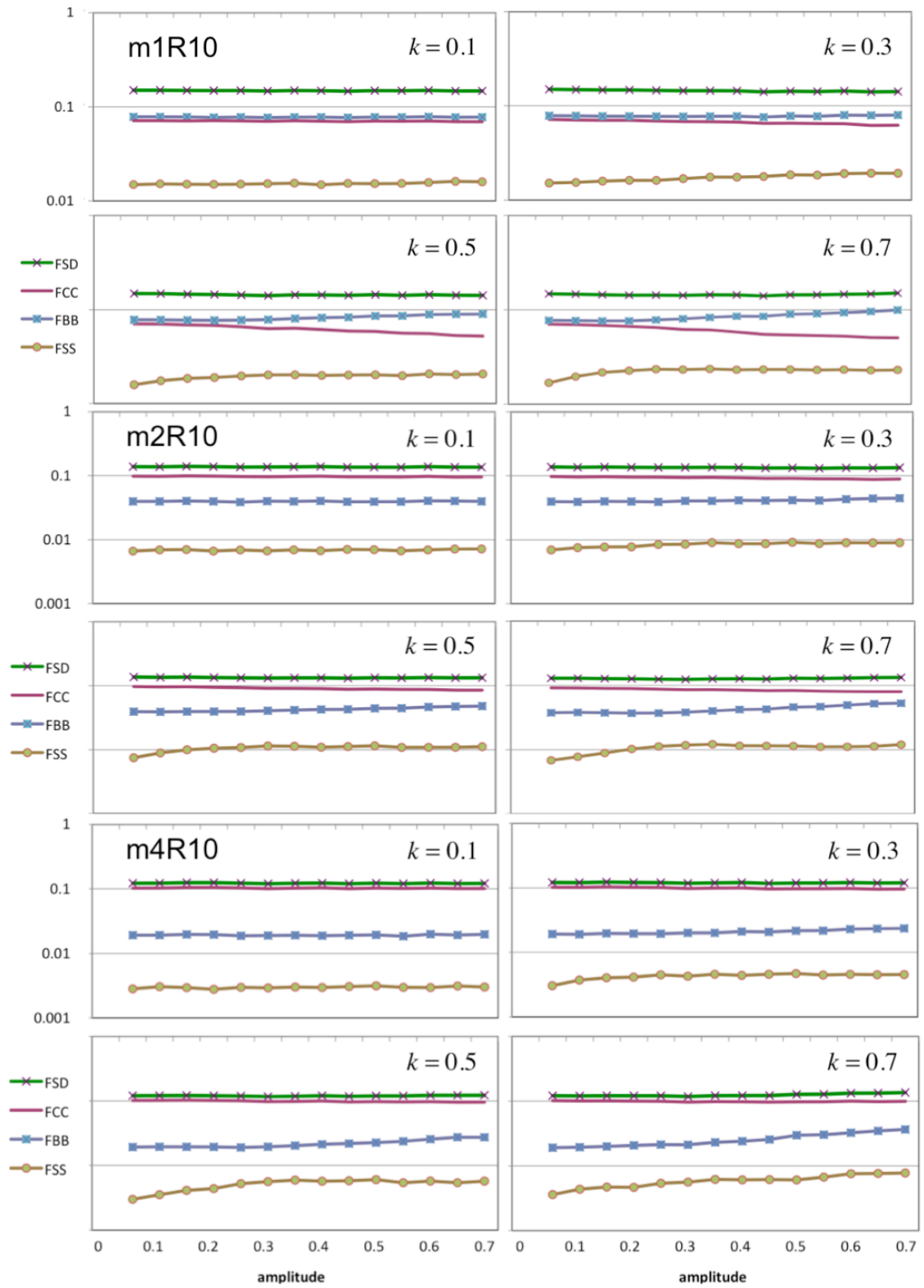


Figure D-18 Accelerated Fluxes at ramp $L_r = 0.10$ grouped by ion mass

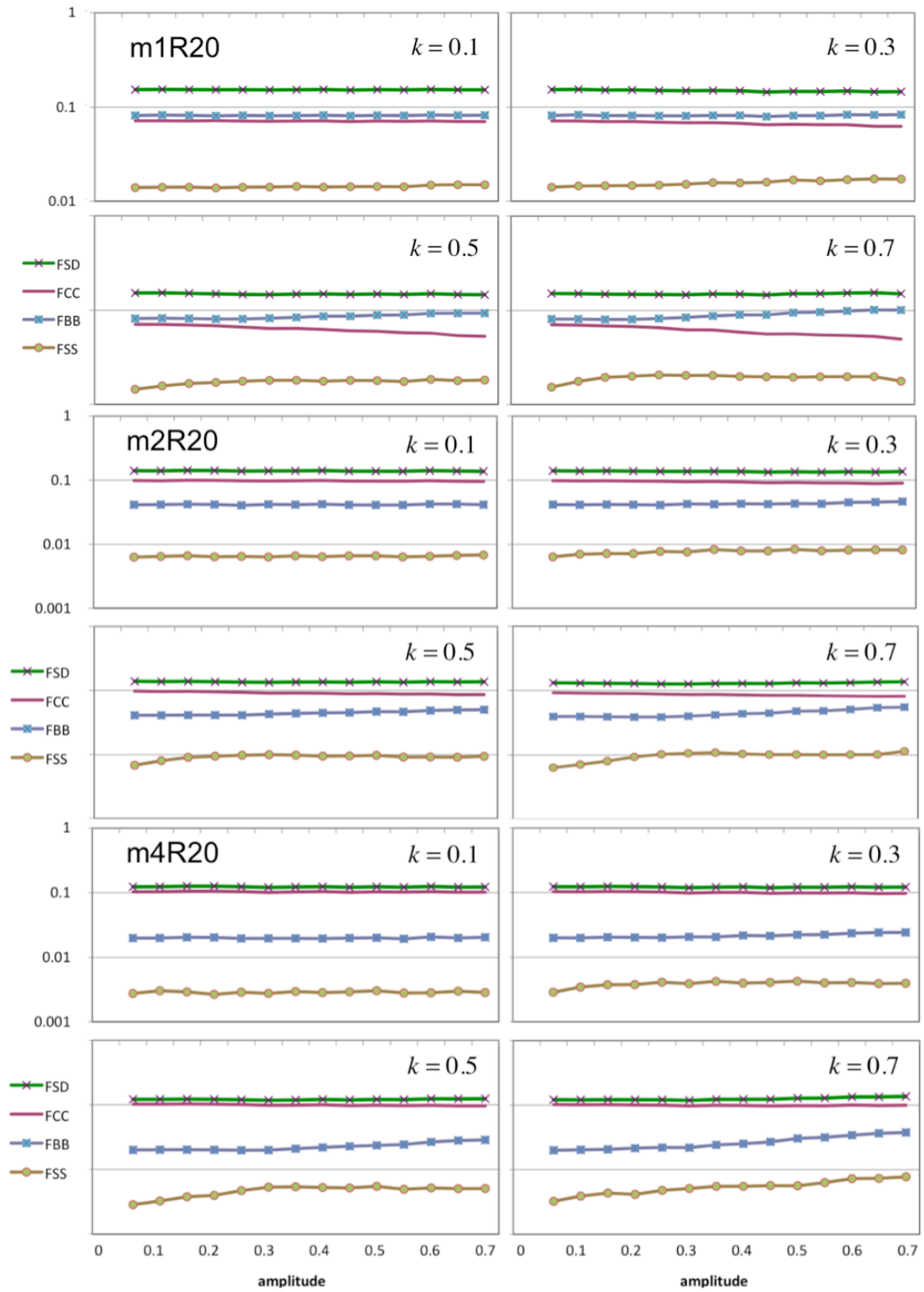
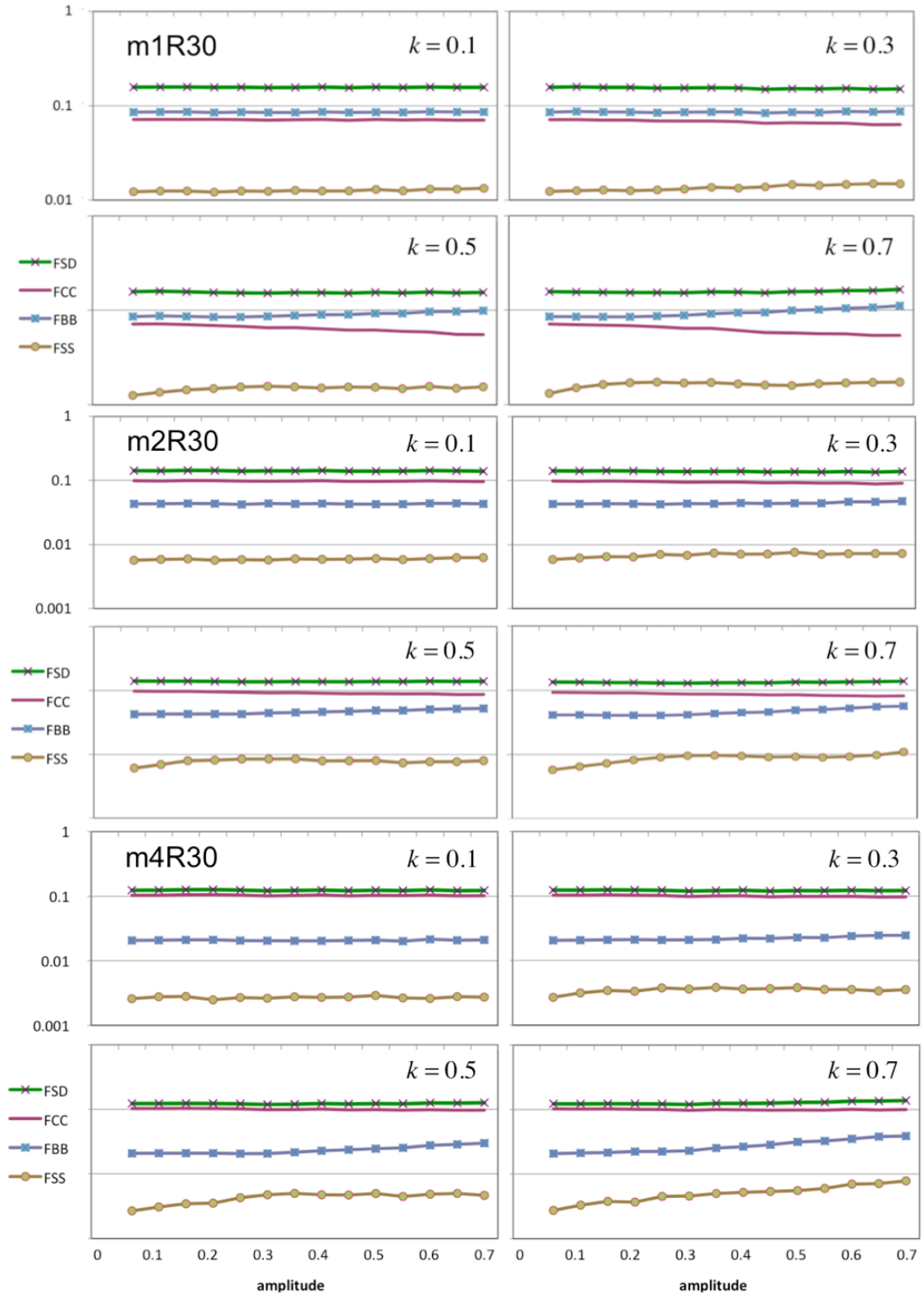


Figure D-19 Accelerated Fluxes at ramp $L_r = 0.20$ grouped by ion mass

Figure D-20 Accelerated Fluxes at ramp $L_r = 0.30$ grouped by ion mass

The following sets beginning with Figure D-21 and ending with Figure D-24 represent the same data, reorganized and presented in an alternate visualization. The data are grouped by acceleration mechanism, FBB, FCC, FSD and FSS as defined above. These are presented in groups of three quad panel sets per page. Each set represents a single ion mass, and the four superimposed curves per panel are identified by ripple wavenumber. The vertical scales are expanded to emphasize the divergence of each mechanism as a function of perturbation strength, and more clearly demonstrate the effects of ripples on each mechanism separately. However, the narrow range of values on the vertical scale indicates the mild parameter dependence of the total fluxes.

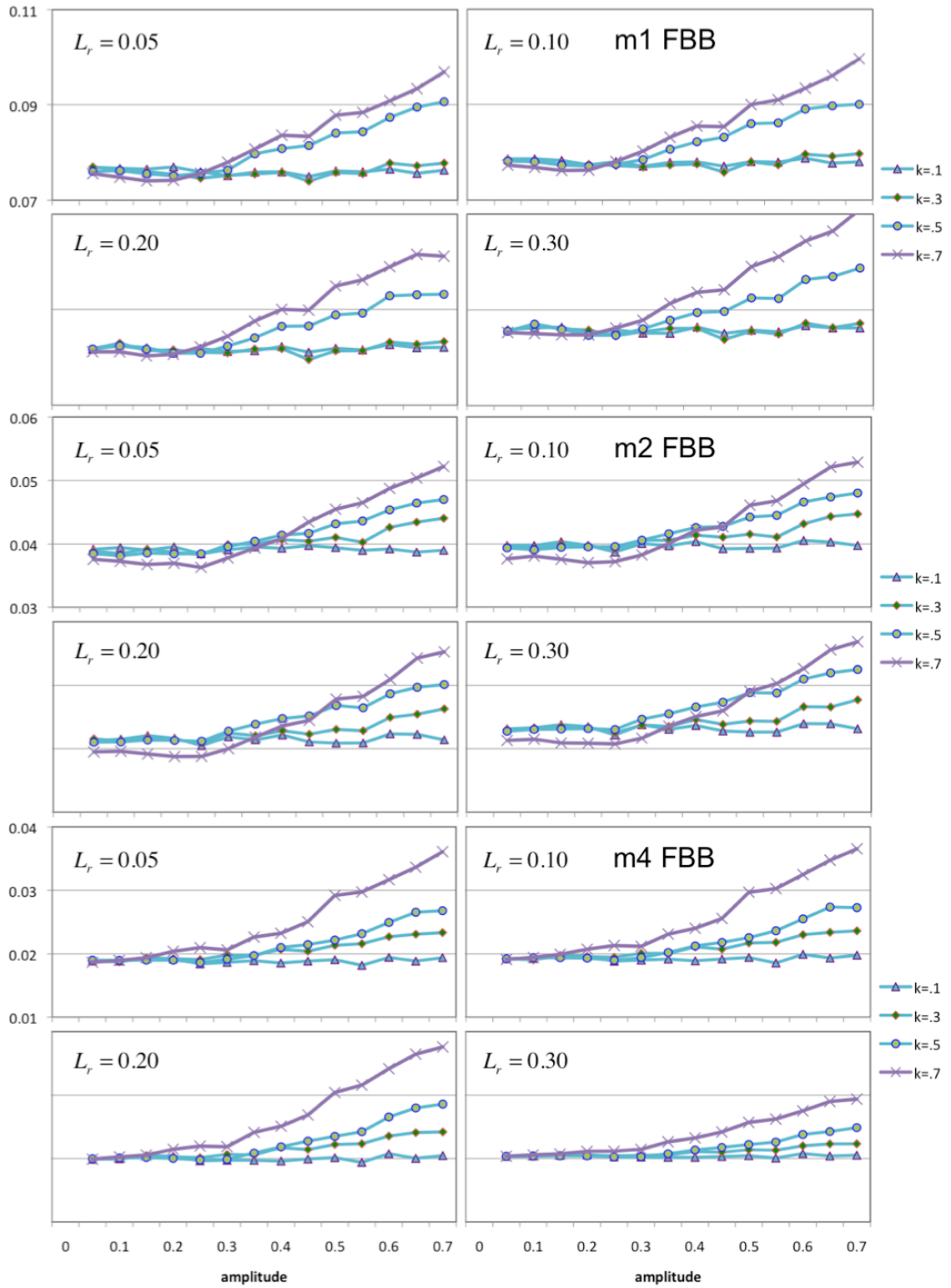


Figure D-21 Single-Bounce Fluxes grouped in quads by ion mass

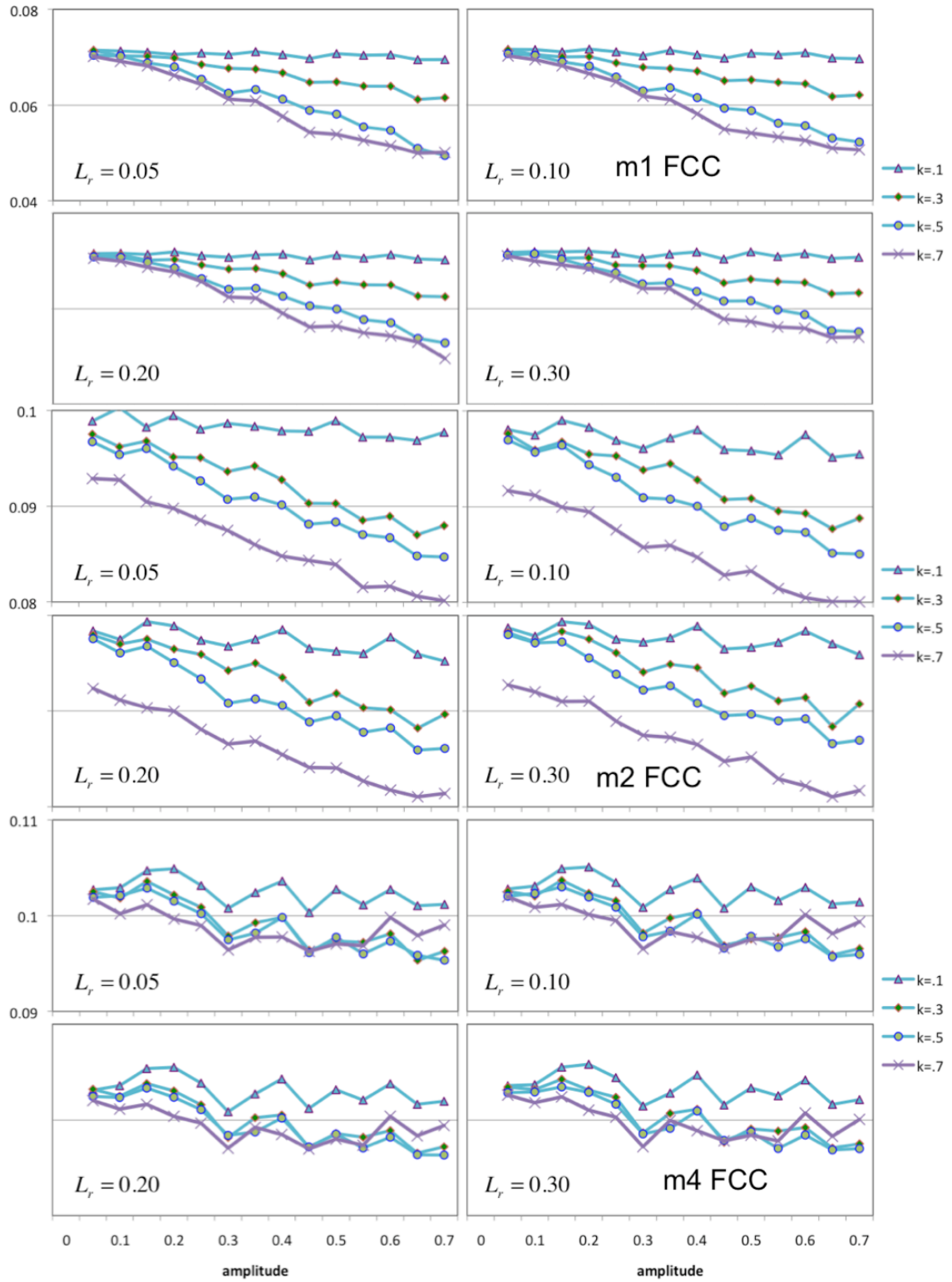


Figure D-22 Crossing Fluxes grouped in quads by ion mass



Figure D-23 Aggregate Shock Drift Fluxes grouped by ion mass

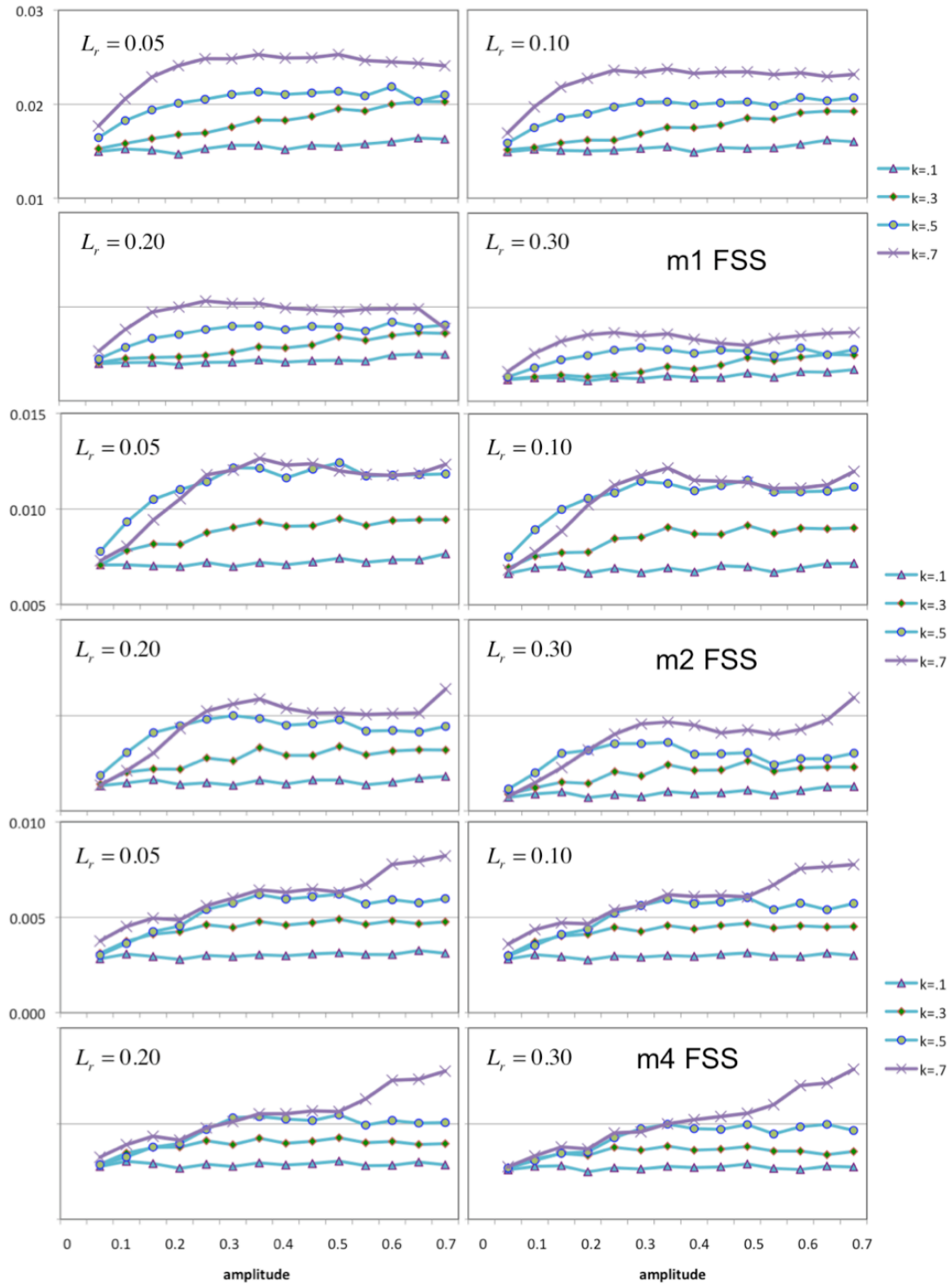


Figure D-24 Shock Surfing Fluxes grouped by ion mass

Surfing Efficiency

The following data sets show the surfing 'efficiency' as defined by the flux of surfing ions promoted above the velocity ceiling of the 'single-bounce' ($F > B$) and 'crossing' ($F > C$) mechanisms, and the function $F3V$ which measures the flux above three times the incident solar wind speed. These are given in quad sets of single mass value with each of the four panels representing a different ramp width, and curves for four perturbation wavelengths plotted together in each.

These quad panel sets beginning with Figure D-25 and ending with Figure D-27 are grouped by mechanism, one to each page, with three masses represented per page – protons, deuterons and singly ionized helium.

All scales are linear.

Following these, we again reproduce the same efficiency data set in a different visualization scheme, with log plots to emphasize the flux suppression of the surfing ion flux tails. Beginning with Figure D-28 and continuing through Figure D-31, we superimpose curves for each efficiency measure in three groups of quad panels per page. Each page represents the same corrugation wavenumber, and the four panels in each group are different ramp widths, while each of the three groups represent a different ion mass.

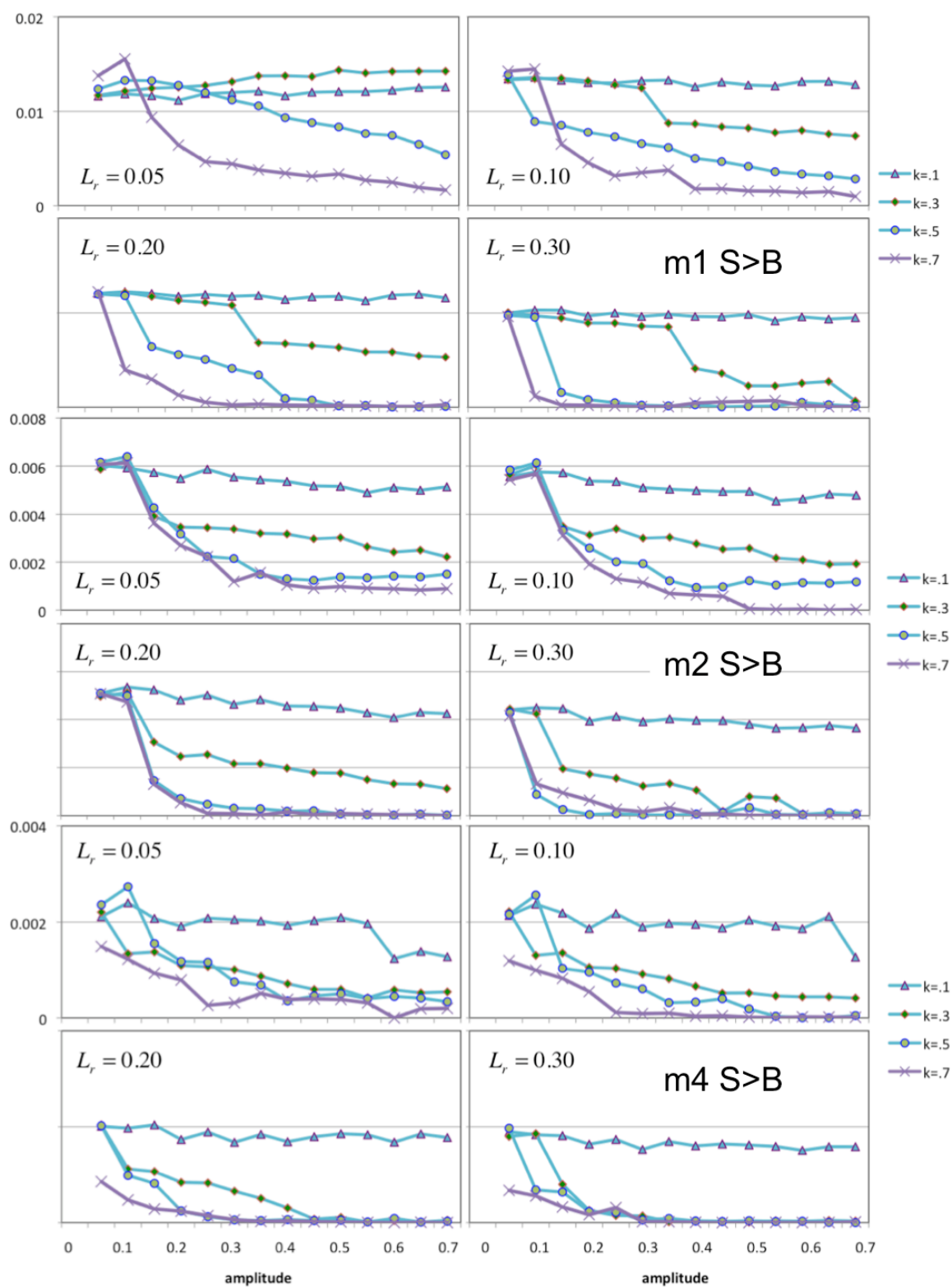


Figure D-25 Surfing Efficiency above the Since Bounce Ceiling

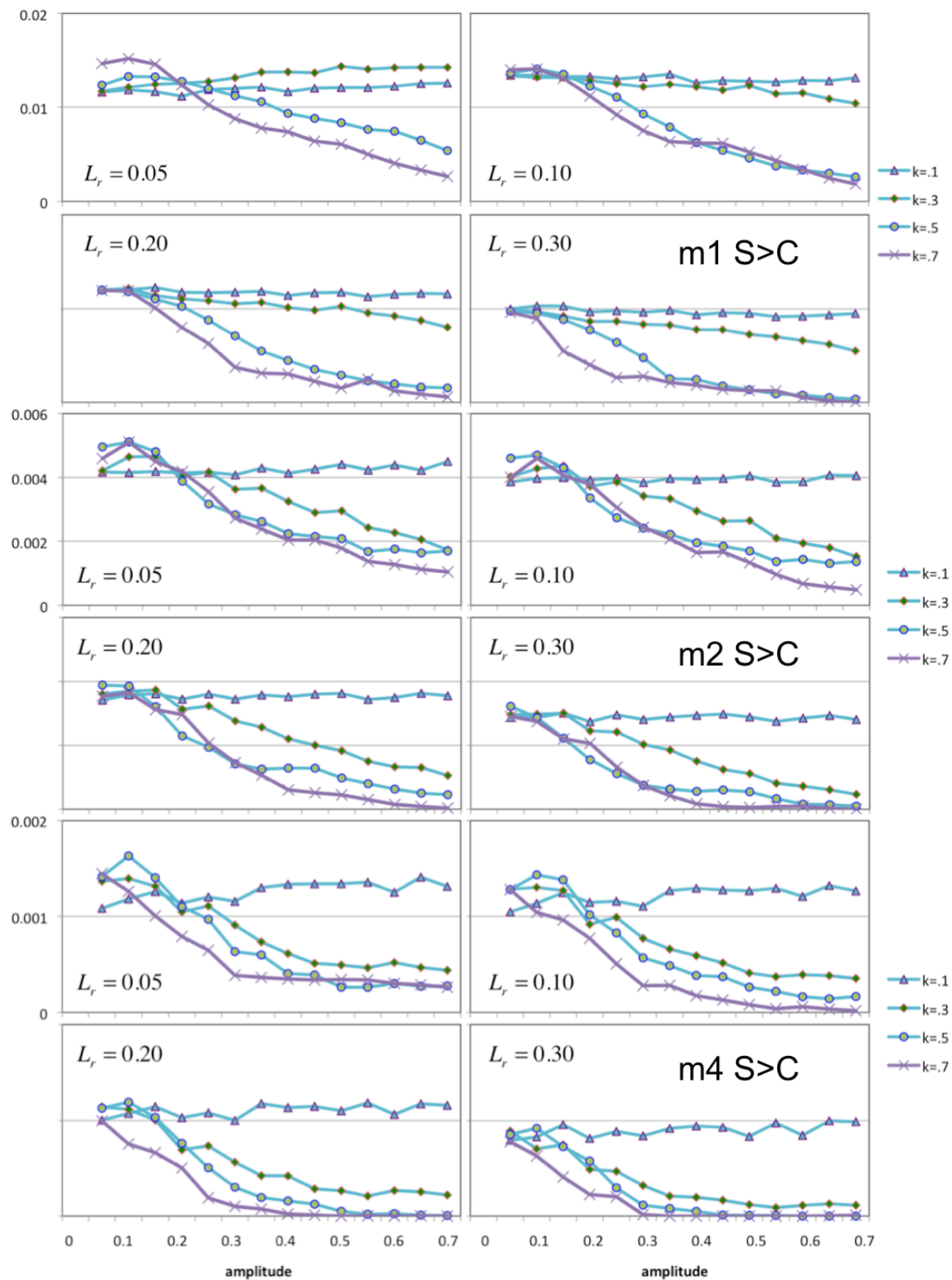


Figure D-26 Surfing Efficiency above the Crossing Ceiling

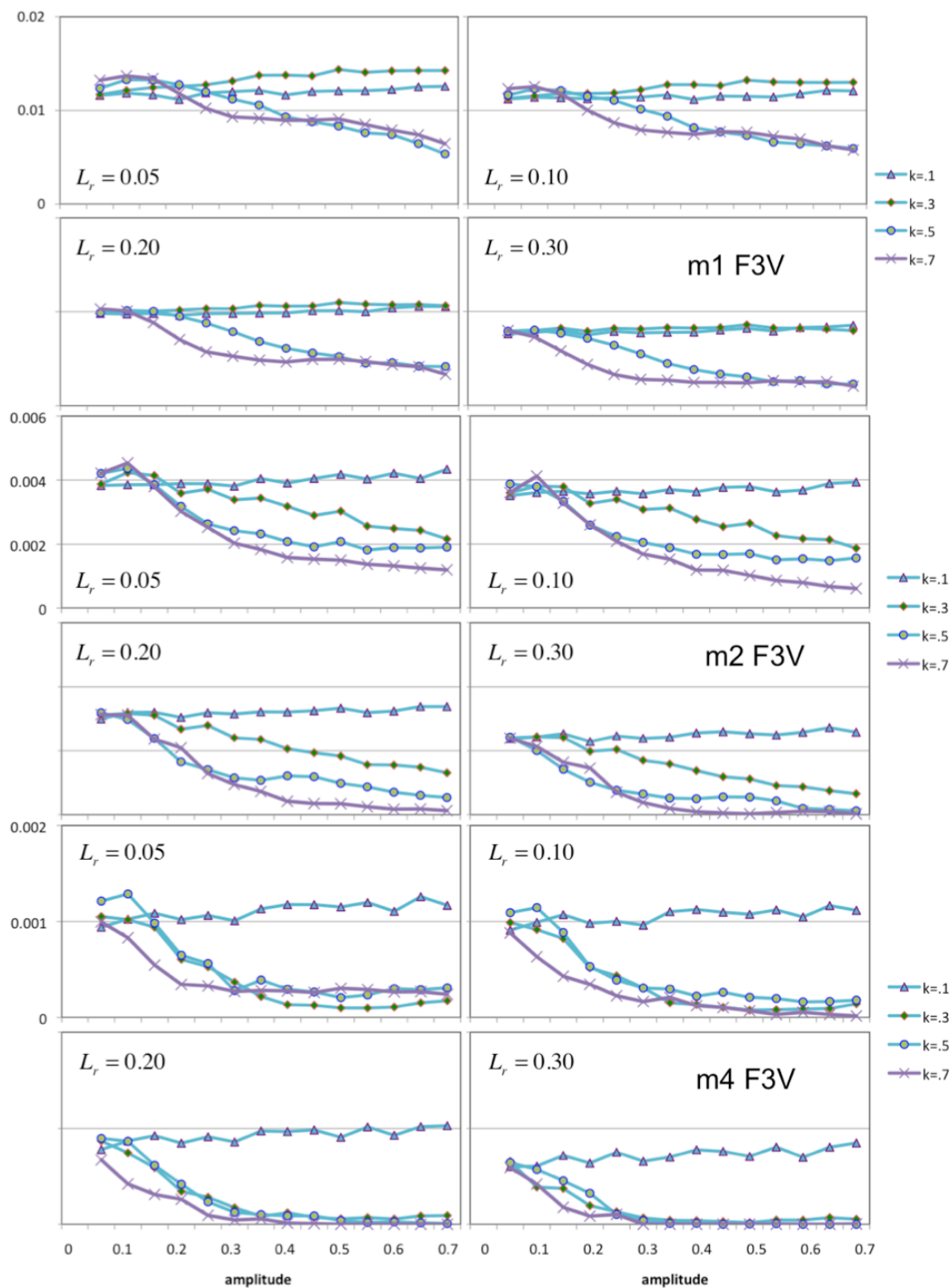
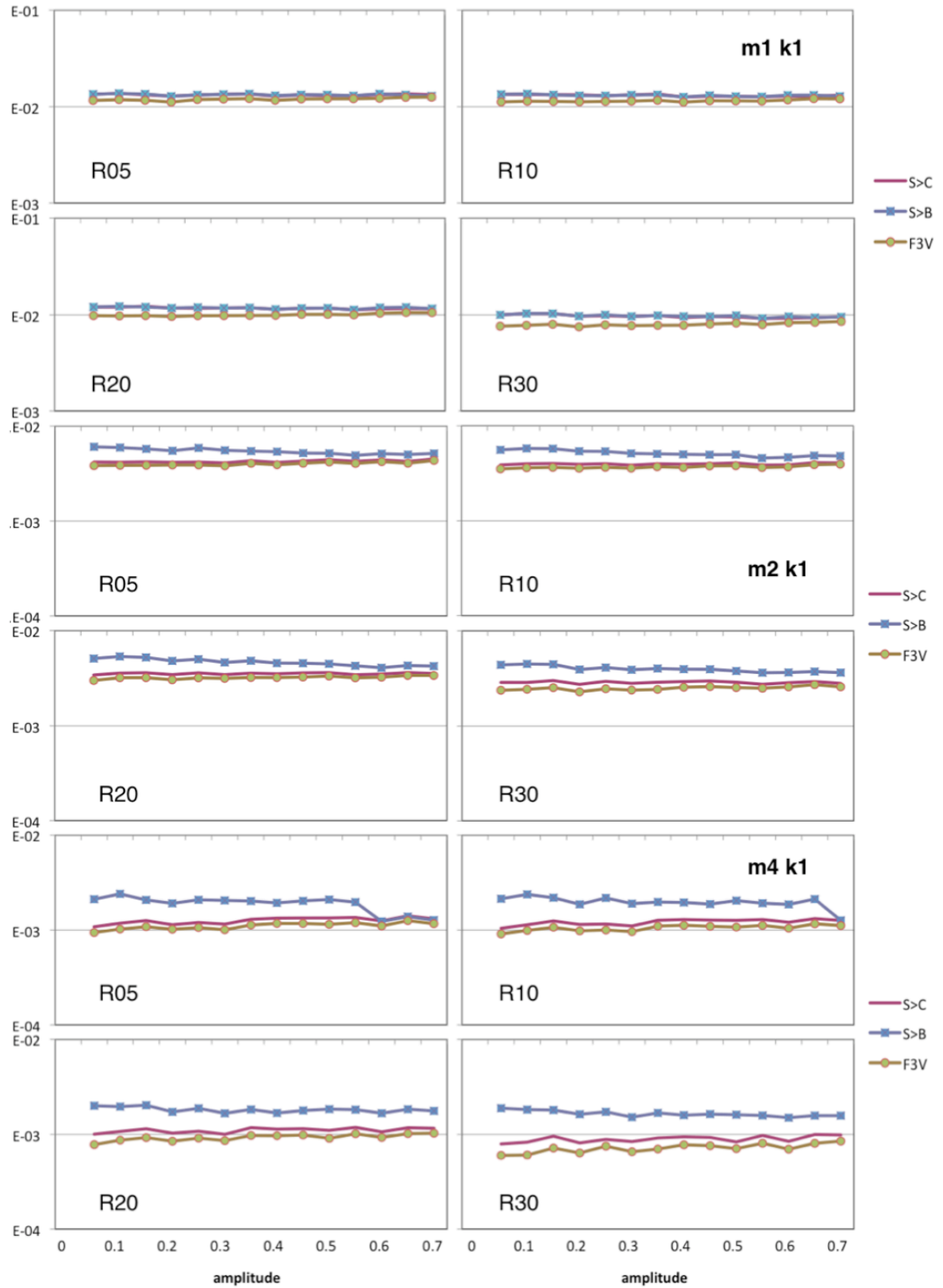
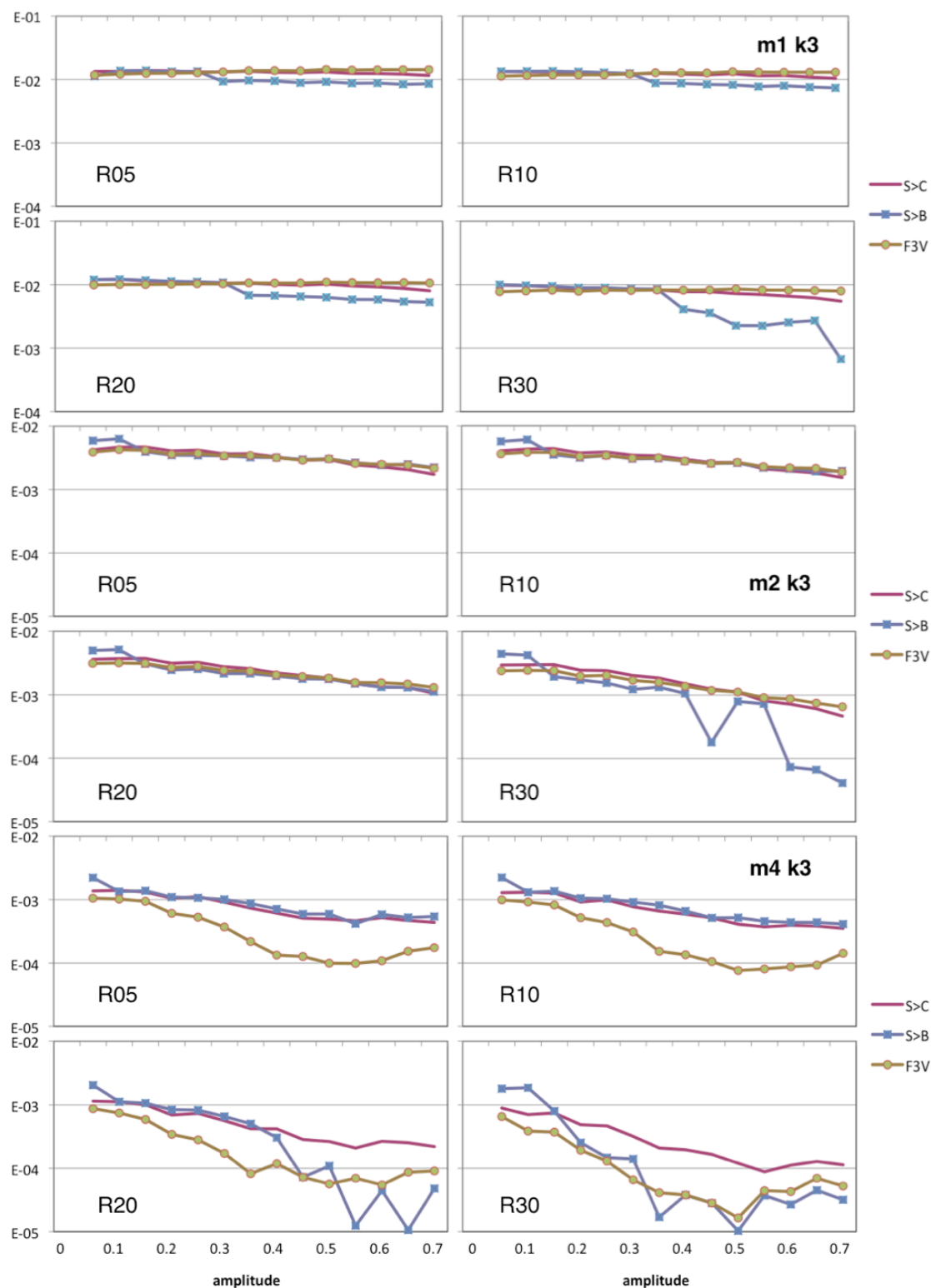


Figure D-27 Surfing Efficiency above three times the SW ram speed

Figure D-28 Surfing Efficiency Log plots for wavenumber $k = 0.1$

Figure D-29 Surfing Efficiency Log plots for wavenumber $k = 0.3$

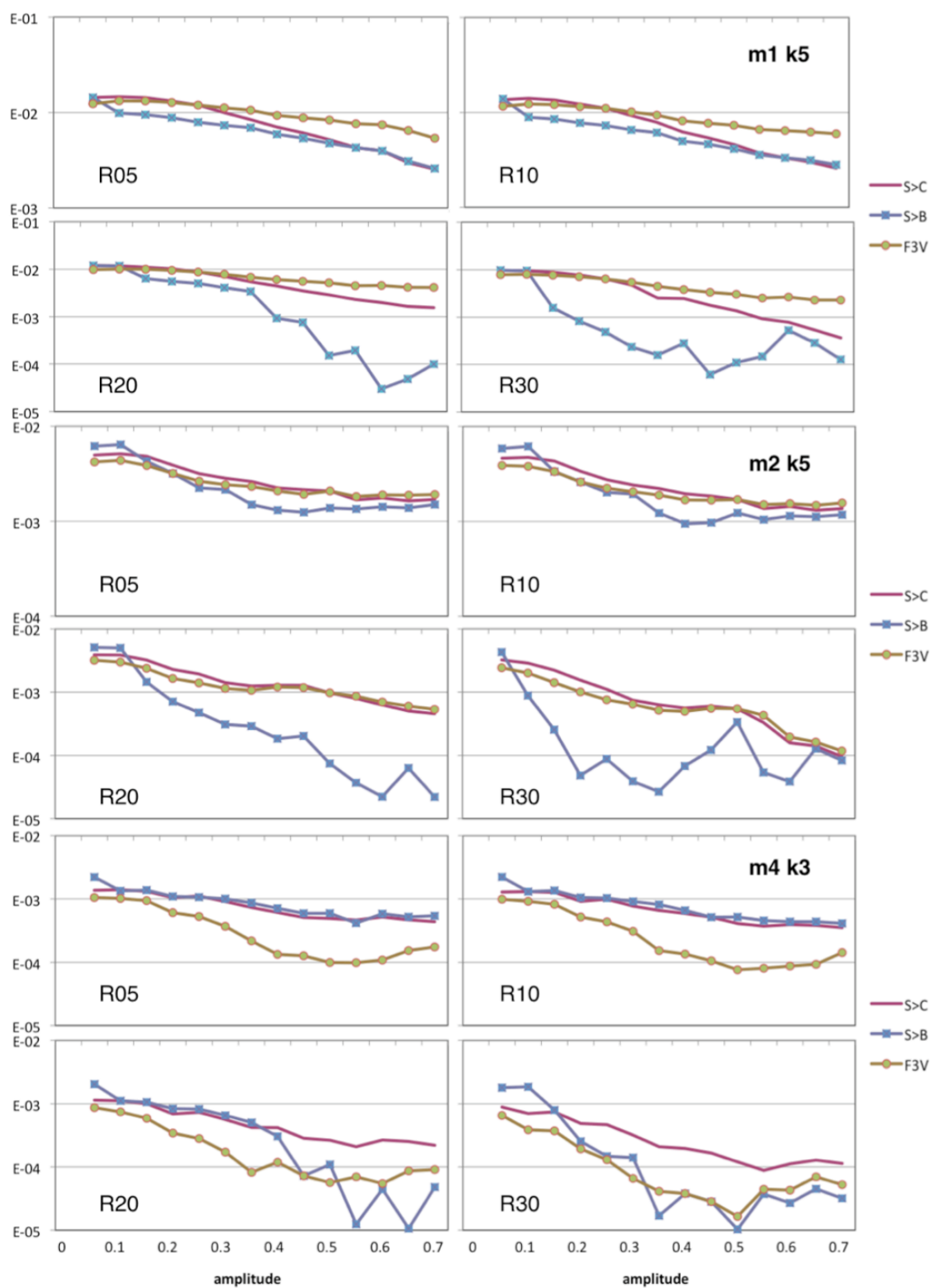
Figure D-30 Surfing Efficiency Log plots for wavenumber $k = 0.5$

Figure D-31 Surfing Efficiency Log plots for wavenumber $k = 0.7$

Appendix E Perturbed Constants of the Motion

We have shown that ripples in the shock front suppress the surfing mechanism, both in maximum attainable velocity and the flux of ions promoted above the shock drift ceiling. We now reconsider the effect of front perturbations upon the mathematical foundations of shock surfing as put forth by Lee et al (1996). We begin by recapitulating the conditions necessary for successful shock surfing derived by the authors and ask whether violation of any or all of these is consistent with our results. Recall that the basic assumptions made were such that, once a surfing ion has exceeded a few initial reflections:

1. The adiabatic invariant $\oint v_x dx$ over each bounce period should hold
2. The energy gain per bounce should rapidly asymptote to $\Delta\epsilon = qE_y\Delta y$ i.e., the longitudinal distance between reflections should become fixed
3. The bounce amplitude should monotonically decrease with increasing longitudinal velocity such that $\Delta x \simeq v_x^2 / \omega_c v_y$
4. Finally, the trapping should be efficient due the relation $v_x \propto v_y^{1/3}$ which follows from the above assumed conditions.

To record each ion's state completely throughout its flight in the simulation would have required more computing power than was available. Instead, we assume that, after a few bounces, aggregate data would be representative, and the following sets of plots were generated from ensemble averages. Upon every reflection, each test particle's state was accumulated in a bin for that bounce number, and the bins indexed and

incremented accordingly. In order to determine *post mortem* the ensemble spread for each bin, a second set of bins accumulated the squares, and both were averaged at the end of the run from which the standard deviations were calculated.

We begin by examining the populations of the bins themselves and show in Figure E-1 below a set of four plots representing H^+ (top pair) and He^+ (bottom pair) at two perturbation strengths, weak $a = k = 0.1$ (left pair) and strong $a = k = 0.5$ (right pair). The function $f(n)$ where $n \in \{1, \dots, n_{\max}\}$ is the bounce ‘bin’ number, is evidently a smooth function when the perturbation is weak (the left hand plots are truncated) but drops rapidly for stronger perturbations when approaching $n_{\max}$ a value that decreases with amplitude, wavenumber and ion mass. It is essential to note that the range is large – over five orders of magnitude – and no distinction was made between the fastest and slowest ions. That is, a particular parameter bin number, say n_j for the adiabatic invariant, will contain information from ions that terminated there as well as for those that continued accelerating and therefore underwent many more reflections. Hence a substantial variance is to be expected in the lowest bins, diminishing as $n \rightarrow n_{\max}$ where only the very fastest and fewest arrive. The paucity of bin entries in the tail is evident and, by comparison, we see that the suppression of $n_{\max}$ at identical perturbation strength is more profound on heavier ions.

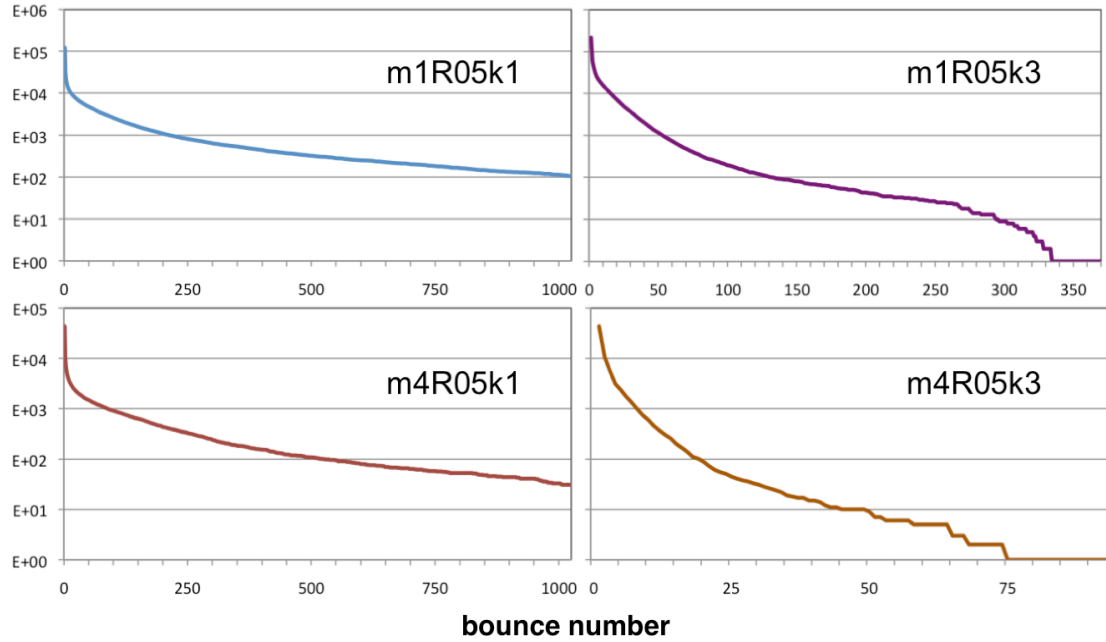


Figure E-1 Bounce Number Bin Counts

Mach 5 shock with thin ramp $L_r = 0.05$ is shown in all panels. Upper panels represent H^+ with He^+ below. Weak perturbation $a = k = 0.1$ is shown on the left, somewhat stronger $a = k = 0.3$ on the right. The functions $f(n)$ where n is the bin number are not normalized, but the total number of particles counted during the simulation run.

We focus our attention primarily on the thinnest ramp physically likely to be encountered $L_r = 0.05$ where we expect to see the fastest accelerated particles. We note that this was in the middle of the limited range explored in the baseline study of the adiabatic invariant, which already showed substantial deviation from theoretical constructs even at the perfectly smooth front.

We now examine the effect of perturbations on the four variables – adiabatic invariant, kinetic energy, bounce period and amplitude. In order to evaluate the correlation between these functions, we present them together in sets of four panels,

grouped at three representative corrugation strengths. We begin with the weak perturbation $a = k = 0.1$ applied to protons as shown in Figure E-2 below. The adiabatic invariant $\langle A \rangle$ is shown top left, kinetic energy $\langle E \rangle$ top right, bounce period $\langle T \rangle$ bottom left and bounce amplitude $\langle X \rangle$ bottom right. Each is plotted against the same bounce number, since they all represent the same ensemble of particles at the given shock and perturbation parameters.

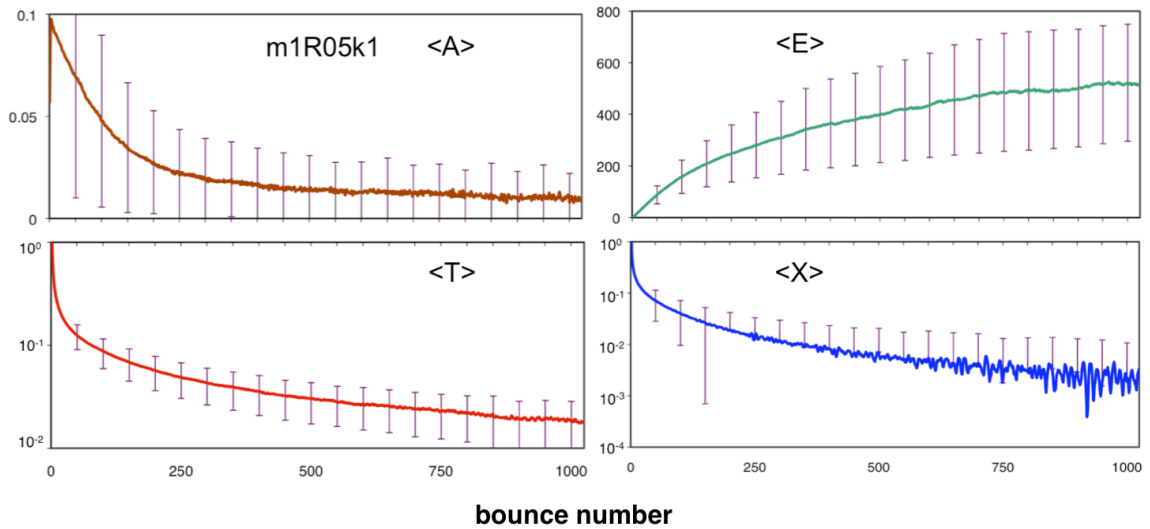


Figure E-2 Four Theoretical Variables defined by Lee et al
 $\langle A \rangle$ is the adiabatic invariant, $\langle E \rangle$ the energy per nucleon relative to the piston energy, $\langle T \rangle$ the bounce period and $\langle X \rangle$ the amplitude of normal excursion. All are functions of the bounce number. This set represents protons $m = 1$ at a thin ramp $L_r = 0.05$ with weak perturbation amplitude and small wavenumber $a = k = 0.1$. Note that the bounce amplitude $\langle X \rangle$ begins to show disruption in the distribution tail.

We first observe that all the functions are smooth, as should be expected for the weak case. We also see a large variance in the data as shown by the error bars. This is due to our method of binning, that is, the accrual in the bins of all surfing ions without

discrimination with respect to the final bounce number $n_{\max}$ of any particular particle.

Note that the limit for the fastest ions is not shown, since it exceeded the maximum of our bin arrays, which were limited to 1024.

We now consider each of the four functions in turn. To compute the Adiabatic Invariant, the normal velocity and displacement are integrated at each time step between bounces, and the sum recorded in the appropriate bounce number bin upon next reflection, the function is smooth and the mean value asymptotes to n^{-1} dependence as in the unperturbed case. The kinetic energy $\langle E \rangle$ is given per nucleon $v^2(n)/u_1^2$ and grows monotonically with bounce number. As in the smooth front case, comparison with the bounce period $\tau(n) \sim n^{-1/2}$ shows that the growth is linear in velocity as it should be under constant acceleration by the convective electric field i.e., $\varepsilon_K(t) \sim t^{1/2}$. The bounce period $\langle T \rangle$ is the elapsed time between reflections, and is also a smooth function of bounce number, in the weak case.

The bounce amplitude $\langle X \rangle$ is the perpendicular position difference between turning points for each particle at each bounce cycle. We see roughly the same functional dependence here, but notice that the perturbation has produced erratic behavior in the tail above $n \sim 700$. This finally gives us an indication as to the source of the large variance in these functions. As was seen in the trajectory plots, ions may be ‘launched’ upstream from the ramp, due to the combined effect of corrugation curvature and induced normal oscillation. The large variation indicates the corruption of the ion’s flight at its last reflection, regardless of the bounce number at termination of its flight.

This trajectory disruption can be clearly seen in the erratic behavior of all four functions as shown in the panels of Figure E-3 below. Here the perturbation strength is

increased to a moderate value for protons $a = k = 0.3$ and several features of interest are notable.

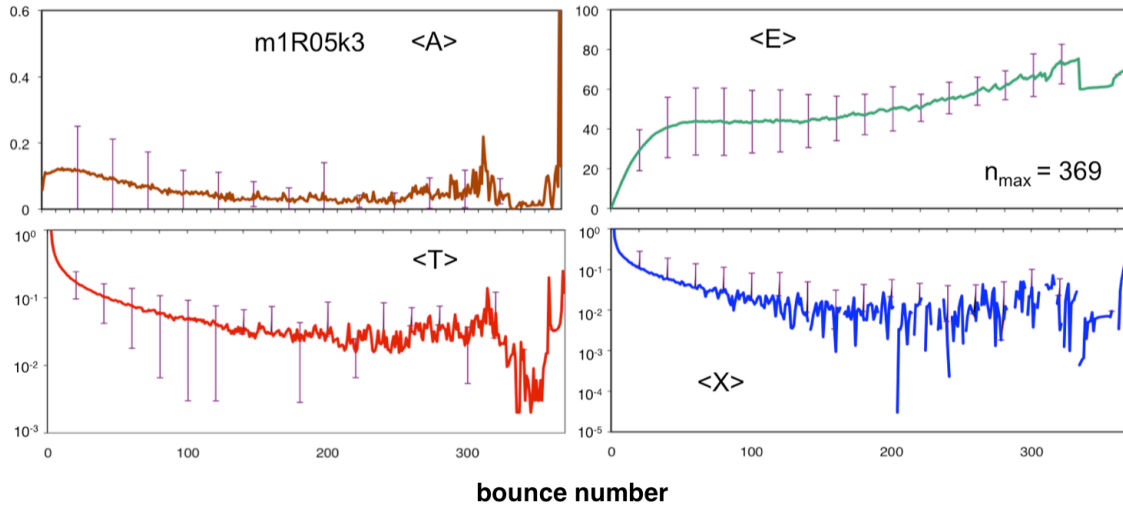


Figure E-3 Moderate perturbation strength for protons
Wavenumber and amplitude parameters are increased to $a = k = 0.3$
which suppresses the maximum bounce number to $n_{\max} = 369$.

The maximum bounce number has been suppressed, which explains the velocity limiting previously seen. The functions also become increasingly erratic, and as the bounce number approaches the maximum bounce number $n_{\max} = 369$, the scattering becomes so severe that the process quickly terminates even for the fastest ions. This is particularly visible in the last bin of the adiabatic invariant, but the smooth trajectory disruption is evident even for slower ions as can be seen in the bounce amplitude curve. Note also that the error bars disappear near $n = 330$, consistent with the corresponding bin plot in Figure E-1, and here in the tail at least, we finally observe data from the end of a single particle's flight.

The trajectory disruption due to shock front corrugation becomes even more evident in the last data set for protons shown in Figure E-4 where the perturbation strength has been increased to $a = k = 0.5$ which is larger, but still reasonably moderate (unity would be very strong.) Here we see that the maximum bounce number $n_{\max}$ has decreased to 43. Both the adiabatic invariant and the bounce period show small variances in their tails. Once again, the can see from the earlier bin plots that the population has dwindled to a few ions there. From the absence of large transient spikes in the data, we surmise that the velocities attained are insufficient to support violent breaches of the front. We note further from the ‘trajectories’ of the tail particles as deduced from oscillations in the period $\langle T \rangle$ and the amplitude $\langle X \rangle$ that normal oscillation resonances with the corrugations have become the predominant mechanism responsible for terminating their flight.

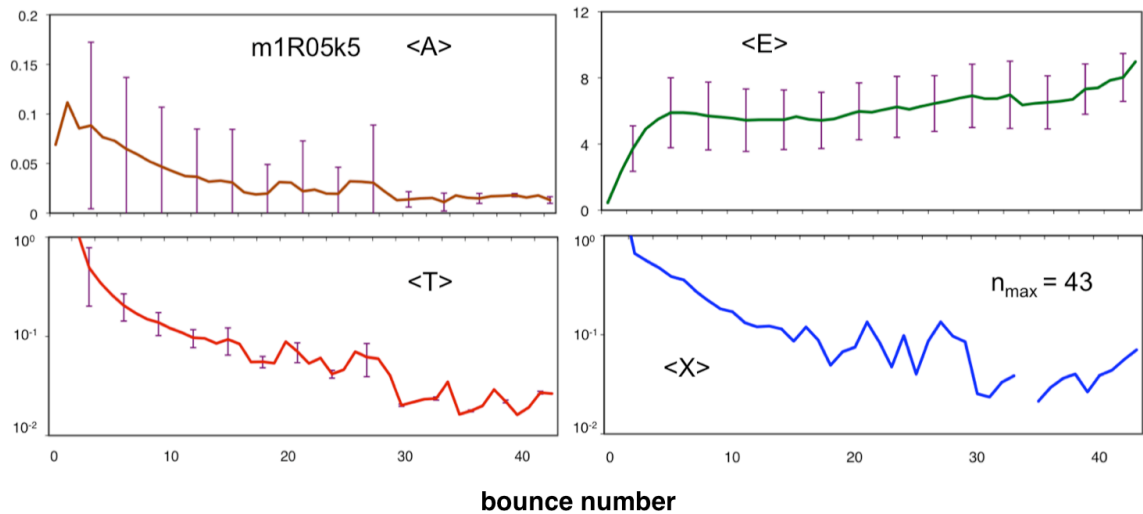


Figure E-4 Large perturbation strength $a = k = 0.5$
The maximum bounce number has been drastically suppressed

We now display a similar set of figures for the heavier He^+ but over a correspondingly smaller range of perturbation strengths in amplitude and wavenumber. We begin in Figure E-5 with $a = k = 0.1$ and see that, as was true for protons at this mild strength, all but the bounce amplitude appear well behaved. For the latter, we see substantial deviation in the tail, even though the variance remains relatively unchanged. This may indicate large deviations for those ions whose flights are prematurely terminated.

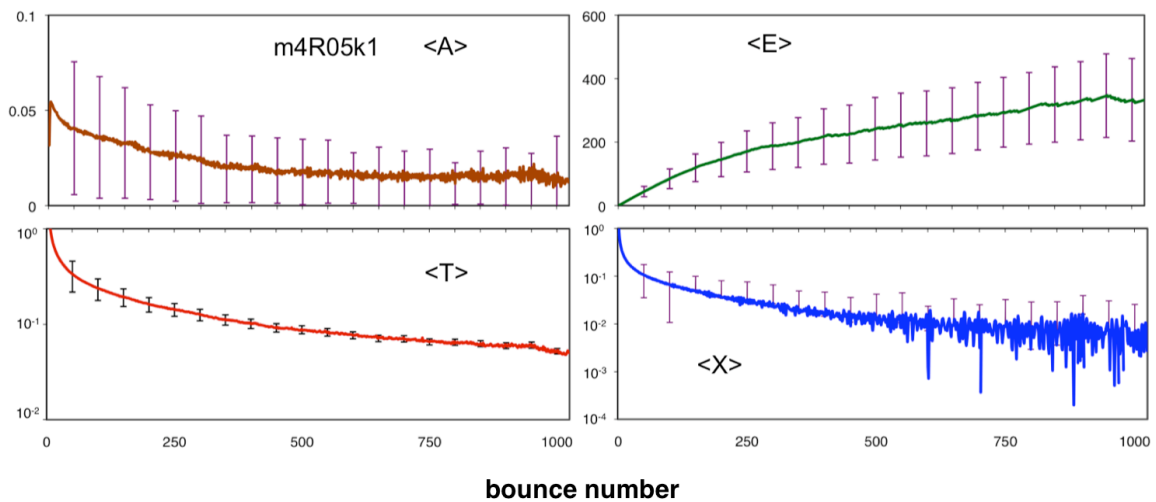


Figure E-5 Lee variables for He^+

ions of mass $m = 4$ at a weakly perturbed front where $a = k = 0.1$ begin to show deviation in the bounce amplitude earlier than do protons.

In Figure E-6 we show perturbation strength $a = k = 0.2$ and note a commensurate and possibly even greater ‘disruption’ in the trajectories for this heavier ion than we saw for protons at greater perturbations strength. At bounce number approximately $n = 528$ the tails of all four variables become highly erratic.

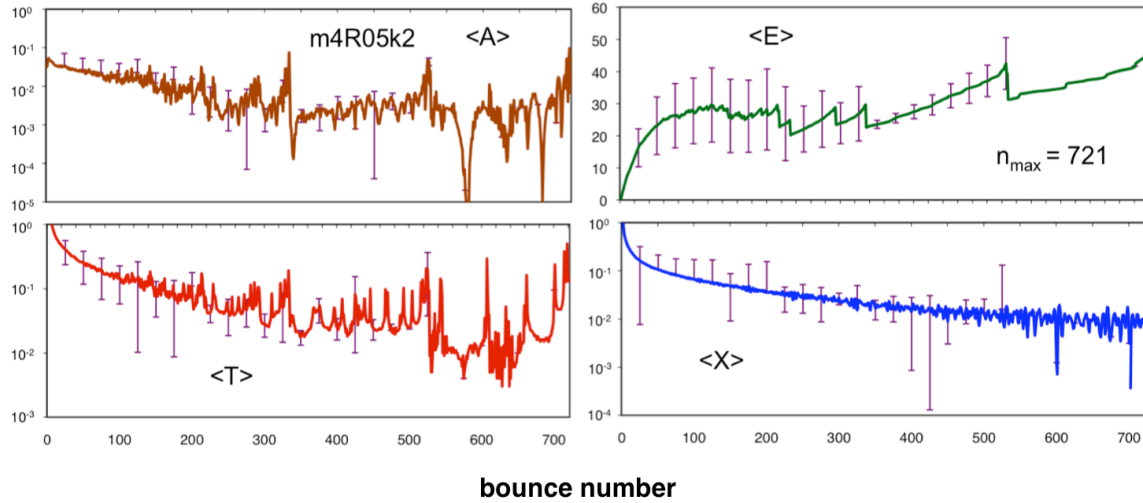


Figure E-6 Singly Ionized Helium at $a = k = 0.2$

Modest perturbation strength for He^+ ions. After a few hundred bounces, the bounce period (bottom left) erratically increases. Longer time-of flight leads to greater gyro-phase rotation at the next shock encounter.

Finally, we offer a last set of He^+ in Figure E-7 at strength $a = k = 0.3$ and compare it with protons at the same perturbation strength. Here the maximum bounce number has been suppressed to $n_{\text{max}} = 93$ which shows that this perturbation is strong for Helium while only moderate disruptive $n_{\text{max}} = 369$ for protons at the same strength. Once again, we see a correlation between the onset of trajectory disruption at around $n = 50$ with large spikes at $n = 75$ which signals the end of acceleration for all but the single remaining particle which was fortunate enough to begin the process in precisely the right location in velocity space. But even for this lone particle, we see large variation in all four panels. The bottom two are particularly revealing as they show a dramatic increase in both bounce period and amplitude. Clearly, the particle has left the ramp, and behaves as if on a typical single-bounce trajectory.

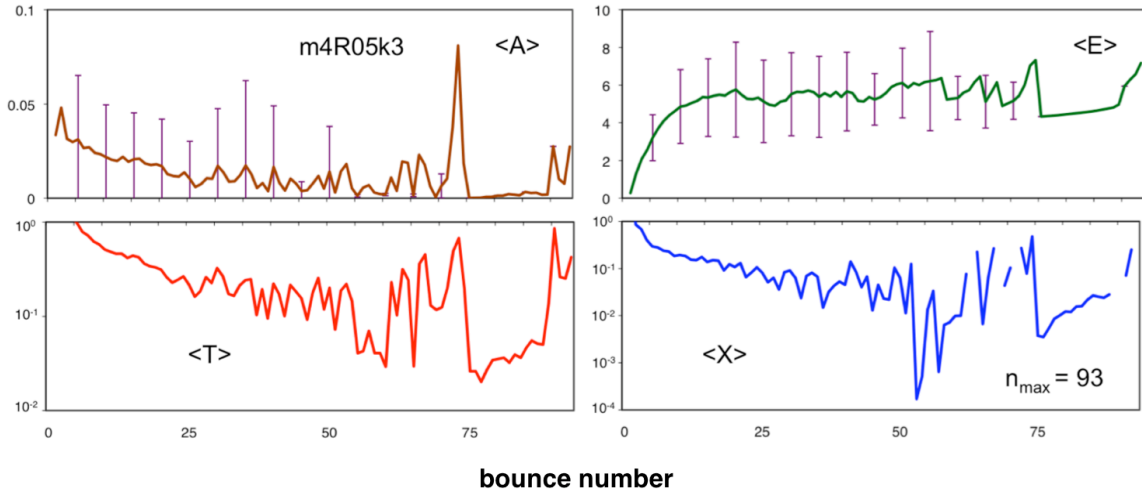


Figure E-7 Singly Ionized Helium at $a = k = 0.3$

Trajectories of He^+ ions become violently disrupted much earlier. By bounce number 75, the ensemble is decimated, and the bounce period (bottom left) of remaining particles increases to an ion cyclotron period.

In presenting our analysis of these statistical ensembles, we reiterate that our counting method is not ideal since far better and more complete information would be gleaned from the histories of individual particles or, at least, families of particles with the same total number of reflections. By its nature, our method integrates out much information of interest, providing only statistical moments. However, we feel that the insight gained through examination of the tails has been revealing enough to support our physical model of surfing disruption.

Discussion

From analysis of the period and amplitude of oscillations normal to the shock, we conclude that the imposition of ripples along the front, quickly cause surfing ion trajectories to become erratic. Heavier ions have difficulty but some do manage, by their greater inertia, to remain near-ramp (but not ramp-bound) longer. However, even those

ions fortunate enough to reach a few dozen reflections also eventually succumb, are ejected far from the front, and return with normal velocity sufficient to breach the shock.

These effects explain not only the velocity maxima suppression, but also the thinning of the accelerated fluxes seen in the distribution functions. Even for perturbations of moderate strength, ions incident on the shock outside a very limited region of velocity space are soon driven by perpendicular 'kicks' as they traverse the ripple crests and are thus prematurely 'ejected' in the process.

We have shown that when the front is perturbed, even by a sinusoidal corrugation, the mathematical foundations developed in the Lee model of shock surfing fail. The data show this, the trajectory plots show this, but analysis of the bounce amplitude and period of individual particles perhaps demonstrate more clearly why the mechanism fails.

Appendix F Standing Waves

Standing Wave Drift

In order to validate the results and analyses from external studies and to provide a basis for deriving a DSA scattering parameter, we performed our own simulations by employing waves standing in the shock frame. We stress that this was an exercise intended for reconnaissance only, as it does not model precisely the same mechanisms as were discussed in the review of turbulent scattering. Even so, we assert that a common parametric dependence should govern their effects, and explore whether front return mechanisms other than DSA may function at small planetary shocks.

For this study, we introduce magnetic waves standing in the shock frame and track the evolution of the test particle velocity distribution within several disjoint regions downstream. In this way, we may observe the temporal evolution of injected pickup ion shells and compute the time elapsed for the ensemble to reach each region, as well as the diffusion rate $\langle \Delta \tilde{v} / \Delta t \rangle$ as a function of coarser time intervals. We record the entire distribution function in each of the multiple regions, and perform post-run analyses on the features of interest.

Test particle simulations were run with strictly magnetic waves imposed over the ambient solar wind fields we had previously employed. Pickup ions were injected at the foot of the ramp, and their progress through the simulation box recorded by integrating the distribution function of all particles in each of the regions. Only the perpendicular velocity component $v_{\perp} = \sqrt{v_x^2 + v_y^2}$ is recorded.

We included the standard shock-normal electric field, but chose a broad ramp width $L_r = 2.0$ in order to limit shock surfing, and thus isolate the initial study to ions with velocities boosted strictly by shock drift. A baseline run was performed without waves as shown in Figure F-1, which demonstrates that the profile remains essentially constant once the ions leave the vicinity of the front.

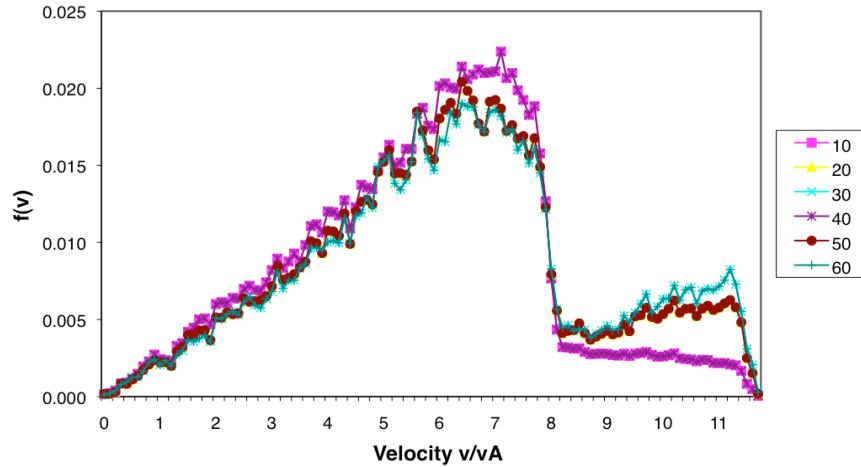


Figure F-1 TP Distribution Functions with No Waves

Plots for six simulations box regions are shown, at distances in multiples of ten ions inertial lengths downstream of the shock. The ramp width $L_r = 2.0$ ensures that only the shock drift mechanism accelerates the ions, as seen in the first region, which includes the ramp itself.

Test particle simulations were also run with waves standing in the fluid frame. Regardless of variation in amplitude and wavenumber, little or no effect was observed on the pickup ion velocity distribution, which was not surprising since Karimabadi (1992) notes that “particle energy in the wave frame is constant.” Simulations were then run to determine the effect of waves standing in the shock frame. First, single waves were imposed along each of the coordinate axes, which varied in amplitude and wavenumber, and the evolution of the distribution function with distance is evident by the example

shown in Figure F-2. Here, standing waves were imposed whose magnetic field varied as $B_y(x) = a_y \cos(kx)$ with extremely long wavelength $k = 0.01$ and large amplitude $a = 2.5$. The amplitude of this particular wave corresponds to a peak magnetic energy density $\varepsilon_B \sim a^2$ or more than six times that of the ambient solar wind.

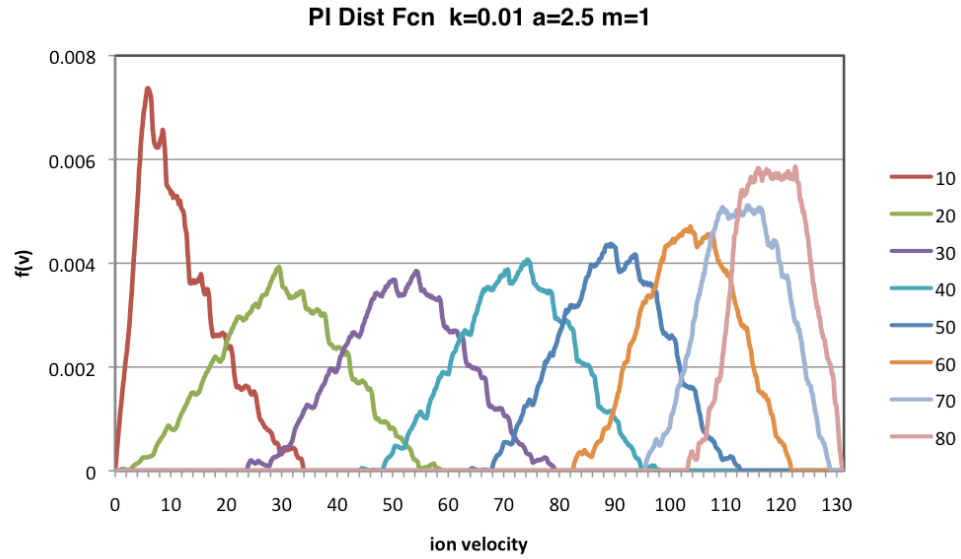


Figure F-2 Distribution Function Evolution

Magnetic waves $B_y(x) = 2.5 \cos(0.01x)$ standing in the shock frame are imposed downstream over the ambient solar wind fields. Distribution functions are integrated in disjoint regions of length $10\lambda_{pi}$. The maximum proton velocity increases with distance, and therefore with time, while the variance appears to narrow. Velocity is given in Alfvénic units.

A polarization study was then performed where the waves were defined in turn upon each of the complementary coordinate pairs e.g., $\delta B_x(y)$ and $\delta B_x(z)$ and results plotted below in Figure F-3. All but the two polarizations $B_y(x)$ and $B_y(z)$ show no effect at all, even at the highest amplitude $a = 2.5$.

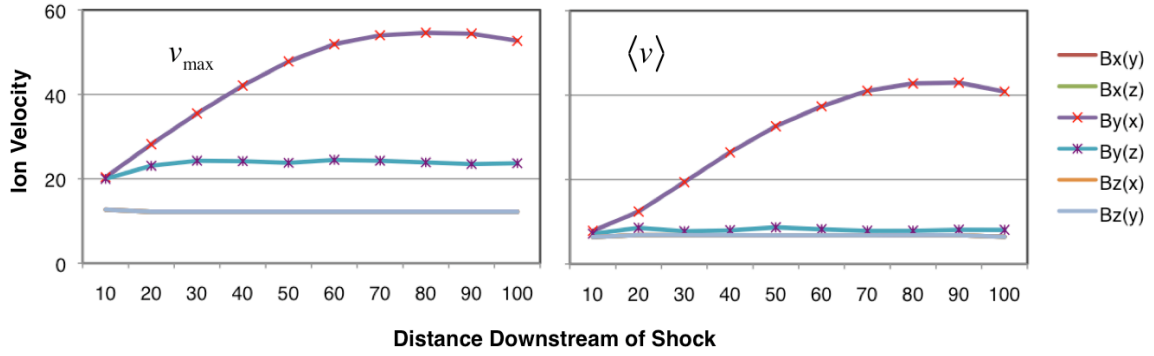


Figure F-3 Vector Polarization Study

The maximum and mean velocities perpendicular to the solar wind magnetic field are integrated in each of ten downstream regions. The wave amplitude is $a = 1.5$ with wavenumber $k = 0.01$. Only the $B_y(z)$ and $B_y(x)$ curves are distinct, while other polarizations are overlaid, indicating identical data with no velocity gain for pickup ions at all.

Since the magnetic wave orientation $B_y(x)$ produces by far the greatest effect, we continue our studies by varying the other parameters holding this field orientation fixed. Data from the wavelength study are shown in Figure F-4 below, and the results appear consistent with the turbulent diffusion studies already cited, namely, that velocity gain varies inversely with wavenumber. However, we see that this is simply due to the imposed harmonic function, and therefore only a local effect. Nevertheless, this localization may still be quite large with respect to ion dynamics near the shock front since, for example, the wavelength corresponding to $k = 0.01$ is $\lambda = 630 c / \omega_{pi}$, several hundred ion inertial lengths. Note also that, for wavelengths shorter than that of the box, the maximum ion velocity does not return to the shock drift ceiling, but oscillates about a higher value dependent upon the wavenumber. In other words, the velocity gain is both secular and persistent.

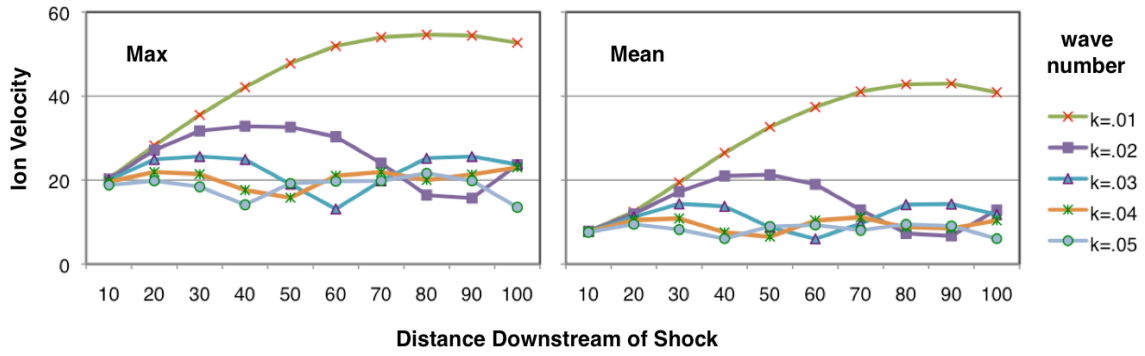


Figure F-4 Wavenumber Study
Diffusion strength varies inversely with wavelength

Moreover, the perturbation must be extremely strong. Very little effect was seen unless the energy density exceeded five times that of the background. The most profound gain occurred for $\delta B_y(x)$ with negligible results for waves along the other coordinate axes, regardless of amplitude and wavenumber.

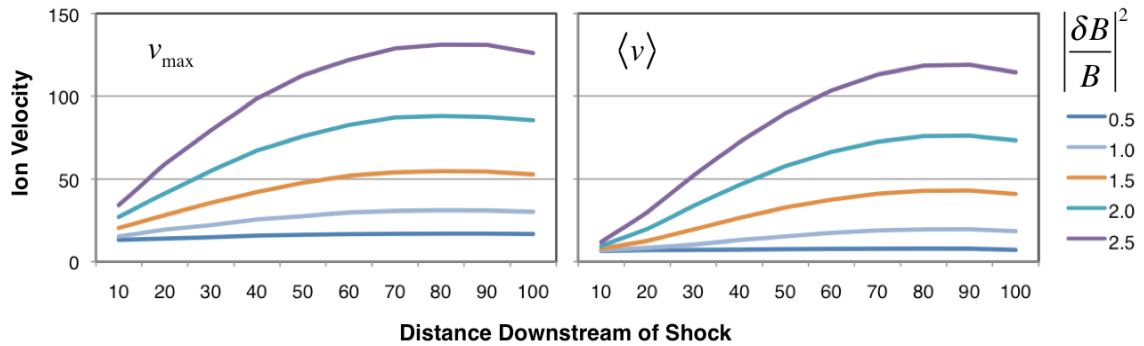


Figure F-5 Amplitude Study

These plots show that both the peak and mean diffusion rates scale with wave energy density, i.e., field amplitude squared

Shock drift ions typically do not return again after a single upstream crossing. Most notably, only the very longest waves $k_y = 0.01$ at the highest amplitude $a_y = 2.5$ were able to return ions to the front a second time. Note also that, as found by Giacalone and Ellison, a wavelength-dependent velocity ceiling for this mechanism exists.

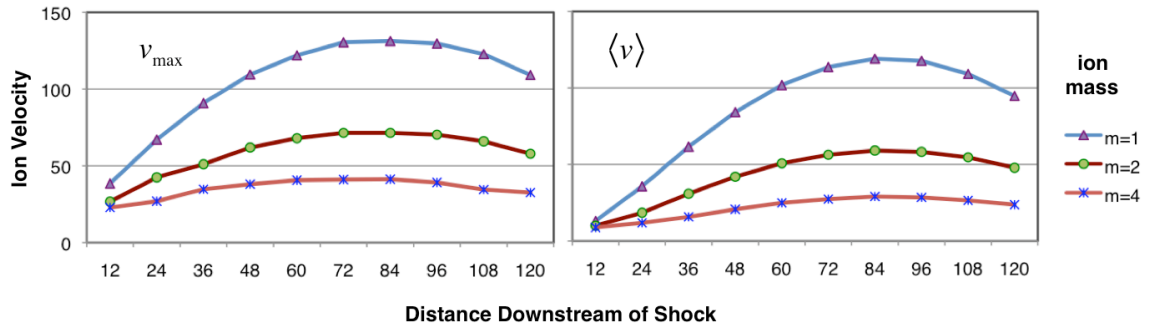


Figure F-6 Mass Study

This study verifies that the diffusion rates scales inversely with mass

Simulations were also run for the same wavelength at several amplitudes, varying the ambient magnetic field up to 15 degrees from the perpendicular. As seen below in Figure F-7, velocity gain becomes significantly suppressed at an angle of only $\Delta\theta = 5^\circ$ and becomes progressively worse with increasing deviation from the perpendicular. The same scaling was seen at all amplitudes.

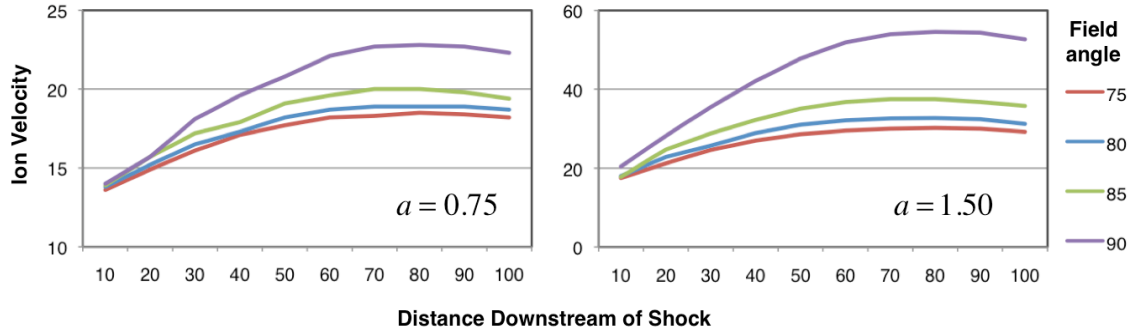


Figure F-7 Oblique Shock Study

Velocity maxima are displayed at two wave amplitudes with a family of curves representing the same set of field angles θ_{Bn} . Velocity diffusion decreases with deviation from the perpendicular.

We also conducted a short series of simulations with time-dependent waves at the polarization and amplitude that produced the greatest static gain. Only frequencies well below the ion cyclotron frequency affected the velocity distribution, while those approaching and exceeding the ion cyclotron frequency $\omega_{\delta B} > \Omega_{cp}$ produced vanishingly small results.

Multiple Waves

Sets of one, two and three wave-vector combinations were then imposed, each of whose energy was set to one-third, one-half and unity with respect to the ambient solar wind magnetic energy density. Three groups of simulations were run at different wavelength scales up to that of the magnetosphere. To ensure that $\nabla \cdot \mathbf{B} = 0$, each vector's wavenumber was defined along a coordinate axis as for the single-wave studies above, and its magnetic vector amplitude defined along another. The perpendicular pair was then rotated through pre-defined sets of Euler Angles. The divergence of a sum equals the sum the divergences, and the divergence is coordinate-independent and thus

invariant under fixed rotation. The results were not substantially different from those of the single-wave studies. Only ensembles where the total energy exceeded the ambient field energy and whose field vectors were predominantly in the y -direction succeeded in producing any velocity gain at all.

Conclusions

For the case of magnetic waves standing in the shock frame, pickup ion velocity diffusion scales directly with the wave energy density, directly with the wavelength and inversely with ion mass. A velocity ceiling is reached which is exceeded regardless of further elapsed time, and is directly proportional to the maximum wavelength of the magnetic perturbations. Wave polarization is a factor, and the strongest effect is obtained when the wave vector is parallel to the fluid flow, and the field is perpendicular to both the shock-normal and the ambient magnetic field of the fluid. Finally, diffusion is sufficient to return pickup ions to the front from downstream only when the perturbation energy density exceeds a factor five times above that of the background magnetic field. We offer below an additional derivation of velocity 'diffusion' due to the magnetic gradient drift.

Supplemental Drift Derivation

Recall that the magnetic gradient drift (shock drift mechanism) contains the vector cross product $\mathbf{v}_{\nabla B} = \mathbf{B} \times \nabla B$. At a purely perpendicular shock in the absence of waves or other perturbations yields a drift through the front along the direction of the convective electric field. Since $\mathbf{u} = u\hat{\mathbf{x}}$, $\mathbf{B}_1 = B_1\hat{\mathbf{z}}$ and $\nabla B = \partial_x B\hat{\mathbf{x}} \simeq (B_2 - B_1)/L_r\hat{\mathbf{x}}$ we have $\mathbf{v}_{\nabla B} \sim \mathbf{B} \times \nabla B = \hat{\mathbf{z}} \times \hat{\mathbf{x}} = \hat{\mathbf{y}}$ with the convective field unit vector direction $\hat{\mathbf{E}}_C = -\hat{\mathbf{u}} \times \hat{\mathbf{B}} = -\hat{\mathbf{x}} \times \hat{\mathbf{z}} = \hat{\mathbf{y}}$ so that the drift along the electric field direction results in a net increase in ion energy, $\Delta\epsilon = qE_C\Delta y$.

Now consider a perturbed region downstream of the shock where the field consists of the uniform ambient magnetic field $\mathbf{B}_2 = B_2\hat{\mathbf{z}}$ with a superimposed standing wave $\delta\mathbf{B}(x, y, z)$ of arbitrary polarization. The total field is $\mathbf{B}(x, y, z) = B_2\hat{\mathbf{z}} + \delta\mathbf{B}(x, y, z)$ and, since the background field is uniform, $\nabla B_2 = 0$ so that $\mathbf{B} \times \nabla B = (\mathbf{B} + \delta\mathbf{B}) \times \nabla \delta B$.

A static field of arbitrary direction may be expressed as

$$\delta\mathbf{B}(x, y, z) = \delta B_x(x, y, z)\hat{\mathbf{x}} + \delta B_y(x, y, z)\hat{\mathbf{y}} + \delta B_z(x, y, z)\hat{\mathbf{z}}$$

and, taking $B(x, y, z) = |\mathbf{B}(x, y, z)|$, we have the cross product (abbreviated for clarity)

$$\mathbf{B} \times \nabla B = \det \begin{vmatrix} \hat{\mathbf{x}} & \hat{\mathbf{y}} & \hat{\mathbf{z}} \\ B_x & B_y & B_z \\ \partial_x B & \partial_y B & \partial_z B \end{vmatrix} = \begin{pmatrix} (B_y \partial_z B - B_z \partial_y B) \hat{\mathbf{x}} \\ (B_z \partial_x B - B_x \partial_z B) \hat{\mathbf{y}} \\ (B_x \partial_y B - B_y \partial_x B) \hat{\mathbf{z}} \end{pmatrix}$$

Note that the imposed field perturbation $\delta\mathbf{B}(x, y, z)$ stands in the shock frame and therefore does not contribute to the convective electric field, which remains in the $\hat{\mathbf{y}}$ direction. Hence, both the $\hat{\mathbf{x}}$ and $\hat{\mathbf{z}}$ components produce nothing more than drifts of the

guiding center, while the $\hat{\mathbf{y}}$ component alone leads to energy gain. Writing that

component out explicitly, we have $(B_{z0} + \delta B_z) \frac{\partial \delta B}{\partial x} - \delta B_x \frac{\partial \delta B}{\partial z}$. Assuming the functional

form $\delta \mathbf{B} = ae^{i\mathbf{k} \cdot \mathbf{r}} \hat{\mathbf{r}}$, gives $\delta B_{x_i} \frac{\partial \delta B}{\partial x_j} = a^2 k_{x_j} e^{i(2\mathbf{k} \cdot \mathbf{r})}$ which is diminished in strength – first,

since the very small wavenumber k_{x_j} is now a multiplicative factor and also because the

wavenumber in the argument of the harmonic function is effectively doubled. Hence,

even in the case of large amplitude magnetic fluctuations, we may dismiss the

contribution from second term and, taking the fluid flow $\mathbf{u} = u\hat{x}$ into account, we find the

most likely candidate for the perturbation in the remaining term $B_{z0} \frac{\partial \delta B}{\partial x}$ is the

Magnetosonic wave $\delta \mathbf{B} = \delta B(x) \hat{\mathbf{y}} = ae^{i(k_x x - \omega t)} \hat{\mathbf{y}}$.

We again remind the reader that this exercise is intended to provide us with yet another basis for constructing the scattering strength parameter, which may then be applied either directly into the Webb formalism or, by way of a Master Equation into the Fokker-Plank equation. At the very least, it gives us justification for treating velocity diffusion in this case as a two-dimensional random walk with variable step size, dependent upon ion charge and mass, perturbation field strength and wavelength and, perhaps most importantly, the instantaneous ion cyclotron velocity.

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