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Romeo, Orlando Marco

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An Interdisciplinary Approach to Novel In Situ Measurements of Solar and Planetary
Electrons and Magnetic Fields

by

Orlando Marco Romeo

A dissertation submitted in partial satisfaction of the

requirements for the degree of

Doctor of Philosophy

in

Earth and Planetary Science

in the

Graduate Division

of the

University of California, Berkeley

Committee in charge:

Professor Michael Manga, Chair

Doctor Davin Larson

Professor Bruce Buffett

Professor Stuart Bale

Fall 2024

An Interdisciplinary Approach to Novel In Situ Measurements of Solar and Planetary
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Abstract

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For millennia, humanity has studied celestial objects through ground-based observations, including the Sun, planets, moons, and comets. Besides the orbits of these astronomical bodies, we have also observed space-based phenomena from the Earth, such as sunspots, solar eclipses, and aurorae. The emergence of space exploration and technological advancements in the mid-20th century enabled us to then directly probe the realm beyond our planet through in situ measurements. These innovations in direct, localized observations contributed significantly to addressing numerous questions related to the solar wind, planetary magnetospheres, and space weather. However, it was only around 20 years ago that the term *Heliophysics* was finally established to describe the interdisciplinary branch of science that connects the study of the Sun to its interactions throughout the solar system. This field was broadened to include the entire heliosphere, a vast region of space, influenced by the solar wind and the solar magnetic field, extending beyond the orbits of the outer planets. By framing our scientific approach to connect the solar corona, interplanetary medium, and geospace, we can gain a more cohesive understanding of the entire system.

Space-based technology for in situ measurements continues to advance, especially in studying the inner heliosphere, defined as the region extending from the Sun to the orbit of Jupiter. One example is the Parker Solar Probe (PSP) mission, which has traveled closer to the Sun than any previous spacecraft to study solar wind, solar activity, and the mechanisms that drive these phenomena. PSP includes a whole suite of instruments to examine the solar environment down to about $10 R_S$ (solar radii) from the Sun. The mission primarily focuses on in situ measurements of the solar wind from intervals of quiescent conditions to large solar eruptive events. Another example is the Interior Exploration using Seismic Investigation, Geodesy and Heat Transport (InSight) mission on the surface of Mars to study the interior structure and seismic activity of the planet. The lander includes a magnetometer to provide the first surface magnetic field measurements on Mars. Additional magnetic field observa-

tions were later conducted by the Zhurong mission at 16 different locations on the Martian surface. This dissertation is dedicated to investigating these novel in situ measurements in different areas of Heliophysics, specifically focusing on the evolution of electrons and magnetic fields in the inner heliosphere. By analyzing spacecraft and lander measurements, along with numerical simulations, we investigate various environments between the Sun and Mars to connect different systems in the heliosphere for future studies.

We first aim to investigate the mechanisms that govern solar wind acceleration and energy balance near the Sun using PSP in situ measurements. These processes are theorized to be strongly influenced by the solar wind heat flux and ambipolar electric field, which can be studied through electron velocity distribution functions (eVDFs). While the electron core and halo populations of the eVDF are nearly isotropic, the strahl beam component is more anisotropic and typically magnetic field aligned. Due to the asymmetry and non-thermal nature of the electron strahl, this population mainly contributes to the overall heat flux in the near-Sun environment. The strahl beam can undergo adiabatic focusing by the solar magnetic field or pitch angle scattering due to Coulomb collisions and plasma wave instabilities. We can study the evolution of the electron strahl over a wide range of heliocentric distances under ambient solar wind conditions using the onboard Solar Probe ANalyzer-Electron (SPAN-E) instruments. Multiple calibration procedures are introduced to properly estimate various parameters of the solar wind electrons, such as electron density and temperature. As a proxy for strahl scattering, we fit to the beam distribution to obtain an angular strahl width and direction. We find the strahl beam to narrow with increasing electron energy. In addition, the strahl width is observed to vary with heliocentric distance, electron density, and electron collisional age. Our study suggests that close to the Sun, Coulomb collisions primarily scatter the strahl electrons, until whistler waves become the dominant mechanism above $30 R_S$ from the Sun.

In the following chapter, our focus shifts to large solar eruptive events observed by PSP, specifically Interplanetary Coronal Mass Ejections (ICMEs). ICMEs originate from the Sun as Coronal Mass Ejections (CMEs), large structures of coronal plasma and magnetic fields that travel through the interplanetary medium. These structures can interact with planetary environments to cause space weather phenomena, such as geomagnetic storms, potentially resulting in widespread disruptions to power systems and communication networks. Extensive research has been conducted using in situ instruments to study these events as they reach Earth or other planets. However, directly observing ICMEs close to the Sun will improve our understanding of their evolution and provide better constraints for numerical models. An example of such an event occurred on 05 September 2022 when PSP passed through a topologically complex, fast ICME near the Sun. Fortunately, this ICME did not impact Earth, but had the capacity to heavily affect geospace on a similar scale to the famous 1859 Carrington Event. PSP traveled through the flank and wake of the structure, collecting in situ measurements of the bulk plasma and magnetic field within $14 R_S$ from the Sun. We observe classical ICME signatures compared to structures that reach 1 AU

(astronomical unit) near Earth, along with two magnetic inversion line crossings. We also utilize remote-sensing observations from PSP, providing a novel visual perspective of the ICME close to the Sun. The event is modeled as two individual CMEs propagating from the Sun to achieve the best agreement with both in situ and coronagraph observations. In addition, a global reconfiguration of the Heliospheric Current Sheet (HCS) is identified after the eruption, suggesting that CMEs play a significant role in the evolution of the HCS. As the HCS extends across the heliosphere, these solar eruptions can therefore influence regions of space well beyond their direct travel path.

The manner in which solar wind and eruptive structures interact with other astronomical bodies depend on the presence of a global planetary magnetic field. This interaction can result in the formation of planetary magnetospheres, which may shield a planet from the substantial influx of particles and energy associated with ICMEs. While Mars currently does not possess an active geodynamo, it still exhibits an induced magnetosphere due to the solar wind and interplanetary magnetic field (IMF). However, an ancient dynamo may have existed in the past due to evidence of localized, remanent crustal magnetization, especially in the southern hemisphere. Thus, investigating these crustal fields may provide insight on the evolution of Martian magnetization. The final chapter of this dissertation analyzes the first in situ surface measurements of the Martian crustal magnetic field at the InSight and Zhurong landing sites. We model the Martian crustal magnetization with Monte Carlo simulations based on the altitude profile of crustal magnetic fields above each mission site. Our results provide improved constraints on magnetization coherence and depth scales near the surface below InSight and Zhurong. Furthermore, we develop new models of subsurface magnetization based on geological and topographic features to relate different regions of magnetization with distinct periods in Martian history. Our study suggests the occurrence of at least one magnetic polar field reversal in the past, along with evidence of a potential reactivation of a new dynamo (now extinct) following a weakened or dormant Martian geodynamo.

In each chapter, we utilize new in situ measurements conducted in different locations across the inner heliosphere. While the results and discussion in each section contribute to their respective sub-fields, it is also crucial to highlight their connection in future work. Understanding the mechanisms that control electron dynamics in the ambient solar wind is essential when studying electron physics within ICME structures. This research can then provide insight on ICME dynamics observed closer to the Sun, including their evolution and interactions with planetary environments. In the case of Mars, these interactions may have played a significant role in the climate evolution of the planet after the loss of its global magnetic field. Constraining the timing of the ancient Martian geodynamo through current and future measurements of crustal fields then becomes essential in better understanding this process. As our technological capabilities improve with time, novel in situ instrumentation could be extended to the outer heliosphere, further enhancing our knowledge of Heliophysics as an interconnected system.

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Chapter 1

Introduction

1.1 The Heliosphere

Solar Structure and Dynamics

Ionized plasma is the most abundant state of ordinary matter, constituting roughly 99% of the observable universe. Stellar bodies are one major source of plasma for planetary systems. The Sun accounts for over 99% of the total mass in the Solar System, composed of approximately 92.1% hydrogen, 7.8% helium, and less than 0.1% heavier elements (Moldwin 2022). Figure 1.1 illustrates the standard model of the Sun, with the solar interior separated into three primary regions (Phillips 1995). In the inner core, nuclear fusion converts hydrogen into helium, generating substantial amounts of energy. This region, extending roughly 2×10^5 km from the center, reaches temperatures exceeding 1.5×10^7 K with a density of $150\text{--}160 \text{ g cm}^{-3}$. The produced energy is then transported radially outward from the core to the radiative zone, where energy is transferred primarily by radiation. Due to the inefficient transport of emitted photons through dense plasma (up to 20 g cm^{-3}), the temperature in this region drops to 2×10^6 K once the radiation reaches 5×10^5 km from the center. The final region, the convection zone, is approximately 7×10^5 km thick and is driven by turbulent convective currents that transport energy from the inner layers (Stix 2002). These convection cells are important to solar dynamics as they govern the motion of solar matter and magnetic fields near the surface (Miroshnichenko 2023). At the solar surface, the density decreases to about $10^{-7} \text{ g cm}^{-3}$, along with the temperature to 5700 K.

Beyond the interior, the solar atmosphere comprises the photosphere, chromosphere, and corona (Moldwin 2022). The photosphere is the optically visible layer of the Sun with a thickness of 500 km and density of $5 \times 10^{-7} \text{ g cm}^{-3}$. This corresponds to a particle number density of $3 \times 10^{17} \text{ cm}^{-3}$, which decreases by almost two orders of magnitude at the boundary of the chromosphere (Miroshnichenko 2023). The chromosphere has a thickness of roughly 1500 km, reaching temperatures from about 6×10^3 to 2×10^4 K. The plasma density decreases with the height of the chromosphere, approaching values of $10^{-15} \text{ g cm}^{-3}$. The final outermost layer, the corona, exhibits temperatures exceeding 10^6 K.

Despite these high temperatures, the particle number density of the corona ranges between 10^9 and 10^4 cm^{-3} from its base to a distance of $10 R_S$ (solar radii). Due to the low density, the upper atmosphere is only visible during a total solar eclipse, appearing as a faint halo. The longstanding mystery of why the corona is significantly hotter than atmospheric layers closer to the solar surface still remains unresolved (Sigalotti et al. 2023).

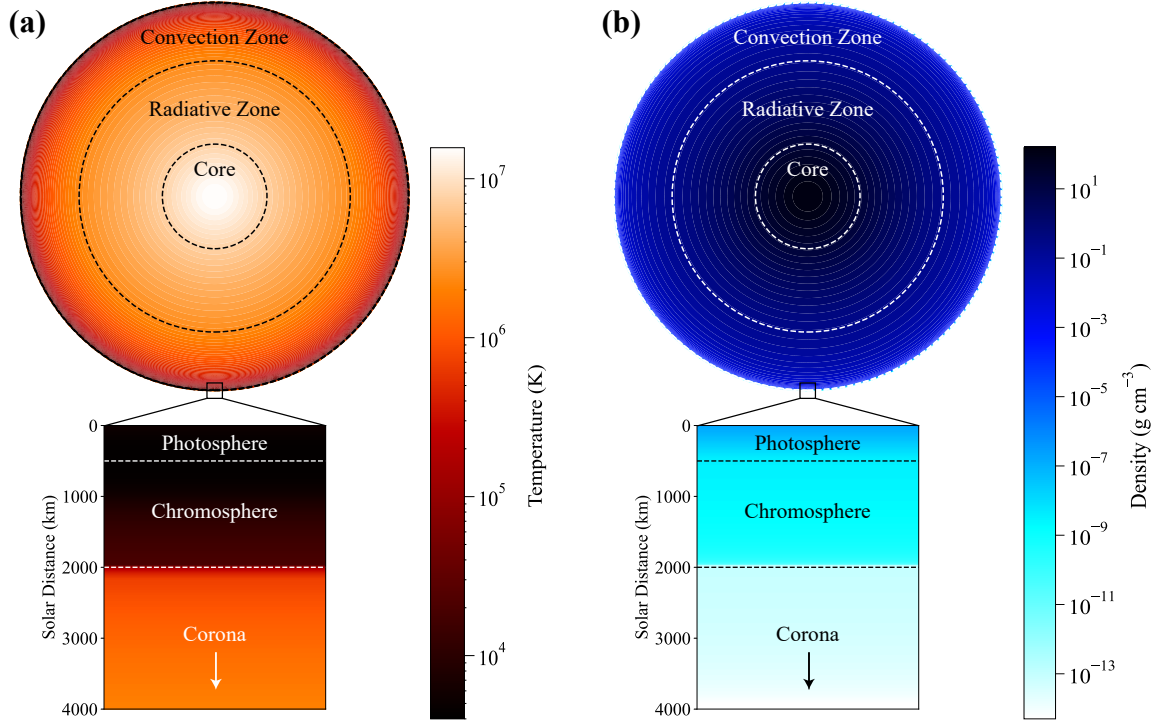


Figure 1.1: Simplified solar structure from the Sun's interior to atmosphere, where the (a) temperature (K) and (b) plasma density (g cm^{-3}) gradients are shown for each region. Solar models sourced from Christensen-Dalsgaard et al. (1996) and Lang (2001).

The Sun also generates a global magnetic field via hydromagnetic dynamo processes in its interior, where convecting solar material causes magnetic flux tubes to rise toward the surface (Aschwanden 2014). In the photosphere, strong and concentrated magnetic fields form sunspots, reaching intensities of several kilogauss. Sunspots tend to be cooler than the surrounding areas as the intense fields inhibit convective motion, appearing as dark regions on the surface. Beyond the photosphere, the magnetic field strength decreases in the solar corona on the order of tens to hundreds of gauss, but becomes more complex forming structures that extend far into space. The coronal magnetic field is the primary driver of various observed features, such as streamers, plumes, and loops (Hood et al. 2011). In addition, magnetic reconnection can occur in the solar corona, where opposing magnetic field lines break and reconnect to release large amounts of energy, effectively converting magnetic

energy into kinetic energy (Parker 1963). This process can originate from arches of opposing field lines connecting two sunspots on the surface, corresponding to the footprints of the flux rope, and can create various phenomena like solar flares and coronal mass ejections (see Section 1.2).

Another property of the solar corona is that it is continuously expanding into interplanetary space because, as demonstrated by Parker (1958), the solar plasma could not be in hydrostatic equilibrium with the gravitational field of the Sun. This results in the solar wind, a continuous stream of primarily electrons and protons, with a lower abundance of alpha particles and heavier ions. The solar wind can be categorized into fast ($v \gtrsim 400 \text{ km s}^{-1}$) and slow ($v \lesssim 400 \text{ km s}^{-1}$) streams, which emanate from different solar sources. Open field lines originating from coronal holes into the solar surface and extending toward the interplanetary medium are the source of the fast solar wind. Instead, slow solar wind sources are more complex and are theorized to originate from multiple phenomena, such as solar transient releases from helmet streamers, accelerated plasma from rapidly expanding magnetic fields near coronal hole boundaries, and/or continual plasma exchanges between loops and open magnetic fields (Rouillard et al. 2021). When these streams interact, they can form structures known as stream interaction regions (SIRs), which can generate interplanetary shocks.

The differential rotation of the Sun, which is theorized to also influence the magnetic dynamo, causes the solar wind to expand outward in a Parker spiral formation instead of a radial path from the source (Aschwanden 2014). The solar angular velocity ω is roughly $2.9 \times 10^{-6} \text{ rad s}^{-1}$ near the equator, with a slower speed closer to the poles. Due to the high electrical conductivity of the solar wind, it carries the solar magnetic field which is “frozen” into the plasma as it expands, forming the interplanetary magnetic field (IMF). The IMF also forms a Parker spiral, whose shape is dependent on the heliocentric distance r and solar wind speed v_{SW} . Based on the equations derived by Parker (1958), the magnetic field $\mathbf{B}$ at a given point (r, θ, ϕ) in a Sun-centered spherical coordinate system is

$$\mathbf{B}_r(r, \phi, \theta) = \mathbf{B}(r_0, \phi_0, \theta) \left(\frac{r_0}{r} \right)^2 \quad (1.1)$$

$$\mathbf{B}_\phi(r, \phi, \theta) = -\mathbf{B}(r_0, \phi_0, \theta) \left(\frac{\omega r_0^2}{v_{SW} r} \right) \sin(\theta) \quad (1.2)$$

$$\mathbf{B}_\theta = 0 \quad (1.3)$$

where $\mathbf{B}_r$, $\mathbf{B}_\phi$, and $\mathbf{B}_\theta$ represent the radial, longitudinal, and latitudinal components of the magnetic field, respectively. We define ϕ_0 as a given longitude at a reference distance r_0 from the Sun. This corresponds to the spiral magnetic field line equation given by

$$\phi = \phi_0 - \left[\frac{(r - r_0) \omega \cos(\theta)}{v_{SW}} \right] \quad (1.4)$$

where we have accounted for the counter-clockwise rotation of the Sun. Figure 1.2 portrays the IMF in the equatorial plane for the slow and fast solar wind streams. Near Earth at

1 AU (astronomical unit), the average IMF line is inclined roughly 45° in the equatorial plane for $v_{SW} \sim 400 \text{ km s}^{-1}$. For $|\theta| > 0^\circ$, the IMF forms helices wrapped around cones of constant latitude, approaching radial lines at the poles. Due to the expansion of open IMF lines oriented in opposite directions from the Sun, a sector boundary is formed, known as the Heliospheric Current Sheet (HCS), that divides regions of opposing polarity (Hoeksema et al. 1983). The regional extent of the solar wind and IMF is known as the heliosphere, encompassing the entire solar system that shapes the interplanetary environment.

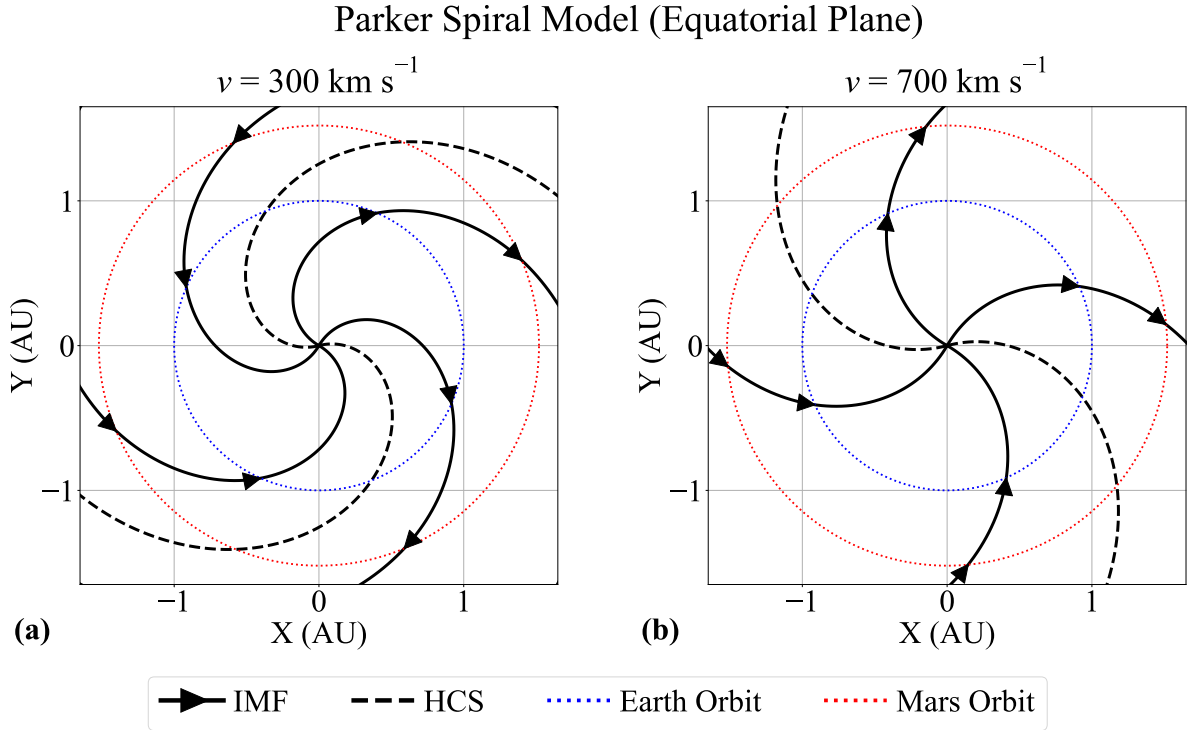


Figure 1.2: Parker spiral model of the interplanetary magnetic field (IMF) lines (black) computed for **(a)** slow ($v = 300 \text{ km s}^{-1}$) and **(b)** fast ($v = 700 \text{ km s}^{-1}$) solar wind streams. The black arrows indicate the direction of the IMF, where the Heliospheric Current Sheet (dashed black line) denotes a sector boundary between opposing IMF directions. Circular orbits at heliocentric distances to Earth and Mars are shown by the blue and red dotted lines, respectively.

Solar Wind Distributions

The first direct in situ measurements of the solar wind were obtained by a cylindrical electrostatic analyzer (ESA) on board Mariner 2, the first NASA space mission to perform a planetary encounter (Neugebauer et al. 1962). Since that discovery, numerous missions have

studied the solar wind with ESAs and other plasma instruments throughout the heliosphere. We can apply a statistical approach using kinetic theory by considering the plasma as an ensemble of particles to derive various macroscopic properties from these solar wind measurements (Kittel 2004). These parameters include the density n , temperature T , and pressure P of the plasma ensemble. For an individual particle in the ensemble, we define its phase space as the multidimensional space representing all possible states of the particle's position $\mathbf{r}$ and momentum $\mathbf{p}$. For a plasma consisting of N particles in 3D space, the phase space density $f(\mathbf{r}, \mathbf{p}, t)$, also known as the distribution function, defines the number of particles within a small volume element $d^3\mathbf{r}d^3\mathbf{p}$ at time t . By calculating moments of the distribution function $f(\mathbf{r}, \mathbf{v}, t)$ in velocity space, i.e., integrating the distribution function over velocity space $d^3\mathbf{v}$, one can derive the macroscopic properties of the plasma.

For a plasma in thermal equilibrium with no external potentials, the velocity distribution function (VDF) for particle species s can be described as a Maxwellian distribution given by

$$f(\mathbf{v}) = n_s \left(\frac{m_s}{2\pi k_B T_s} \right)^{3/2} \exp \left(-\frac{m_s(\mathbf{v} - \mathbf{u})^2}{2k_B T_s} \right) \quad (1.5)$$

where n_s is the species number density, m_s is the particle species mass, k_B is the Boltzmann constant, T_s is the isotropic plasma species temperature, and $\mathbf{u}$ is the plasma flow velocity (Gurnett et al. 2017). However, particle species in the solar wind are not always in thermal equilibrium to warrant the use of a Maxwellian distribution (see below). We can instead use other mathematical forms of the distribution function to attempt to model the kinetic physics of the solar wind. For example, in the presence of a magnetic field, a bi-Maxwellian function may provide a more accurate representation of the plasma, defined as

$$f(\mathbf{v}) = \frac{n_s}{T_{s\perp} \sqrt{T_{s\parallel}}} \left(\frac{m_s}{2\pi k_B} \right)^{3/2} \exp \left(-\frac{m_s(v_{\parallel} - u_{\parallel})^2}{2k_B T_{s\parallel}} \right) \exp \left(-\frac{m_s(v_{\perp} - u_{\perp})^2}{2k_B T_{s\perp}} \right) \quad (1.6)$$

where the components $\parallel$ and $\perp$ are parallel and perpendicular to the magnetic field direction, respectively. This distribution function is both anisotropic and gyrotropic, where $f(\mathbf{v})$ does not depend on the phase angle of the gyromotion around the magnetic field direction. In the case that $u_{\parallel}$ or $u_{\perp}$ vanishes because the plasma drifts either perpendicular or parallel to the field direction, the resulting distribution function is known as a drifting Maxwellian in that given direction (Kallenrode 2013). For higher energies, certain particle species are better modeled by a bi-Kappa (or bi-Lorentzian) distribution (Summers et al. 1991), given by

$$f(\mathbf{v}) = \frac{n_s}{v_{\kappa\parallel} v_{\kappa\perp}^2} \left(\frac{1}{\pi\kappa} \right)^{3/2} \frac{\Gamma(\kappa+1)}{\Gamma(\kappa-1/2)} \left[1 + \frac{(v_{\parallel} - u_{\parallel})^2}{\kappa v_{\kappa\parallel}^2} + \frac{(v_{\perp} - u_{\perp})^2}{\kappa v_{\kappa\perp}^2} \right]^{-\kappa-1} \quad (1.7)$$

where $\Gamma(z) = \int_0^\infty s^{z-1} e^{-s} ds$ is the Gamma function for an integration variable s and the components of the kappa thermal speed v_κ are defined as

$$v_{\kappa\parallel} = \sqrt{\left(\frac{\kappa - 3/2}{\kappa}\right) \frac{2k_B T_{s\parallel}}{m_s}} \quad v_{\kappa\perp} = \sqrt{\left(\frac{\kappa - 3/2}{\kappa}\right) \frac{2k_B T_{s\perp}}{m_s}} \quad (1.8)$$

for $\kappa > 3/2$, where κ is the spectral index that quantifies the deviation from a Maxwellian distribution. For high energies, the distribution is approximated by an inverse power law of the form $f \propto E^{-\kappa-1}$ instead of the exponential decay from a Maxwellian function. As κ approaches infinity, the distribution function converges toward a Maxwellian. Thus, κ characterizes the slope of the energy spectrum for the suprathermal component of the VDF.

Solar wind protons and alpha particles cannot be sufficiently modeled with a simple isotropic Maxwellian function, in part due to weak collisionality. A bi-Maxwellian VDF provides a better approximation, especially as the IMF establishes a preferred direction in space. Anisotropic proton temperatures in the fast wind tend to follow $T_{p\perp} > T_{p\parallel}$, while the opposite holds true in slow solar wind streams (Marsch 2006). For both types of streams, non-Maxwellian features are observed such as field-aligned proton beams with drift speeds comparable to the local Alfvén speed, defined as $v_A = B/\sqrt{4\pi n_s m_s}$. This may be an indication that the VDF is strongly affected by ion beam streaming instabilities through wave-particle interactions (Gosling 2014).

Solar wind electron VDFs can be divided into three different populations (Feldman et al. 1975; Pilipp et al. 1987), shown in Figure 1.3. The cold thermal population, referred to as core electrons, is nearly isotropic and can be modeled with a bi-Maxwellian distribution function. The ratio between parallel and perpendicular core temperatures is close to 1.1 at 1 AU, with higher ratios for lower density intervals (Gosling 2014). The core population makes up roughly 95% of the total electron density in the solar wind at 1 AU (Maksimovic et al. 2005; Štverák et al. 2009). The remaining suprathermal portion can be separated into a halo and strahl population. Halo electrons are more energetic than the core population and are best modeled by a Kappa distribution function (Maksimovic et al. 1997; Scudder 1992). The strahl population is composed of magnetic field aligned electrons propagating away from the Sun, carrying the majority of the solar wind heat flux (Pilipp et al. 1987; Scudder et al. 1979b).

Previous studies have shown that the solar wind electron VDF evolves with distance from the Sun (Berčič et al. 2019; Maksimovic et al. 2005; Štverák et al. 2009), suggesting that initially free-streaming strahl electrons could form the halo as they expand from the solar corona (Berčič et al. 2020, 2021; Boldyrev et al. 2019; Landi et al. 2012). The width of the strahl beam, defined by the pitch angle distribution at a given energy, can change under different dominating kinetic effects as the solar wind propagates through the heliosphere. Strahl electrons are adiabatically focused by a spatially weakening magnetic field as they move anti-sunward due to the conservation of the first adiabatic invariant. Conversely, the beam width of these suprathermal electrons can be broadened as a result of electron pitch-angle scattering mechanisms. These scattering processes include Coulomb collisions (Berčič

et al. 2021; Boldyrev et al. 2019; Scudder et al. 1979b; Štverák et al. 2008) and particle instabilities (Landi et al. 2012; Štverák et al. 2008; Vasko et al. 2019), increasing the angular width of the strahl beam. A combination of these mechanisms could control the strahl width at any given time in the solar wind, as the dominance of each scattering process could vary based on solar wind conditions and heliocentric distance. Recent heliospheric missions, such as Parker Solar Probe (PSP; Fox et al. 2016), could provide evidence for the dominating scattering mechanisms closer to the Sun within 0.3 AU.

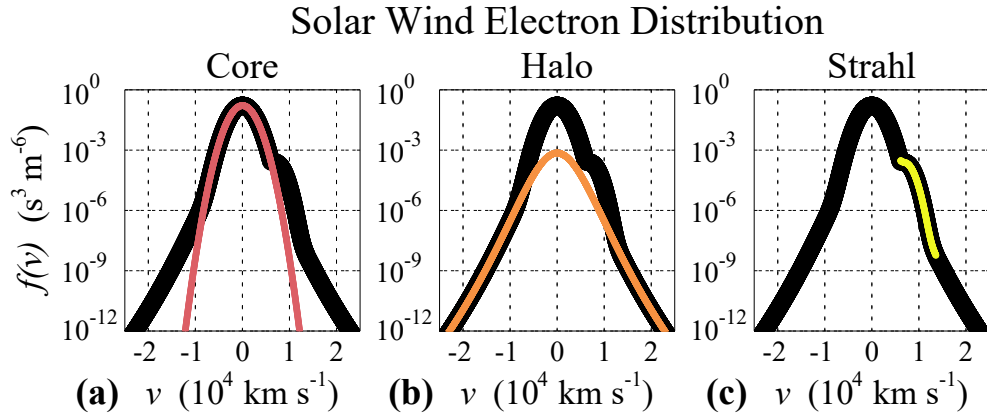


Figure 1.3: Illustration of a typical solar wind electron distribution function $f(v)$ for velocity v , composed of the (a) core (red), (b) halo (orange), and (c) strahl (yellow) populations.

Solar Periodicity and Evolution

The Sun exhibits variability across different temporal cycles, linked to the structure and dynamics of solar magnetic fields. For example, the number of sunspots on the Sun has a periodicity of ~ 11 years. The average sunspot number S_n is associated with solar activity and has been demonstrated to correlate with parameters such as the total surface magnetic flux and total solar irradiance. The period of intense solar activity is referred to as solar maximum, whereas quiet Sun conditions are observed during solar minimum. Enhanced solar activity also corresponds to an increase in the occurrence rate of solar eruptive events (see Section 1.2). Figure 1.4a portrays the annual sunspot number from 1700 to 2024 obtained from various solar observatories and compiled by the SILSO World Data Center (2024). Each cycle can be divided into the growth phase toward solar maximum and the decline phase toward solar minimum. Figures 1.4b through 1.4d show the power spectral density (PSD) of yearly sunspot number for the growth, decline, and combined phases. The growth phase can last 2 to 5 years, while the decline phase can last between 5 and 12 years (Miroshnichenko 2023). However, the combined power spectrum of both phases has a clear peak near 11 years. The solar cycles have been numbered since 1755, with solar cycle 25 currently

approaching solar maximum. The orientation of the solar magnetic field also changes every two solar cycles, with polar reversals occurring every 22 years, known as the Hale cycle. The solar magnetic field resembles a dipole during solar minimum, while the field becomes more complex as the polarity reverses during solar maximum (Moldwin 2022).

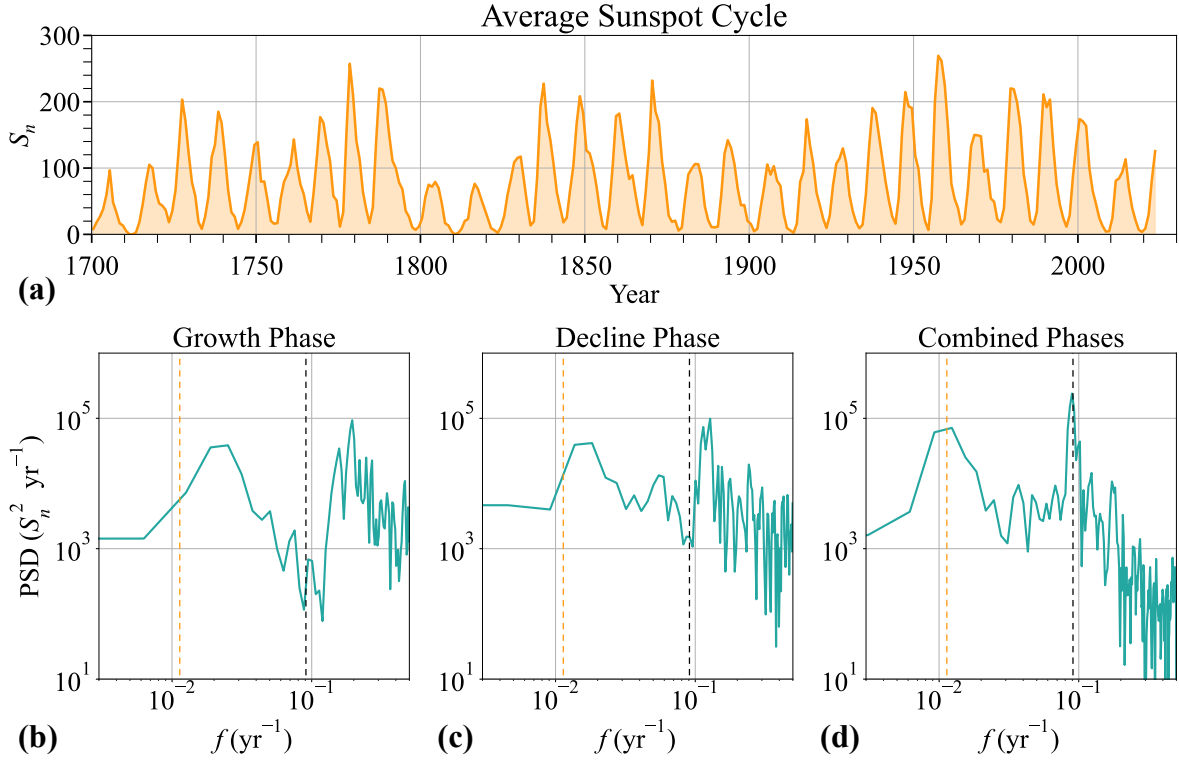


Figure 1.4: **(a)** Annual average total sunspot number (S_n) from 1700 to 2024. The sunspot numbers before 1749 were reconstructed from indirect geophysical data. The power spectral density (PSD; teal) is computed for the **(b)** growth phase, **(c)** decline phase, and **(d)** combined phases of the cycles. Two spectral peaks are found in (d) corresponding to a period of 88 years (vertical orange dashed line) and 11 years (vertical black dashed line). Data provided by SILSO World Data Center (1700–2024), Royal Observatory of Belgium, Brussels.

Figure 1.4 reveals a longer cycle that oscillates roughly every 88 years, indicated by the vertical orange dashed line. This is known as the Gleissberg cycle, with a period between 80 and 130 years, and is related to the amplitude variation of the 11 year cycle (Gleissberg 1958). There have been periods of weak solar cycles, such as the Maunder minimum (1645 to 1715) which has been correlated to the “Little Ice Age” in Europe. Longer cycles have been discovered using dendrochronology with measurements of radiocarbon (^{14}C) concentration in tree rings. Enhanced solar activity is more effective in shielding cosmic rays from Earth that can bombard atmospheric nitrogen to create ^{14}C . Therefore, solar maximum periods are

correlated to low intervals in ^{14}C production, with a 20 to 60 year delay due to the residence time of ^{14}C in the atmosphere-ocean system (Stuiver et al. 1980). Solar activity has also been inferred to vary on time scales of about 200 years (known as the Suess cycle) and 2000 years (Damon et al. 1991; Sonett et al. 1990).

Beyond humanity's existence, the Sun has evolved on time scales on the order of billions of years. Compared to the present Sun, the solar luminosity was roughly 75% when life began on Earth (~ 3.7 Ga) based on the Yonsei-Yale theoretical model (Kim et al. 2002). In addition, younger stars (~ 0.1 Gyr in age) rotate more rapidly than our Sun currently by an order of magnitude (Güdel et al. 1997). Due to the smaller rotation period, a young Sun would produce a stronger magnetic dynamo, corresponding to enhanced high-energy emissions. Older stars tend to have slower rotation speeds due to the loss of angular momentum from the release of magnetized solar wind (Durney 1972). The strengths of FUV, UV, and XUV (EUV and X-ray) emission for the young Sun are theorized to be stronger than the present emission of the Sun by a factor of 5–10, 10–100, and 100–1000, respectively (Ribas et al. 2005). Figure 1.5 shows the solar evolution of XUV flux, luminosity and rotation period normalized to current solar conditions as a function of stellar age. Solar XUV flux, characterized by wavelengths between 0.1 and 120 nm, was roughly 7 times higher when life first appeared on Earth. Solar wind intensity was also higher for a younger Sun by a factor of 10 to 1000, but future work is required to better resolve these estimates (Ribas 2009).

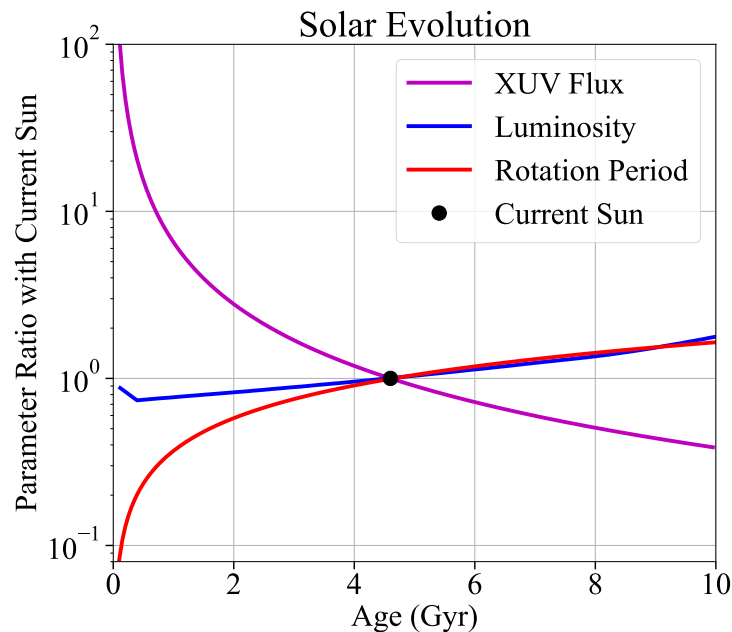


Figure 1.5: Solar evolution as a function of time. The XUV flux (0.1–120 nm; purple), luminosity (blue), and rotation period (red) are normalized to current solar conditions. Adapted from Ribas (2009).

1.2 Solar Eruptive Phenomena

Coronal Mass Ejections

Solar eruptive events, such as those produced by magnetic reconnection, can expel large amounts of energy and mass. An example of such an event is a Coronal Mass Ejection (CME), which releases magnetized plasma from the corona (Miroshnichenko 2023). CMEs are typically associated with active regions on the Sun, due to the concentration of strong magnetic fields. Once the trigger mechanism initiates the release of accumulated magnetic energy, the resulting magnetic instability rapidly grows, leading to an outward eruption accelerated from the Sun (Thompson et al. 2021).

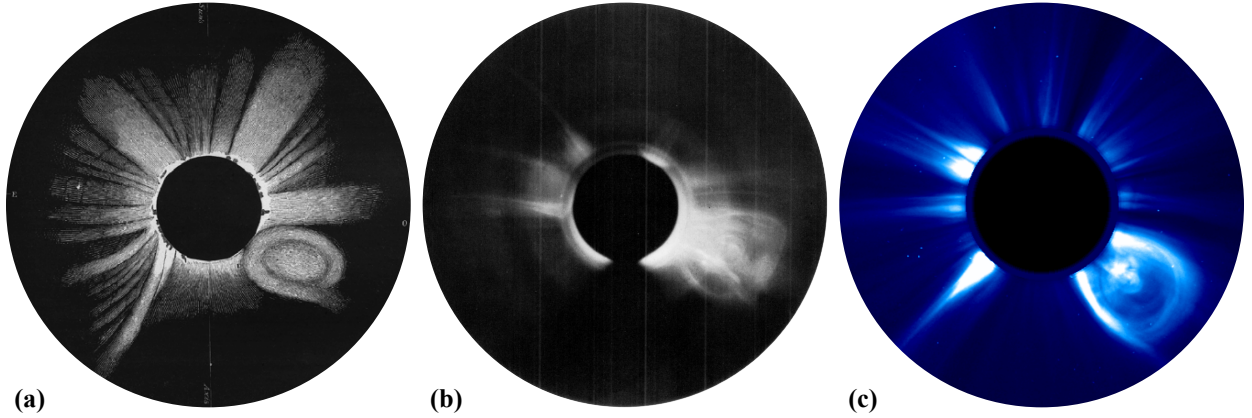


Figure 1.6: Coronal Mass Ejection (CME) observations over 150 years. (a) Illustration by Guglielmo Tempel of the solar corona in Spain during the 18 July 1860 total solar eclipse with a possible CME. (b) White-light coronagraph observation from the Skylab mission of a CME on 10 June 1973. (c) LASCO C2 image of a CME on 27 January 2012, with a classical three-part structure. Reproduced from Eddy (1974) and the SOHO LASCO CME catalog.

One of the first documented observations of a CME was recorded by Italian Astronomer Guglielmo Tempel in Torreblanca, Spain during the July 1860 total solar eclipse. However, it was over 100 years later that Eddy (1974) compared Tempel’s illustration with white-light coronagraph observations from the Skylab mission (MacQueen et al. 1974) to conclude that Tempel most likely witnessed a CME (see Figure 1.6). Numerous space missions with coronagraphs have been deployed since this discovery, such as the Large Angle Spectroscopic Coronagraph (LASCO; Brueckner et al. 1995) on board the Solar and Heliospheric Observatory (SOHO; Domingo et al. 1995). Figure 1.6 shows a LASCO C2 observation of a CME on 27 January 2012, following the classical three-part structure (Illing et al. 1986). This model is defined to have a bright frontal shell of coronal plasma, a darker inner cavity enclosing the magnetic flux rope, and a bright central core containing the remnants of the eruptive

prominence. It is important to note that this structure is not always observed for more complex CMEs, and can depend on other factors such as the observer’s perspective relative to the event.

Once the CME propagates further into space after most of its acceleration, it is referred to as an Interplanetary Coronal Mass Ejection (ICME). These structures typically move outward from the Sun with speeds ranging from 50 to 2000 km s⁻¹, and in some cases beyond 3000 km s⁻¹ (Thompson et al. 2021). CMEs typically carry between 10¹³ and 10¹⁶ g of solar matter, requiring roughly 10²⁷–10³² erg of energy to overcome the gravitational attraction of the Sun (Kataoka 2022). Based on their ejected mass and occurrence rate, CMEs can contribute up to 10% of the total mass loss from the corona by the solar wind (Stix 2002).

Solar Flares

Solar flares are intense bursts of radiation that can emit electromagnetic energy in almost all observable wavelengths (see Figure 1.7), persisting from seconds to hours (Benz 2017). Similar to CMEs, flares require a trigger mechanism to release magnetic energy and tend to occur near active regions. It is important to note that solar flares do not always coincide with the generation of CMEs and can occur independently based on the nature of the energy release from the Sun. Different types of flares can vary significantly in terms of scale, duration, and emission spectrum (Blackman 1997). Compact or confined flares typically last on the order of minutes, with luminosity values near $\sim 10^{28}$ ergs s⁻¹ and length scales of less than 10³ km. In general, these events are more impulsive and can create bursts of non-thermal electron emission. The second type of solar flare can continue for hours with similar luminosity values, known as long-duration, gradual flares. These eruptions create little non-thermal emission and can span length scales exceeding 10⁵ km and, in some cases, more than one solar radius. This wide range of scales motivates the development of improved instrument resolution and sensitivity, such as the instrumentation on the FOXSI-4 sounding rocket mission (Buitrago-Casas et al. 2021), to further study finer structures.

The duration of a flare can be divided into four different phases, as depicted in Figure 1.7 (Benz 2017). The preflare phase (1-3 minutes) begins with heating of the coronal plasma in soft X-rays and EUV, with the occasional occurrence of non-thermal emission. Soft X-ray precursors are transient brightness enhancements that often occur in loops or unresolved kernels, whereas UV precursors are mainly small-scale transients above active regions, exhibiting a wide range of amplitudes (Su 2010). Energetic electrons and ions are accelerated during the impulsive phase (3-10 minutes), when the majority of energy is released. All wavelengths exhibit an increase in intensity, with substantial hard X-ray emission from foot-point sources in the chromosphere. Soft X-ray, EUV, and H α emission peak in intensity during the flash phase (5-20 minutes) of the flare. Finally, the plasma in the lower corona cools in the decay or gradual phase (>1 hour), with energy transported to other regions of the atmosphere, leading to continued acceleration of particles by various processes such as plasma ejections or shock waves (Su 2010).

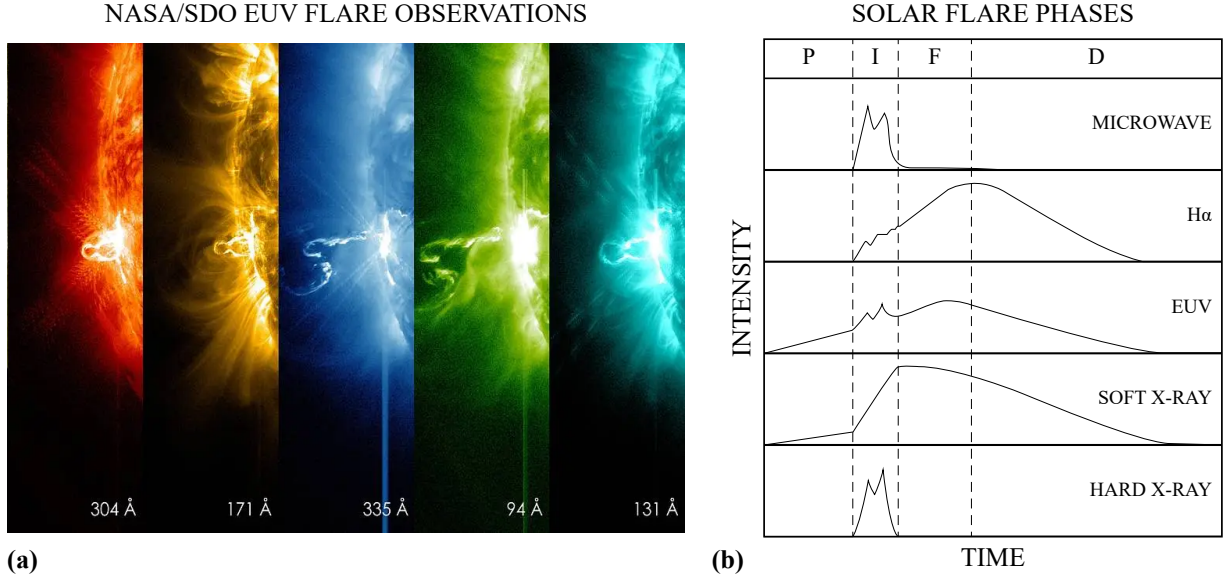


Figure 1.7: **(a)** NASA Solar Dynamics Observatory (SDO) observations of the 24 February 2014 X-class solar flare in EUV emission from 94 to 304 Å. Image courtesy of NASA/SDO. **(b)** A schematic illustration of typical solar flare intensity over time at different wavelengths. The flare phases are depicted at the top of the panel, labeled as preflare (P), impulsive (I), flash (F), and decay (D) phases. Adapted from Benz (2017).

Solar flares can be further divided into different categories based on their X-ray intensity in the wavelength range between 1 and 8 Å (Thompson et al. 2021). The base magnitude values of peak soft X-ray flux for flare classes A, B, C, M, and X are 10^{-8} , 10^{-7} , 10^{-6} , 10^{-5} , and 10^{-4} W m $^{-2}$, respectively. Each class can be further separated into numerical divisions ranging from 1.0 to 9.9, representing the multiplicative factor to the order of magnitude flux (e.g., a X4.2 flare has a flare strength of 4.2×10^{-4} W m $^{-2}$). It is possible to achieve intensities above X10, such as the famous 01 September 1859 Carrington solar flare with an estimated strength of $X45 \pm 5$ (Cliver et al. 2013). The Geostationary Operational Environmental Satellite (GOES) system, operated by NOAA, continuously monitors solar activity in near real-time to classify solar flares.

1.3 Solar Impacts on Planetary Environments

Planetary Magnetospheres

Once the solar wind reaches Earth, it interacts with the global geomagnetic field to form a cavity in which the strength of the planetary magnetic field dominates the solar wind pressure (Russell 1991). The intrinsic field intensity of the planet is ~ 30 μ T at the equator

and roughly $50 \mu\text{T}$ near the poles (Pulkkinen 2007), whereas the IMF strength is about $0.1 \mu\text{T}$ at 1 AU. Figure 1.8a shows the various components of the Earth’s magnetosphere that result from this dynamic. On the dayside (sunward side), a bow shock forms due to the supersonic solar wind encountering the magnetosphere which acts as an obstacle, creating a standing shock wave. This boundary heats and compresses the solar wind plasma as it is deflected around the magnetosphere (Borovsky et al. 2018). The region of shocked solar wind, known as the magnetosheath, contains plasma that is denser and hotter than the ambient solar wind. Another threshold forms once the pressure between the solar wind and the geomagnetic field is balanced, referred to as the magnetopause. Other distinct regions exist within this boundary, such as the radiation belts, the ring current, and the plasmasphere, each characterized by particle energies on the order of MeV, keV, and eV, respectively. On the nightside, the magnetosphere is stretched to develop the magnetotail (extending over hundreds of earth radii), a reservoir of magnetic energy that drives numerous magnetospheric processes (Borovsky et al. 2018). Polar cusps form near the magnetic poles, where the geomagnetic field direction diverges from the dayside to the nightside, enabling solar wind particles to funnel directly into the upper atmosphere. In addition, the upper atmospheric layer of the Earth, known as the thermosphere, is subject to ionization from solar radiation and energetic particles from the magnetosphere, forming the ionosphere.

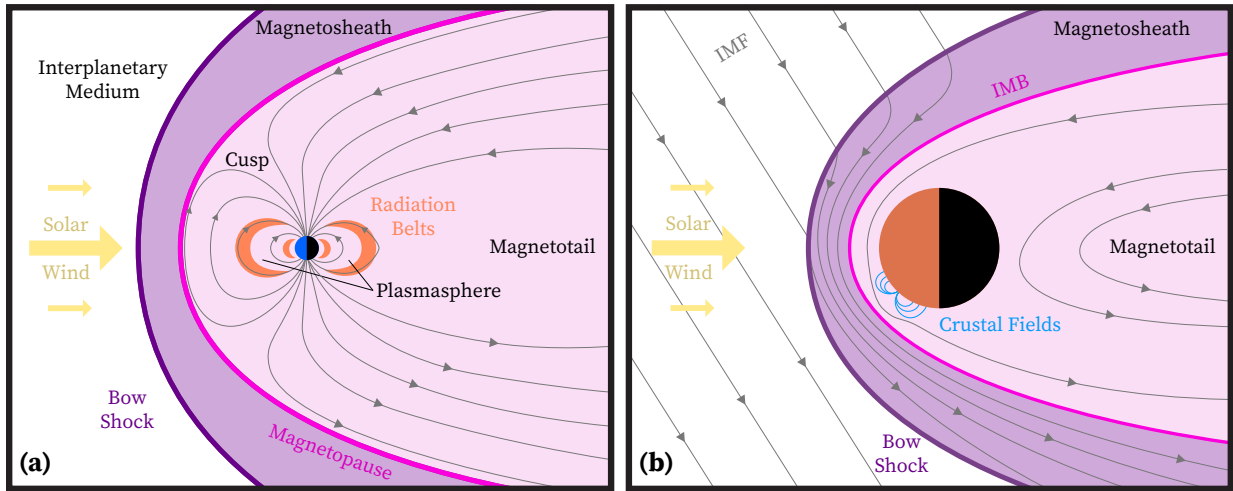


Figure 1.8: Planetary magnetospheres for (a) Earth and (b) Mars (induced). The solar wind (yellow) propagates left to right, carrying the interplanetary magnetic field (grey). Key features include the bow shock (dark purple), magnetosheath (light purple), magnetopause (dark pink) for Earth, induced magnetosphere boundary (IMB; dark pink) for Mars, and the magnetotail. The inner magnetosphere of Earth contains the radiation belts (orange), the polar cusps, and the plasmasphere. Although Mars lacks an intrinsic geodynamo, strong crustal fields (blue) exist near the surface. The dimensions of the figure panels for Earth and Mars are approximately 40×50 earth radii and 8×10 mars radii, respectively.

The magnetosphere is crucial in protecting the Earth from solar energetic particles and radiation originating from the Sun. However, highly energetic solar events can still disrupt conditions in the geospace environment near Earth, known as space weather (Moldwin 2022). For example, CMEs can generate strong interplanetary shock waves and southward IMF to produce geomagnetic storms, which represent major disturbances of the Earth’s magnetosphere. X-class flares can also cause significant disruptions in the ionosphere, leading to radio blackouts and navigation system interference. Other effects of intense solar activity include satellite damage, astronaut safety risks, overloaded power systems, and auroral events on Earth. The March 1989 geomagnetic storm, triggered by a CME, caused widespread disruptions to satellite and communications systems, including a major power outage in Quebec, Canada.

In the absence of a planetary geomagnetic field, the solar wind can interact directly with the upper atmosphere of the planet and generate an induced magnetosphere (Bertucci et al. 2011; Nagy et al. 2004), shown in Figure 1.8b for Mars. Similar to Earth, a bow shock forms around Mars due to the supermagnetosonic solar wind flow. The compressed and decelerated plasma then fills the turbulent magnetosheath within the shock boundary. This region is separated by the induced magnetosphere boundary (IMB), also known as the magnetic pile-up boundary, with draped IMF lines. The induced magnetosphere arises from induced electric currents in the Martian ionosphere, which provide enough magnetic pressure to counteract the solar wind plasma (Ramstad et al. 2020). On the nightside, a magnetotail develops with two lobes that are divided by a plasma sheet. In addition, strong localized magnetic fields remanent from a preexisting geodynamo exist near the surface, which can interact with the solar wind to produce a “hybrid magnetosphere” (Acuña et al. 1998; Brain et al. 2003). These crustal fields can influence the motion of particles to determine if they remain within or escape from the Martian environment (Sánchez-Cano et al. 2022). As the distribution of the crustal fields is non-uniform on the surface, their effect is dependent on planetary rotation.

Space Weather Effects on Mars

Solar radiation and particle emissions were more intense for a young Sun, as described in Section 1.1, leading to a greater impact on planetary environments in the past. Mars is especially vulnerable to solar activity due to the lack of a protective geomagnetic field coupled with the small size of the planet, roughly half the diameter of Earth (Brain 2017). The effects of space weather on Mars include changes in the induced magnetosphere structure, auroral events, atmospheric heating, and atmospheric loss. A large fraction of the Martian atmosphere has been lost to space, which led to the transition from a warm and wet environment to the cold and dry climate of the planet today (Jakosky et al. 2017). It is theorized that an ancient water ocean once existed in the northern hemisphere of Mars, but was gradually lost over time, either through evaporation into the atmosphere or sequestration in the Martian soil (Parker et al. 2021; Wright et al. 2024). Atmospheric erosion can occur through various non-thermal processes, such as sputtering or ion pick-up mechanisms (Ribas 2009). This

issue is also relevant to the future of humanity on Mars, as it poses significant hazards to astronauts. For example, during the 2003 Halloween solar storm, the Mars Radiation Environment Experiment (MARIE) on board the Mars Odyssey orbiter was destroyed because of intense solar radiation (Gopalswamy et al. 2005). Although Mars is technically within the habitable zone with respect to the Sun, it may have been difficult for life to develop on the surface due to these conditions.

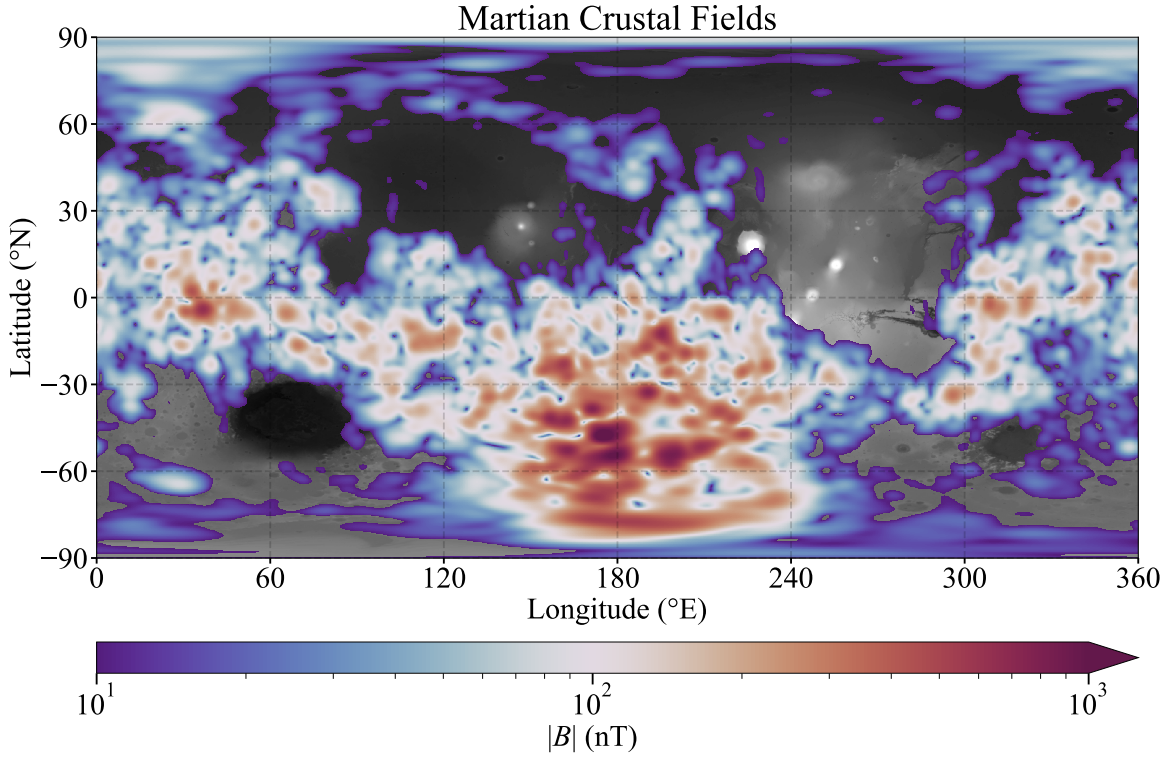


Figure 1.9: Magnetic field intensity ($|B|$) of crustal magnetism on Mars predicted 150 km above the surface by Langlais et al. (2019), displayed over the Mars Orbiter Laser Altimeter (MOLA) topography map (Smith et al. 2001). The model resolution is 0.25° and excludes crustal magnetic field values below 10 nT.

Mars most likely possessed an ancient geodynamo in the past that would have protected the planet from periods of intense solar activity. Stevenson (2010) describes three different processes that could have caused the loss of the dynamo. One possibility is that a reduction in the cooling of the Martian core occurred to the extent that conductive heat loss became the primary mechanism, leading to a subadiabatic, conductive planetary state that could not sustain a dynamo. Another theory is that Mars experienced a transition from an efficient convection mode, such as plate tectonics, to the current stagnant lid mode. This shift would halt the cooling of the mantle and core, leading to the cessation of core convection

and dynamo activity. A third possibility is that the Martian core solidified to a point where the liquid core was too thin to maintain a geodynamo. Of course, the dynamics of the Martian interior may be more complex than these mechanisms. However, the crustal magnetic fields near the surface may offer clues about Martian history. Figure 1.9 shows a global map of the remanent fields on Mars at an altitude of 150 km predicted by the Langlais et al. (2019) model based on orbital measurements. Stronger fields are observed in the southern hemisphere, with little to no magnetization in the northern hemisphere and near impact basins. Studying these crustal magnetic fields will offer valuable insights into planetary evolution and improve astronaut safety from solar radiation during future human exploration missions on the Martian surface.

1.4 Dissertation Outline

Each section of this dissertation analyzes the temporal and spatial evolution of a different portion of the heliosphere, utilizing observations from novel instrumentation. In chapter 2, we investigate the dominating mechanisms that scatter the solar wind electron strahl close to the Sun using Parker Solar Probe in situ measurements. In addition, we find the radial evolution and governing solar wind conditions of the electron strahl beam width. In Chapter 3, we present novel in situ and remote-sensing measurements from Parker Solar Probe of a Carrington-level ICME event on 05 September 2022 (the “Labor Day” ICME). We compare this event with typical ICME signatures at 1 AU near Earth, and examine its role in the global reconfiguration of the HCS over time. In Chapter 4, we utilize new surface magnetometer measurements to determine spatial scales of the crustal magnetization on Mars. We relate different distributions of magnetization to geological regions of different ages, deriving a possible chronology of the Martian dynamo responsible for an ancient global intrinsic magnetic field. Our conclusions are presented in Chapter 5, where we summarize the main results of each section and connect them within the framework of Heliophysics.

Chapter 2

Radial Evolution of Solar Wind Strahl Electrons Observed by PSP

2.1 Abstract

Due to their anisotropic and non-thermal nature, strahl electrons significantly contribute to the solar wind heat flux, which plays a role in the energy budget of the Sun. The strahl electrons form a beam distribution, observed by in situ measurements of the solar wind electron velocity distribution functions (eVDFs). The strahl distribution can adiabatically focus due to the diverging interplanetary magnetic field (IMF) or defocus due to various processes, such as Coulomb collisions or plasma wave instabilities. However, these observations have been previously limited to heliocentric distances above $62 R_S$ (solar radii) from the Sun. With the Parker Solar Probe (PSP) mission, we can measure eVDFs above roughly $10 R_S$ with the Solar Probe ANalyzer Electron (SPAN-E) instruments. In this chapter, we calibrate the SPAN-E particle detectors to determine bulk plasma properties, such as electron density and temperature. We then determine the electron strahl beam width and direction to investigate the evolution of these electrons from the Sun. We find the strahl beam narrows with increasing electron energy. In addition, the strahl width is observed to vary with heliocentric distance, heliospheric current sheet distance, electron density, and electron collisional age. Our study suggests that close to the Sun, magnetic focusing dominates Coulomb collisions to narrow the strahl beam less than $\sim 30 R_S$ from the Sun.

2.2 Introduction

The connection between the Sun and other bodies in the solar system is partially mediated by the solar wind, a fully ionized plasma mainly composed of protons and electrons that flow away from the solar corona (Parker 1958). The solar wind can interact with planetary magnetospheres, atmospheres, and surfaces throughout the heliosphere. In the case of Earth, extreme space weather events can impact natural systems of the planet and human

infrastructure, from geosynchronous satellites in orbit to power grids on the surface (Hapgood 2018). To better understand these interactions, we must first investigate fundamental processes that control the solar wind as it expands from the Sun. These include energy balance and acceleration mechanisms that occur in the near-Sun environment. Until recently, the Helios 2 mission (Porsche 1981) provided the closest in situ measurements to the Sun, maintaining the record for about four decades. The spacecraft collected solar wind data at a proximity of 0.29 AU (astronomical unit) or $62 R_S$ (solar radii) from the solar center, leaving roughly one-third of the distance between the Sun and Earth unexplored. However, recent technological advancements have enabled us to observe the solar wind closer to our star than ever before.

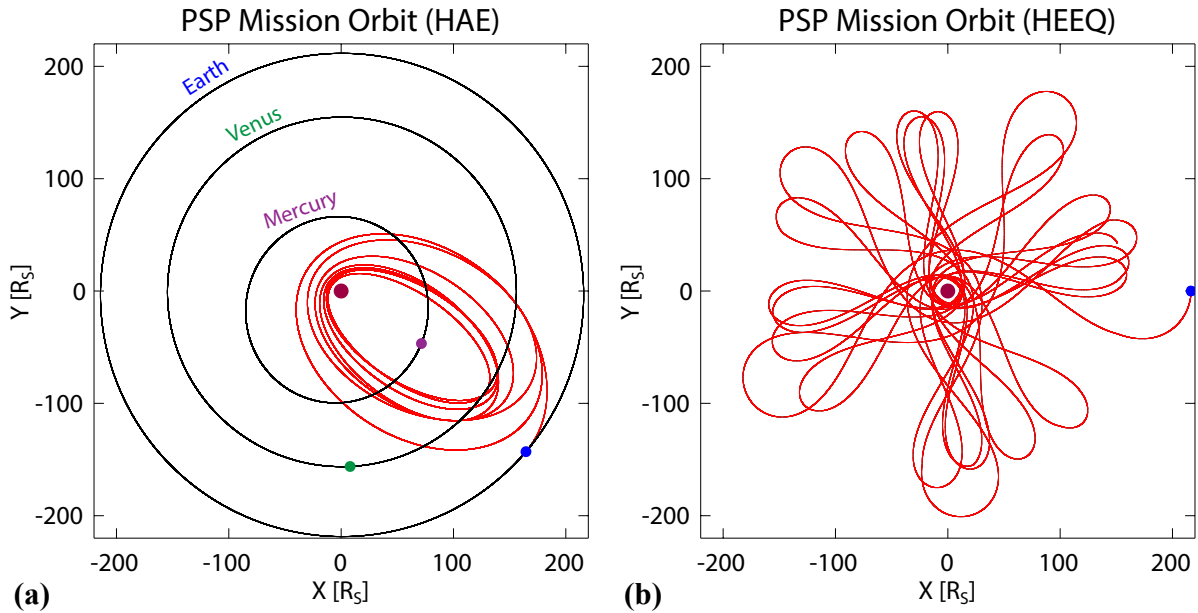


Figure 2.1: Orbital trajectory of Parker Solar Probe (red) from August 2018 to August 2025 projected in the (a) Heliocentric Aries Ecliptic (HAE) and (b) Heliocentric Earth equator (HEEQ) frames. The positions of the Sun (dark red), Mercury (purple), Venus (green), and Earth (blue) are shown at the time of the PSP launch.

The Parker Solar Probe mission (PSP; Fox et al. 2016) was launched on 12 August 2018 to directly measure the solar wind near the Sun. Figure 2.1 depicts the orbital trajectory of PSP in the Heliocentric Aries Ecliptic (HAE) and Heliocentric Earth Equator (HEEQ) frames from launch over the course of seven years. The HAE frame is centered on the Sun, with the x-axis pointed toward the first point in Aries (vernal equinox) and the z-axis toward the ecliptic north pole. The HEEQ frame, also known as the Stonyhurst Heliographic system, is solar centered with the x-axis defined by the intersection of the solar equator and central meridian as seen from Earth and the z-axis toward the solar north pole. During its

seven year primary mission phase, PSP will complete 24 orbital encounters with respect to the Sun. The spacecraft will undergo seven Venusian gravity assists to reduce its orbital perihelion from $35.7 R_S$ to within $10 R_S$, shown in Figure 2.1a. PSP is able to reach these distances due to the thermal protection shield (TPS) with a diameter of 2.4 m, made of 11.43 cm thick carbon-composite (C-C) material with a special white coating. The coating consists of tungsten and white alumina with magnesium oxide, which are air plasma sprayed onto the C-C material (Ercol et al. 2022). The tungsten serves as a barrier layer on the C-C, allowing PSP to withstand temperatures of roughly 1100°C .

With this capability, the mission can provide insight on fundamental plasma processes near the Sun. One of its scientific objectives is to study heating and acceleration of the solar wind from the low solar corona. We can trace the energy flow throughout the inner heliosphere by studying the non-thermal components of the solar wind particle distribution measured by PSP. Within the solar wind, protons transport the bulk of the momentum, while electrons carry the overall heat flux. This is a result of the larger mass of protons relative to electrons, which is also the source of the ambipolar electric field (Pannekoek 1922). The ambipolar electric field in the solar wind is generated by the relative motion of ions and electrons that exists to maintain charge neutrality of the plasma flow. Both the ambipolar electric field and electron heat flux are critical in regulating the energy balance between the coronal thermal energy and solar wind kinetic energy (Cranmer et al. 2009; Landi et al. 2014; Maksimovic et al. 1997; Scudder 1992; Štverák et al. 2015). This conversion of energy is theorized to drive solar wind acceleration in the heliosphere. We can study the non-thermal components related to the heat flux by examining the electron velocity distribution functions (eVDFs) observed by PSP.

As discussed in Chapter 1, the eVDFs are composed of three major populations: the nearly Maxwellian thermal core, the nearly isotropic and broad (non-Maxwellian) halo, and the suprathermal magnetic field aligned strahl (Feldman et al. 1975; Pilipp et al. 1987). Štverák et al. (2009) observed that the eVDFs evolve with distance from the Sun, shown in Figure 2.2. The relative density of the core electrons with respect to the total electron density remains mostly constant with heliocentric distance for both the slow and fast wind. Instead, as the strahl relative density decreases with distance, the halo density tends to increase for both types of solar wind. This may suggest that initially free-streaming strahl electrons expanding outward from the solar corona form the halo (Berčič et al. 2020, 2021; Boldyrev et al. 2019; Landi et al. 2012). While the relative densities of the strahl and halo change with heliocentric distance, these values are typically smaller than core densities by an order of magnitude. However, strahl and halo electron energies are comparatively much larger, allowing these electrons to travel farther through the heliosphere before achieving thermal equilibrium with the ambient solar wind and to carry more of the overall heat flux (Scudder et al. 1979a,b). In addition, particle heating and acceleration of the solar wind are affected by non-Maxwellian anisotropic portions of the eVDF, as they can provide a source of kinetic instabilities and small-scale turbulence (Forslund 1970). Therefore, due to both the asymmetrical and non-thermal nature of eVDFs, the strahl is the dominant contributor to the overall heat flux in the near-Sun environment.

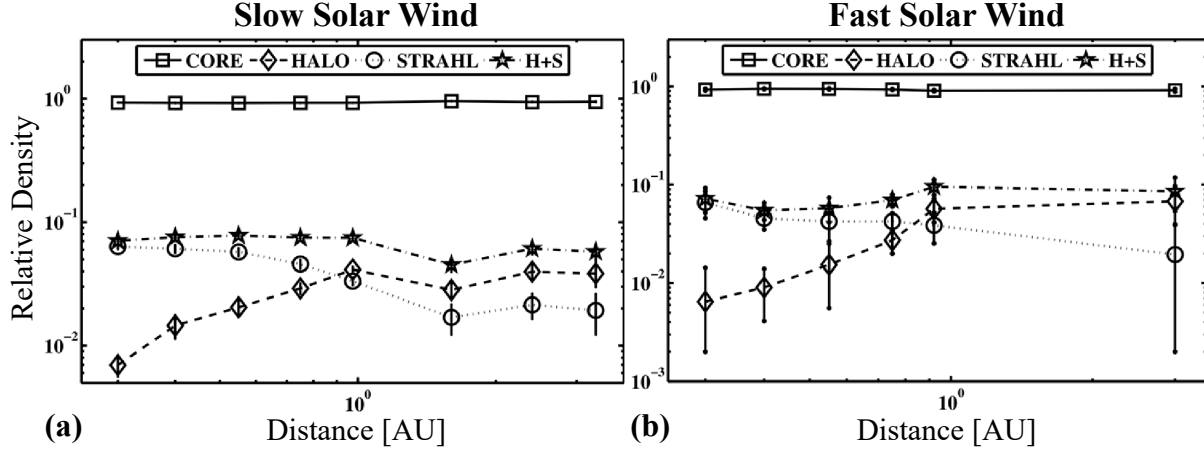


Figure 2.2: Relative electron density of the core (square), halo (diamond), and strahl (circle) populations as a function of radial distance from the Sun measured by the Helios mission for the (a) slow and (b) fast solar wind. While the strahl density decreases with distance below 1 AU, the halo population density increases. The sum of these non-thermal populations (star) appears to remain roughly constant with distance. Adapted from Štverák et al. (2009).

The strahl population is considered a runaway electron group originating from 5 to 15 R_S from the Sun. When referring to the strahl, it is convenient to define the pitch angle α as $\cos(\alpha) = (\mathbf{v} \cdot \mathbf{B}) / (|\mathbf{v}| |\mathbf{B}|)$, which quantifies the angle between the magnetic field direction and solar wind velocity. Figure 2.3 illustrates a typical electron distribution in 2D velocity space with respect to the magnetic field. Figure 2.3a shows the isotropic core and higher energy halo populations with a narrow beam of strahl electrons in the direction of the magnetic field. A wider strahl distribution is portrayed in Figure 2.3b, spread over a larger range of pitch angles. The shape of the strahl beam can be influenced by two dominating kinetic effects in the solar wind. The strahl can be adiabatically focused as the particles move outward into a spatially weakening magnetic field. This weakening occurs because magnetic field lines spread out as the solar wind expands from the Sun, following an approximate inverse square law (see equation 1.1 in Chapter 1). Due to the conservation of the first adiabatic invariant (the magnetic moment μ), the change in the magnetic field strength causes the particles to focus along the field direction. This can be demonstrated by equating the magnetic moment to the magnetic flux Φ_B of a particle gyrotropically rotating around the magnetic field, given by

$$\mu = \frac{|q|}{2} \omega_c \rho_c^2 = \frac{|q|}{2} \left[\frac{|q| |\mathbf{B}|}{m} \right] \rho_c^2 = \frac{q^2}{2\pi m} [\pi \mathbf{B} \rho_c^2] = \frac{q^2}{2\pi m} \Phi_B \quad (2.1)$$

where q and m are the charge and mass of the particle, ω_c and ρ_c are the particle cyclotron frequency and radius, and $\mathbf{B}$ is the magnetic field (Gurnett et al. 2017). The cyclotron

radius, or gyroradius, is the radius of circular motion of a charged particle orbiting around the magnetic field direction. The conservation of the magnetic moment signifies that a fixed magnetic flux is maintained throughout cyclotron orbits. Thus, if the magnetic field magnitude decreases, the cyclotron radius must also decrease to maintain constant magnetic fluxes outward from the Sun. Another method of demonstrating this conservation is to define μ in terms of the perpendicular kinetic energy $E_{\perp}$ and magnetic field B , given by

$$\mu = \frac{mv_{\perp}^2}{2B} = \frac{E_{\perp}}{B} \quad (2.2)$$

where the perpendicular kinetic energy and magnetic field strength must balance each other to ensure that the magnetic moment is conserved. As the magnetic field strength decreases with distance from the Sun, the perpendicular kinetic energy of the particle must also decrease with larger heliocentric distances.

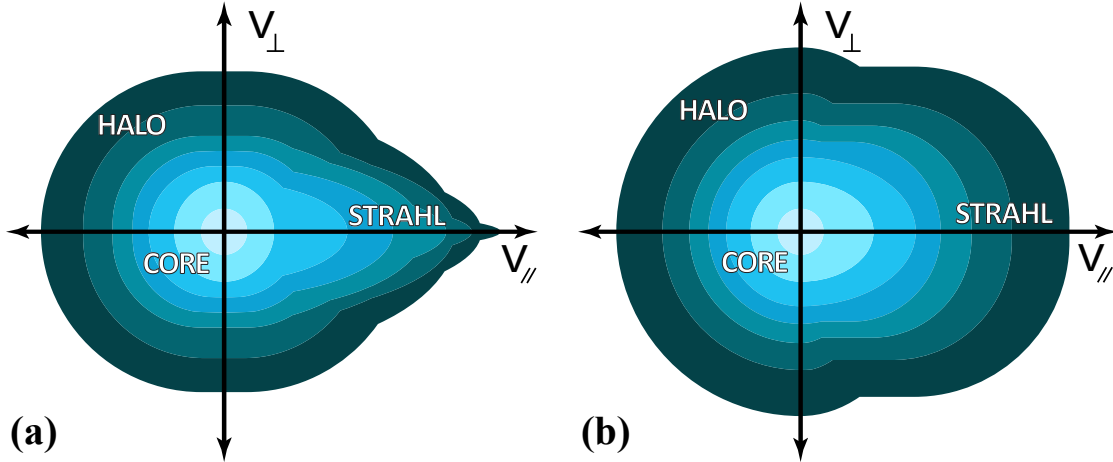


Figure 2.3: Illustration of a typical solar wind electron 2D velocity distribution, where $v_{\perp}$ and $v_{\parallel}$ are the velocities perpendicular and parallel to the magnetic field. **(a)** An example eVDF showing the core, halo and strahl populations when the strahl beam is narrow due to adiabatic focusing by the magnetic field. **(b)** Another eVDF with a broad strahl electron beam due to scattering mechanisms, possibly Coulomb collisions or wave scattering.

The strahl electrons can also broaden due to electron pitch angle scattering mechanisms such as Coulomb collisions (Berčič et al. 2021; Boldyrev et al. 2019; Horaites et al. 2018; Salem et al. 2003; Scudder et al. 1979b; Štverák et al. 2008). We can quantify the amount of collisionality in the solar wind by computing the electron collisional frequency $\nu_e = 2.91 \times 10^{-6} (n_e/T_e^{3/2}) \ln(\Lambda)$, where $\ln(\Lambda)$ is the Coulomb logarithm. Furthermore, we can calculate the electron collisional age $A_e \propto \nu_{e\perp} r / v_{SW}$ based on the radial distance r and solar wind speed v_{SW} (Salem et al. 2003). The collisional age quantifies the number of transverse collisions experienced by a thermal electron during the solar wind expansion over the scale

of the density gradient. Particle instabilities that lead to plasma waves can also scatter strahl electrons (Gary et al. 1975; Landi et al. 2012; Štverák et al. 2008; Vocks et al. 2005). Recent observations from PSP (Cattell et al. 2021; Halekas et al. 2021; Jagarlamudi et al. 2021) suggest oblique whistler wave modes (Horaites et al. 2018; Micera et al. 2020; Vasko et al. 2019; Verscharen et al. 2019) are the dominant plasma instability that scatters the electron strahl. However, recent work has indicated an absence of whistler waves for heliocentric distances below $28 R_S$, as well as for regions with certain values of electron plasma beta ($\beta_e < 0.1$ and $\beta_e > 10$) (Cattell et al. 2022). This implies that Coulomb collisions, or other plasma instabilities, could be the dominant scattering mechanism for solar wind acceleration near the Sun. This scattering could potentially depend on a combination of these processes, as the importance of each individual mechanism could vary based on solar wind conditions.

In this chapter, we investigate the electron scattering mechanisms below $30 R_S$ with PSP to better understand the fundamental processes that drive solar wind acceleration. First, we perform multiple calibration steps of the electron instrument on board PSP in Section 2.3 to obtain accurate estimates of solar wind electron parameters, such as densities and temperatures of the different eVDF populations. In Section 2.4, we fit to the strahl beam distribution to quantify the strahl scattering in terms of strahl width σ . The resulting relationships between strahl scattering and various solar wind parameters are shown in Section 2.5. The discussion and conclusions are presented in Sections 2.6 and 2.7, respectively.

2.3 Calibration of PSP SPAN-E Instruments

SPAN-E Instrument Functionality

The Solar Wind Electrons, Alphas, and Protons (SWEAP; Kasper et al. 2016) suite of instruments on board PSP include two electron electrostatic analyzers (ESAs), known collectively as the Solar Probe ANalyzer-Electron (SPAN-E; Whittlesey et al. 2020) instruments, to measure bulk plasma electrons below roughly 30 keV. SPAN-A is mounted on the ram side of the spacecraft bus and includes separate ion and electron ESAs: SPAN-I (Livi et al. 2022) for ions and SPAN-Ae for electrons. SPAN-B is a single electron ESA placed on the anti-ram side of PSP for maximum Field-of-View (FOV) coverage of the solar wind electrons. Figure 2.4 depicts the SPAN-A and SPAN-B instruments detached from the spacecraft bus. In addition, the spacecraft is equipped with the FIELDS instrument suite (Bale et al. 2016), which includes four electric field antennas near the TPS, two magnetometers, and a search coil. FIELDS is used in conjunction with SPAN-E to determine the magnetic field vector and generate electron pitch angle distributions (ePADs).

In this section, we present a brief overview of the functionality of the SPAN-E instruments. For a comprehensive review of spherical ESAs, more information is provided in Carlson et al. (1982) and Carlson et al. (1998). The basic principle of an ESA is that incoming particles of a specific energy-charge ratio and path angle are guided through the instrument based on voltages of concentric toroidal hemispheres. Particles are then detected

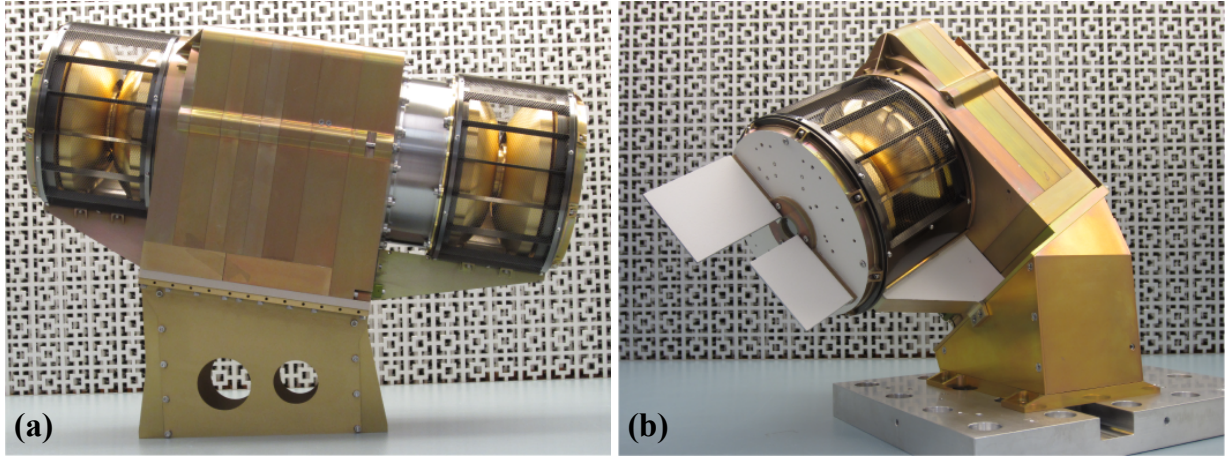


Figure 2.4: **(a)** SPAN-A composite instrument, with both SPAN-Ae (left) for electrons and SPAN-I (right) for ions. The Time-of-Flight (ToF) component is shown attached to SPAN-I within the silver-colored cylinder. **(b)** SPAN-B instrument with radiator fins. The defectors (gold) are shown in each analyzer. Adapted from Whittlesey et al. (2020).

as “counts” by multiple instrument anodes and other electronics to generate eVDFs. The eVDFs are measured as a function of the particle’s initial entry elevation angle θ , azimuth angle ϕ , and energy E . The elevation is measured from -90° to 90° and the azimuth is within the range 0° to 360° .

Figure 2.5 illustrates a cross section of a typical 3D toroidal top-hat analyzer, with an example electron trajectory shown in orange. Near the vertical aperture grids, the trajectory of a solar wind electron is pulled towards the center of the instrument based on the voltages applied to the upper and lower deflectors. The aperture grids outside of the instrument are grounded to isolate electric fields generated by the deflector voltages from the solar wind plasma. The deflector voltages (V_{D1} and V_{D2}) control the acceptable incident elevation angle θ of the electrons, which are then guided toward the opening at the top of the concentric hemispheres. In the case of SPAN-E, the possible elevation angles are within the range of $\pm 60^\circ$. The particle then reaches the hemispherical apex below the top cap, which allows particles to enter the hemispheres with varying azimuthal direction.

An electron will continue to travel between the two hemispheres without impacting the surfaces if its energy is within a given energy passband set by the instrument. This passband is determined by the voltages of the hemispheres and the size of the gap between the inner and outer hemispheres. The inner hemisphere is set to some analyzer bias voltage (V_A), while the outer hemisphere is typically held at ground. The lower section of the outer hemisphere can be set to a different voltage, known as the spoiler voltage V_S , to further reduce the width of the energy passband. The voltage difference between the hemispheres creates an electric field which only allows electrons with a specific energy to pass between the nested toroidal

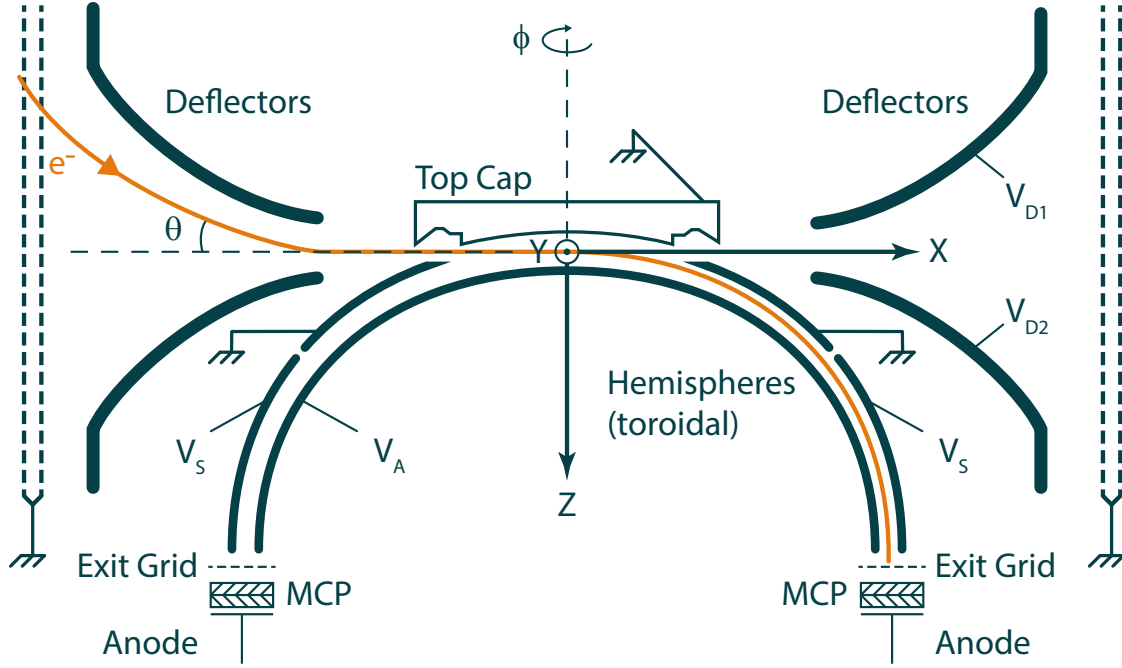


Figure 2.5: Cross section of a toroidal top-hat analyzer with a top cap, deflectors, hemispheres, and spoiler section. The orange curve portrays an electron travel path that reaches an instrument anode at angle ϕ based on various voltages. V_A is the analyzer bias (setting the electron energy), and V_{D1} and V_{D2} are the two deflector biases (setting the incoming θ angle). The spoiler voltage V_S can narrow the electron energy passband. The aperture grounded grids (vertical dashed lines) are fixed outside of the deflectors.

hemispheres. An ESA designed with a smaller gap distance leads to a smaller passband ($\Delta E/E$), resulting in higher energy resolution and a lower geometric factor.

If an electron reaches the bottom of the exit grid successfully, it will be recorded on a microchannel plate (MCP) above each instrument anode. The exit grid is kept at ground to contain electric fields from within the hemispheres. The MCP converts a single charged particle into a cascade of millions of electrons. The anode board then collects electron charge clouds from the MCP as a digital signal to eventually be recorded as one particle count. The anode board, shown in Figure 2.6, is composed of discrete metal pads with a certain angular width, each placed at a given center azimuth angle ϕ relative to the instrument center. SPAN is fabricated with 16 individual anodes, either with 6° or 24° resolution, for a total azimuthal range of 240° . The location of the narrow anodes were determined based on the expected direction of the strahl relative to each instrument. By varying the deflector and hemisphere voltages, a distribution of electrons can be measured and sorted by its velocity in terms of E , θ , and ϕ . A collective measurement can be performed in 0.218 seconds as SPAN-E sweeps through 32 energy and 8 deflector voltage steps (Whittlesey et al. 2020).

In the case of PSP, the spacecraft will traverse a broad range of heliocentric distances, and thus encounter a wide variation in particle fluxes. To account for this range, SPAN utilizes a mechanical attenuator (not shown in Figure 2.5), a thin piece of metal with small slits aligned with the axis of symmetry. By positioning the attenuator into the instrument aperture, the incoming particle flux is reduced by a factor of ten. This component can be actuated close to perihelion for each orbital encounter to avoid saturation, as particle flux is expected to be higher closer to the Sun. Another method of attenuation is shown in Figure 2.5, known as the spoiler section. The lower portion of the outer hemisphere can be set to ground or swept through voltages from 0 to 80 V separately from the rest of the outer hemisphere. When this spoiler section is set to a positive voltage, the energy passband is reduced, leading to a smaller $\Delta E/E$ and lower geometric factor. This form of attenuation reduces the particle flux by a factor proportional to the ratio of the spoiler voltage (V_S) to the inner hemisphere voltage (V_A).

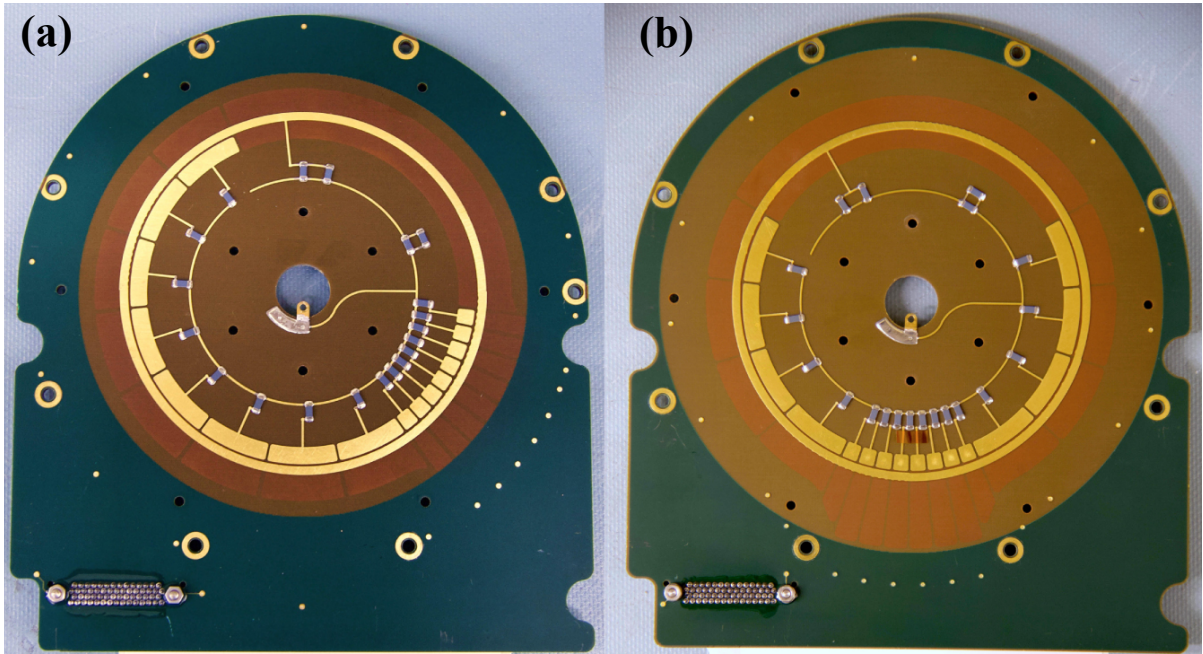


Figure 2.6: Anode boards for (a) SPAN-A and (b) SPAN-B, where the anodes are shown as curvilinear trapezoid cells within the outer gold ring spanning 240° in the azimuthal FOV. The 16 discrete anodes are shown for each instrument, either with 6° or 24° resolution. Adapted from Whittlesey et al. (2020).

SPAN-E Measurements

Level 1 files from SPAN-E measurements are compiled and sorted by particle counts as a function of energy, deflection, and anode bin steps over a given accumulation time. No instrument calibration is performed to generate these files. Instead, Level 2 files are created from Level 1 data by converting measurements from instrument units to physical units. We begin by converting deflection and anode bin steps into incident elevation θ and azimuth ϕ of the incoming particles. The number of counts measured by the instrument at a specific energy E_0 , elevation θ_0 , and azimuth ϕ_0 passband is expressed as $C(E_0, \theta_0, \phi_0)$. Based on the relation $E = m_e v^2/2$ for kinetic energy, where m_e is the mass of an electron, we can represent the measured counts as $C(v_0, \theta_0, \phi_0)$ in velocity space. Following Wüest et al. (2007) and Collinson et al. (2012), the number of observable counts can be related to the distribution function $f(\mathbf{v}, \mathbf{x}, t)$ of the ambient plasma particles and the particle flux entering the instrument aperture over time, defined as

$$C(v_0, \theta_0, \phi_0) = \int dt \int \int d^3\mathbf{v} (\hat{v} \cdot d\mathbf{A}) v R(\mathbf{v}, \mathbf{x}, t) f(\mathbf{v}, \mathbf{x}, t) \quad (2.3)$$

where the instrument response function $R(\mathbf{v}, \mathbf{x}, t)$ and the distribution function $f(\mathbf{v}, \mathbf{x}, t)$ are integrated over the accumulated measurement time, the surface area of the instrument aperture, and the incoming particle velocities. The term $(\hat{v} \cdot d\mathbf{A})v$ can be considered as the particle flux through the entrance aperture, where $d\mathbf{A}$ is the differential area element vector normal to the detector aperture. The detector response R describes the ability of the instrument to measure incident particles within a specific passband, and is necessary when converting data from counts to physical units. The response is a function of phase space and is a dimensionless quantity, where R should be maximized at the center values of the specified passband (E_0, θ_0, ϕ_0) .

To simplify equation 2.3, we assume the measurement integration time is small and approximately equal to Δt such that the incident particle phase space density does not vary significantly over the accumulation time, where $f(\mathbf{v}, \mathbf{x}, t) \rightarrow f(\mathbf{v}, \mathbf{x})$. In addition, we approximate the aperture area as A , which is small enough such that the phase space density is constant over the entrance aperture $f(\mathbf{v}, \mathbf{x}) \rightarrow f(\mathbf{v})$. We apply these assumptions to equation 2.3, and convert from velocity to energy space based on the expression $dE = mv dv$, to obtain

$$\begin{aligned} C(E_0, \theta_0, \phi_0) &= \Delta t \int \int \int v^2 dv \cos \theta d\theta d\phi (A \cos \theta \cos \phi) v \bar{R}(E, \theta, \phi) f(E, \theta, \phi) \\ &= A \Delta t \int \int \int \left(\frac{2E}{m_e} \right)^{3/2} \frac{dE}{\sqrt{2m_e E}} \cos^2 \theta d\theta \cos \phi d\phi \bar{R}(E, \theta, \phi) f(E, \theta, \phi) \quad (2.4) \\ &= \frac{2A \Delta t}{m_e^2} \int dE E \int d\theta \cos^2 \theta \int d\phi \cos \phi \bar{R}(E, \theta, \phi) f(E, \theta, \phi) \end{aligned}$$

where $d^3\mathbf{v} = v^2 dv \cos \theta d\theta d\phi$ and $(\hat{v} \cdot d\mathbf{A})v = vA \cos \theta \cos \phi$. We define $\bar{R}$ as the average of the instrument response function $R(E, \theta, \phi, \mathbf{x}, t)$ over the entrance aperture area and the

accumulation time, where $R(E, \theta, \phi, \mathbf{x}, t) \rightarrow \bar{R}(E, \theta, \phi)$. With an average response function, we define the geometric factor of the instrument as

$$G = \frac{A}{E_0^2} \int dE E \int d\theta \cos^2 \theta \int d\phi \cos \phi \bar{R}(E, \theta, \phi) \quad (2.5)$$

The geometric factor of an ESA defines the sensitivity of the instrument in detecting particles. In practice, the geometric factor can be estimated using particle ray-tracing simulations of the instrument. Further calibration can be performed prior to flight by firing a particle beam toward the ESA at various incident angles and energies to compare the response with the simulation. The geometric factor can be affected by manufacturing imperfections in the instrument and changes in detector efficiency over time. In-flight calibration, discussed later in this section, is used to improve the estimation of anode efficiencies for the geometric factor when converting between plasma units.

Finally, we assume that the incident phase space density does not change significantly with the response function, rewriting equation 2.4 as

$$C(E_0, \theta_0, \phi_0) = \frac{2A\Delta t}{m_e^2} f(E_0, \theta_0, \phi_0) \int dE E \int d\theta \cos^2 \theta \int d\phi \cos \phi \bar{R}(E, \theta, \phi) \quad (2.6)$$

We can then approximate the measured counts using the geometric factor from equation 2.5 to obtain a final expression relating particle counts to the distribution function, described as

$$\begin{aligned} C(E_0, \theta_0, \phi_0) &= \frac{2A\Delta t}{m_e^2} f(E_0, \theta_0, \phi_0) G \frac{E_0^2}{A} \\ &= \frac{2E_0^2}{m_e^2} f(E_0, \theta_0, \phi_0) G\Delta t \end{aligned} \quad (2.7)$$

However, some of our analysis in later sections is performed in terms of differential energy flux J_E , with units of eV / cm² / s / ster / eV, or particle number flux $J = J_E/E$. The differential energy flux J_E is related to the distribution function of incident plasma particles by

$$f(E_0, \theta_0, \phi_0) = \frac{m_e^2 J_E(E_0, \theta_0, \phi_0)}{2E_0^2} \quad (2.8)$$

Substituting the above expression into equation 2.7, we obtain

$$C(E_0, \theta_0, \phi_0) = J_E(E_0, \theta_0, \phi_0) G\Delta t \quad (2.9)$$

With an estimate of the geometric factor from pre-flight calibration, we can convert the particle counts into differential energy flux using equation 2.9. Figures 2.7a and 2.7b show the full FOV for SPAN-A and SPAN-B, respectively, in differential energy flux on 2021/08/10 at

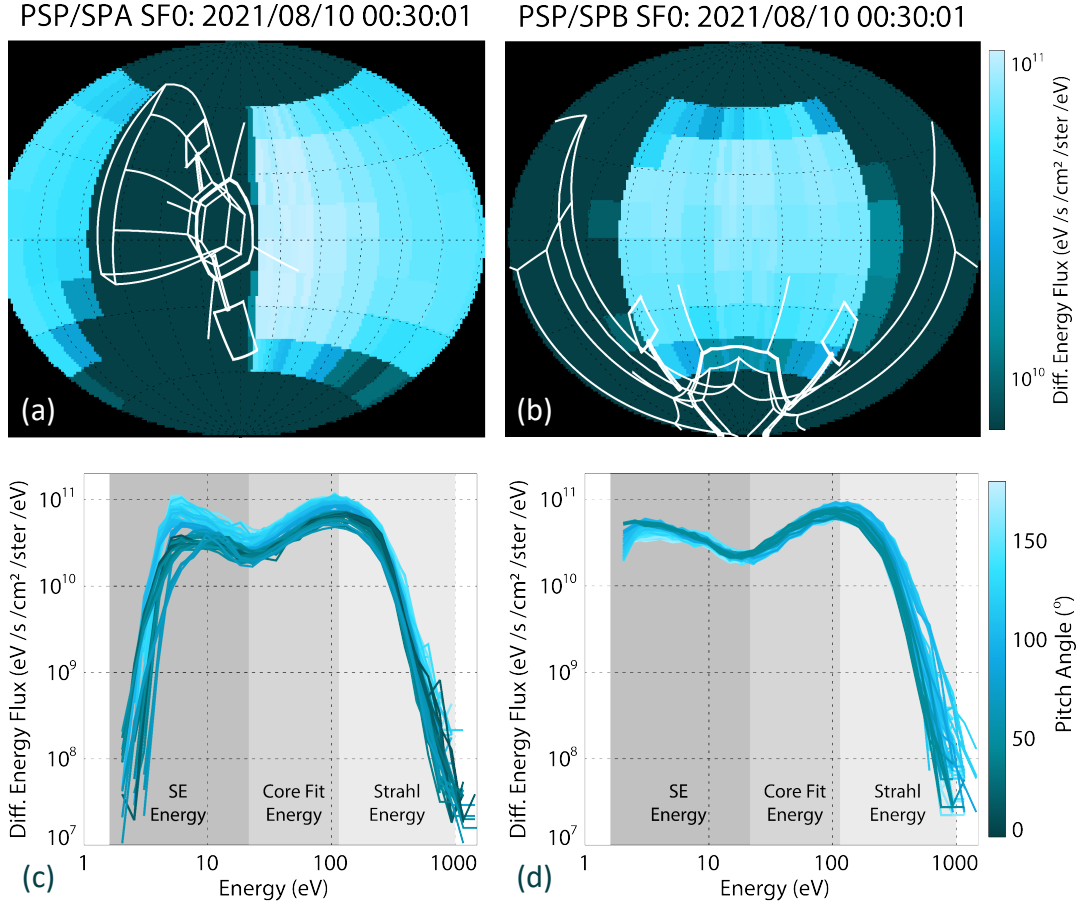


Figure 2.7: Differential energy flux for electron energies near 314 eV in the (a) SPAN-A and (b) SPAN-B FOVs, measured on 2021/08/10 at 00:30:01 UT. The spacecraft obstruction is projected (white) onto each FOV map, accounting for solar panel flap and feather angles. Differential energy flux as a function of energy for each pitch angle measured by (c) SPAN-A and (d) SPAN-B. The grey shading (dark to light) distinguishes the energy range of secondary electron (SE) contamination, fitted core electrons, and strahl electrons.

00:30 UT. The electron measurements are presented in each respective instrument coordinate frame, where the elevation and azimuth correspond to the latitude and longitude directions on the FOV map. A wireframe depiction of the spacecraft bus is projected onto each FOV to illustrate the amount of obstruction experienced by each instrument. It is important to note that the solar panels can rotate by more than 50° during a solar encounter, which can increase or decrease the level of obstruction to each instrument. This interference, along with other factors related to detector efficiency, results in decreased instrument sensitivity at the outer edge of the FOV for both ESAs, as shown in Figure 2.7. The calibration of the

outer anodes and deflection angles will be performed in future work.

Another method of displaying SPAN measurements is shown in Figures 2.7c and 2.7d for SPAN-A and SPAN-B, respectively. The differential energy flux is plotted as a function of energy, where each curve represents a different pitch angle direction. Three different electron populations measured by both instruments are depicted in ascending order of energy. Secondary electrons (SE), created from primary electrons (typically above 50 eV) impacting the spacecraft and instrument surfaces, contaminate the lower energy range below ~ 25 eV. The majority of the SE population is theorized to be generated within the instrument, known as internal secondary electrons. In the example shown, the distribution of core electrons is observed to peak around 105 eV. For each measurement, the peak of the core electrons is used to separate the fitting of the core and strahl populations, discussed later in this section. We limit our analysis to energies above the SE distribution (typically 25 eV) and below roughly 2 keV to characterize the electron core and strahl populations.

Electron Core Distribution Fitting

As described in Chapter 1, the electron core population is isotropic in the solar wind and can be modeled with a bi-Maxwellian distribution. We modify equation 1.6 by defining the velocity $\mathbf{V}$ in the plasma frame as $\mathbf{V} = \mathbf{v} - \mathbf{u}$, and replacing the parallel and perpendicular components with respect to pitch angle α , such that $V_{\parallel}^2 = V^2 \cos^2(\alpha)$ and $V_{\perp}^2 = V^2 \sin^2(\alpha)$. In addition, we compute a temperature ratio $R_T = T_{\parallel}/T_{\perp}$ between the parallel and perpendicular components. The temperature ratio R_T is expected to be close to one for the core distribution, and thus improves the least-squares fitting performance compared to fitting two separate temperatures that can typically range from 10 to 100 eV. After making these substitutions for the electron core distribution f_c , equation 1.6 becomes

$$\begin{aligned} f_c(V) &= \frac{n_c}{T_{\perp} \sqrt{T_{\parallel}}} \left(\frac{m_e}{2\pi k_B} \right)^{3/2} \exp \left(-\frac{m_e}{2k_B} \left[\frac{V^2 \cos^2(\alpha)}{T_{\parallel}} + \frac{V^2 \sin^2(\alpha)}{T_{\perp}} \right] \right) \\ &= n_c R_T \left(\frac{m_e}{2\pi k_B T_{\parallel}} \right)^{3/2} \exp \left(-\frac{m_e V^2}{2k_B T_{\parallel}} [R_T + \cos^2(\alpha)(1 - R_T)] \right) \end{aligned} \quad (2.10)$$

where our fitting parameters are the electron core density n_c , parallel temperature $T_{\parallel}$, temperature ratio R_T , and the electron core bulk velocity $\mathbf{v}_c = \mathbf{u} = \mathbf{v} - \mathbf{V}$. We also define the total electron core temperature as $T_c = \frac{1}{3}(T_{\parallel} + 2T_{\perp})$. The energy at which the core distribution peaks in differential energy flux can be estimated by $E_c = 2T_c$. Figure 2.8 shows electron measurements above 20 eV fitted to a core distribution, collected on 2021/08/08 at 00:00:05 UT by SPAN-A and SPAN-B. The strahl distribution can be observed at higher energies above ~ 200 eV for $\alpha > 135^\circ$.

The fitting of the core distribution is typically performed in units of counts instead of the distribution function ($\text{s}^3 / \text{km}^3 / \text{cm}^3$). Therefore, we convert the bi-Maxwellian distribution function in equation 2.10 to particle counts using equation 2.7. By converting

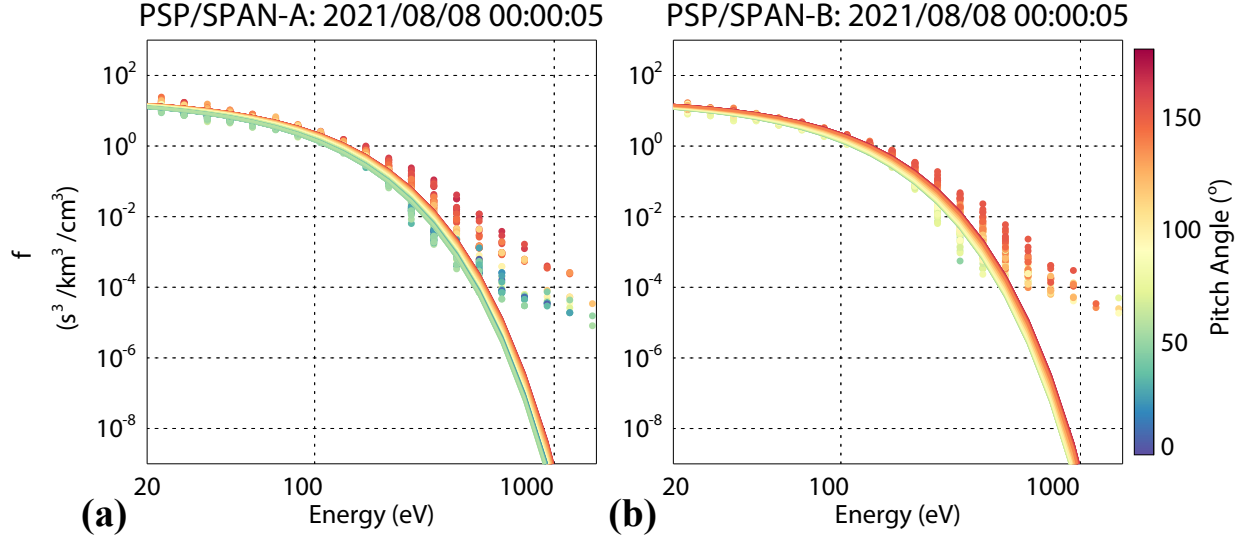


Figure 2.8: Electron distribution function for each measured energy observed by **(a)** SPAN-A and **(b)** SPAN-B on 2021/08/08 at 00:00:05 UT, where the color denotes the pitch angle direction. The corresponding curves are the fitted core distribution function from equation 2.10. In this example, the fitted values for n_c , T_c , and $|\mathbf{v}_c|$ are 588 cm^{-3} , 36.5 eV , and 604 km s^{-2} , respectively.

the function into counts, we can apply any necessary calibrations toward the instrumental energies, FOV, and geometric factor directly in the fitting procedure. It is important to note that the variables E_0 and m_e in equation 2.7 are typically calculated in units of eV and $\text{eV} / \text{km}^2 / \text{s}^2$, which introduces a multiplicative factor of 10^{-5} on the right-hand side of the equation due to the conversion from km to cm.

Additional precautions are needed when fitting the core distribution below 20 eV and above the peak of the core distribution E_c . At lower energies, the SE population contaminates the core component with a semi-Maxwellian distribution centered near 10 eV close to perihelion. To ensure the algorithm does not mistakenly fit the SE component as the core population, we remove all low energies near the SE energy peak E_{SE} . We find the cutoff energy between the SE and core distributions at which the slope of the differential energy flux reverses direction, referred to as E_{cut} . All energies below this boundary are removed from the core fit. At higher energies above E_c , the electron distribution function is dominated by the strahl population, especially near pitch angles close to 0° or 180° . For energies above E_c , the core distribution is fitted to counts measured near $\alpha = 90^\circ \pm 30^\circ$. However, for energies between E_{cut} and E_c , we fit the core distribution to all measured pitch angle directions.

Electron Moments and Spacecraft Potential

Another method of obtaining plasma parameters, such as electron density and temperature, is by computing moments of the distribution function. The general form for the n th moment M_n of $f(\mathbf{v})$ is given by

$$M_n = \int f(\mathbf{v}) \mathbf{v}^n d^3\mathbf{v} \quad (2.11)$$

Similar to the method of converting counts to physical units, we can rewrite equation 2.11 in spherical coordinates from velocity space (v, θ, ϕ) to energy space (E, θ, ϕ) as

$$\begin{aligned} M_n &= \int \int \int f(v, \theta, \phi) \mathbf{v}^n v^2 dv \cos \theta d\theta d\phi \\ &= \int \int \int \frac{m_e^2}{2E^2} J_E(E, \theta, \phi) \mathbf{v}^n \left[\frac{2E}{m_e} \frac{dE}{m_e v} \right] \cos \theta d\theta d\phi \\ &= \int \int \int J_E(E, \theta, \phi) \frac{\mathbf{v}^n}{v} \left[\frac{dE}{E} \right] \cos \theta d\theta d\phi \end{aligned} \quad (2.12)$$

where the angles range from $-90^\circ \leq \theta \leq 90^\circ$ and $0^\circ \leq \phi \leq 360^\circ$. For an ESA instrument, the energy, elevation, and azimuth for incident particles are measured at discrete values. The steps in energy, elevation, and azimuth can be represented by the indices i, j, k , where the integrand in equation 2.12 can be summed discretely over (E_i, θ_j, ϕ_k) , described as

$$M_n = \sum_i \sum_j \sum_k J_{E_{ijk}} \frac{\mathbf{v}_{ijk}^n}{v_i} \left[\frac{\Delta E_i}{E_i} \right] \cos \theta_j \Delta \theta_j \Delta \phi_k \quad (2.13)$$

where ΔE_i , $\Delta \theta_j$ and $\Delta \phi_k$ are the bin sizes for each energy, elevation, and azimuth step set by the instrument, and $J_{E_{ijk}} = J_E(E_i, \theta_j, \phi_k)$ is the indexed differential energy flux. For a top-hat ESA, $\Delta E/E$ should be roughly constant, as described in the previous section. For the zero-order moment, the electron density n is calculated using the equation

$$n = M_0 = \sum_i \sum_j \sum_k \frac{J_{E_{ijk}}}{v_i} \left[\frac{\Delta E_i}{E_i} \right] \cos \theta_j \Delta \theta_j \Delta \phi_k \quad (2.14)$$

To obtain the bulk velocity $\mathbf{u}$, the first-order moment is the number flux density vector, given by

$$\begin{aligned} nu_x &= M_{1x} = \sum_i \sum_j \sum_k J_{E_{ijk}} \cos \theta_j \cos \phi_k \left[\frac{\Delta E_i}{E_i} \right] \cos \theta_j \Delta \theta_j \Delta \phi_k \\ nu_y &= M_{1y} = \sum_i \sum_j \sum_k J_{E_{ijk}} \cos \theta_j \sin \phi_k \left[\frac{\Delta E_i}{E_i} \right] \cos \theta_j \Delta \theta_j \Delta \phi_k \\ nu_z &= M_{1z} = \sum_i \sum_j \sum_k J_{E_{ijk}} \sin \theta_j \left[\frac{\Delta E_i}{E_i} \right] \cos \theta_j \Delta \theta_j \Delta \phi_k \end{aligned} \quad (2.15)$$

To compute the pressure tensor P , the second-order moment is the momentum flux density tensor $\mathbf{\Pi}$, given by

$$\mathbf{\Pi} = M_2 = m_e \sum_i \sum_j \sum_k J_{E_{ijk}} \frac{\mathbf{v}_{ijk}^2}{v_i} \left[\frac{\Delta E_i}{E_i} \right] \cos \theta_j \Delta \theta_j \Delta \phi_k \quad (2.16)$$

To then obtain the pressure, we can transform equation 2.16 by the bulk velocity ($\mathbf{P} = \mathbf{\Pi} - mn\mathbf{u}\mathbf{u}$), given by

$$\begin{aligned} P_{xx} &= m_e \sum_i \sum_j \sum_k J_{E_{ijk}} v_i \cos^2 \theta_j \cos^2 \phi_k \left[\frac{\Delta E_i}{E_i} \right] \cos \theta_j \Delta \theta_j \Delta \phi_k - m_e n u_x^2 \\ P_{yy} &= m_e \sum_i \sum_j \sum_k J_{E_{ijk}} v_i \cos^2 \theta_j \sin^2 \phi_k \left[\frac{\Delta E_i}{E_i} \right] \cos \theta_j \Delta \theta_j \Delta \phi_k - m_e n u_y^2 \\ P_{zz} &= m_e \sum_i \sum_j \sum_k J_{E_{ijk}} v_i \sin^2 \theta_j \left[\frac{\Delta E_i}{E_i} \right] \cos \theta_j \Delta \theta_j \Delta \phi_k - m_e n u_z^2 \\ P_{xy} &= m_e \sum_i \sum_j \sum_k J_{E_{ijk}} v_i \cos^2 \theta_j \cos \phi_k \sin \phi_k \left[\frac{\Delta E_i}{E_i} \right] \cos \theta_j \Delta \theta_j \Delta \phi_k - m_e n u_x u_y \\ P_{yz} &= m_e \sum_i \sum_j \sum_k J_{E_{ijk}} v_i \cos \theta_j \sin \theta_j \sin \phi_k \left[\frac{\Delta E_i}{E_i} \right] \cos \theta_j \Delta \theta_j \Delta \phi_k - m_e n u_y u_z \\ P_{xz} &= m_e \sum_i \sum_j \sum_k J_{E_{ijk}} v_i \cos \theta_j \sin \theta_j \cos \phi_k \left[\frac{\Delta E_i}{E_i} \right] \cos \theta_j \Delta \theta_j \Delta \phi_k - m_e n u_x u_z \end{aligned} \quad (2.17)$$

where the temperature tensor T in units of eV can be calculated using the ideal gas law $\mathbf{P} = n\mathbf{T}$. Finally, we use the third-order moment to compute the energy flux density vector $\mathbf{Q} = \frac{m_e}{2} M_3$. We can relate this quantity to the heat flux vector $\mathbf{H}$, given by

$$\mathbf{H} = \mathbf{Q} - \mathbf{u} \cdot \mathbf{P} - \frac{1}{2} \mathbf{u} \text{Tr}(\mathbf{P}) \quad (2.18)$$

where $\text{Tr}(\mathbf{P})$ is the trace of the pressure tensor, with the scalar pressure $P = \text{Tr}(\mathbf{P})/3$. These moments can be generalized to ions to determine the total moment quantities, such as $n = n_e + n_i$ and $\mathbf{u} = (n_e m_e \mathbf{u}_e + n_i m_i \mathbf{u}_i) / (n_e m_e + n_i m_i)$.

Equations 2.13 through 2.17 are valid in the absence of external potentials. However, once a spacecraft is positioned in interplanetary space, it is continuously bombarded by solar radiation that generates photoelectrons and solar wind particles that create secondary electrons. The electrons emitted from these interactions with the spacecraft generate currents that lead to a final spacecraft potential (Mott-Smith et al. 1926). The net current is balanced (sums to zero) at equilibrium for a specific spacecraft electrostatic floating potential Φ relative to the ambient plasma potential, shown in Figure 2.9a. The incoming currents

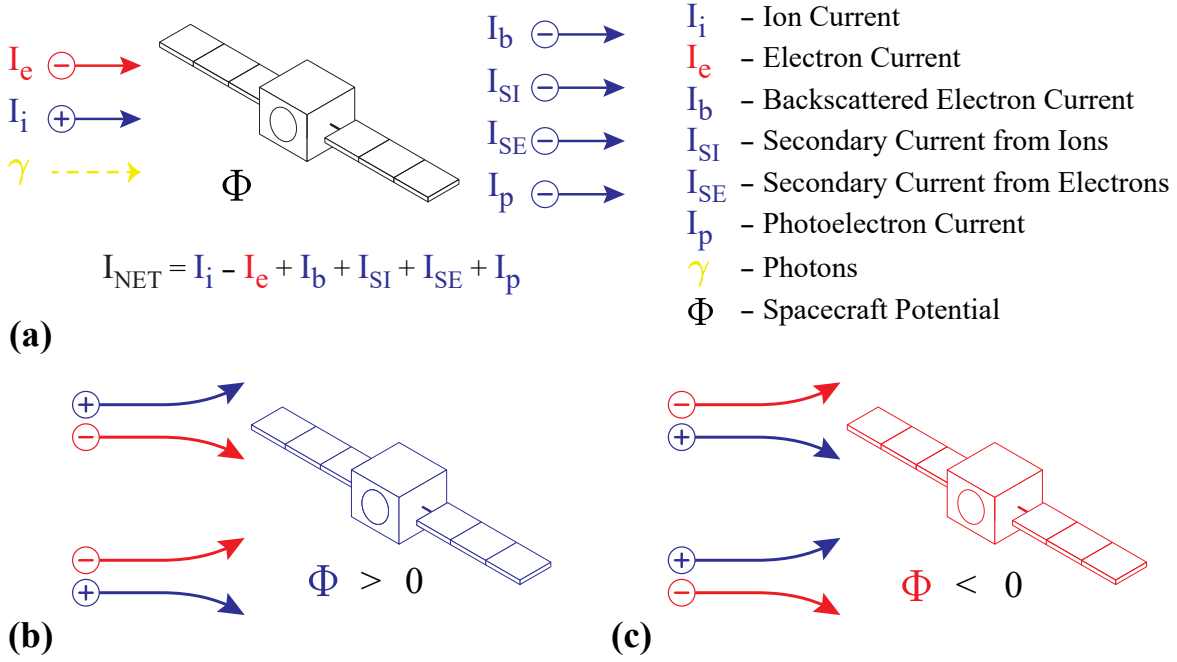


Figure 2.9: (a) An illustration of a spacecraft in interplanetary space experiencing various types of positive (blue) and negative (red) currents due to ions, electrons, and photons, where the net current will reach equilibrium at a given spacecraft potential Φ . (b) For a positive spacecraft potential (blue), electrons are attracted toward the spacecraft, while ions are repelled away, bending their corresponding trajectories. (c) For a negative spacecraft potential (red), ions are attracted toward the spacecraft, while electrons are repelled away.

are due to solar wind ions (I_i) and electrons (I_e), which create a positive (blue) and negative (red) current, respectively. The outgoing currents are due to backscattered electrons (I_b), secondary electrons from ions (I_{SI}), and secondary electrons from electrons (I_{SE}). In addition, the photoelectron current (I_p) is generated from the photoelectric effect by solar ultraviolet (UV) rays. As photons (γ) hit the spacecraft's surface, their energy is transferred to electrons in the material, causing them to overcome the surface binding energy, known as the work function, and be emitted into the surrounding space. If the spacecraft is not in direct sunlight, e.g., when PSP performs a Venus gravity assist maneuver behind the planet relative to the Sun, the photoelectron current is diminished, causing the spacecraft to charge more negatively. Otherwise, the spacecraft will charge positively, depending on the ambient plasma density and temperature. Other currents may exist that are not listed in Figure 2.9a, such as the thermionic electron current generated by electrons emitted from a heated spacecraft body. Except for the current I_e , all the listed currents positively charge the spacecraft relative to the solar wind plasma. A significant spacecraft potential can cause damage to onboard electronics and lead to inaccurate estimates of the plasma density.

Assuming the spacecraft is a sphere with an average radius r_{SC} of 1.0 m (the PSP spacecraft bus dimensions are 3.0 m $\times$ 1.0 m and the TPS diameter is 2.4 m), the surface capacitance C_{SC} of the spacecraft is given by

$$C_{SC} = 4\pi\epsilon_0 r_{SC} = 111.2 \text{ pF} \quad (2.19)$$

where ϵ_0 is the vacuum permittivity of free space (8.85 pF m⁻¹). The capacitance can be related to the spacecraft potential Φ based on charge Q on the spacecraft by $C_{SC}\Phi = Q$. By taking the time derivative of both sides of the expression, we can obtain a relation of the net current I_{net} to the change in the potential $\Delta\Phi$ over time, given by

$$\begin{aligned} C_{SC} \frac{d\Phi}{dt} &= \frac{dQ}{dt} \\ C_{SC} \frac{\Delta\Phi}{\Delta t} &= I_{net} \end{aligned} \quad (2.20)$$

where $I_{net} = 4\pi r_{SC}^2 j_{net}$, with j_{net} set as the net current density. For $\Delta\Phi = 100$ V and $I_{net} = 10 \mu\text{A}$, the time needed to reach charging equilibrium is about 1 millisecond. Spacecraft at geosynchronous altitudes tend to require a few milliseconds to reach current balance (Lai 2012). In general, the ambient electron flux is much higher than the ambient ion flux, due to the lower electron mass. The mass ratio between protons and electrons is $m_p/m_e \approx 1836$, which leads to a velocity ratio of $v_p/v_e \approx 43$ from equating the kinetic energy of each particle species ($\frac{1}{2}m_e v_e^2 = \frac{1}{2}m_p v_p^2$) for equal densities. Thus, if the electron flux is greater than the ion flux ($n_e v_e > n_i v_i$), the spacecraft potential is typically negative (Lai 2012). However, the photoelectron current can dominate over the electron current, causing the value of Φ to become positive.

For a positive spacecraft potential ($\Phi > 0$), shown in Figure 2.9b, electrons will be attracted toward the spacecraft, while ions will be repelled away. Essentially, electrons will gain energy equivalent to $q\Phi$ as they travel toward the spacecraft, enabling the analyzer to measure almost zero-energy electrons with energy $q\Phi$. Ions with energy less than $q\Phi$ cannot be detected at the spacecraft, and are typically modeled at the low-energy end of the spectrum with a theoretical distribution function. The opposite phenomenon occurs in Figure 2.9c for a negative spacecraft potential ($\Phi < 0$), with repelled electrons and attracted ions. Each scenario results in the formation of a plasma sheath around the spacecraft, where the density of the attractive particles is higher compared to the repelled species. The particle trajectories are bent so that the attractive particle species will focus toward the detector, measuring higher fluxes than the true flux in the ambient plasma. This effect tends to be more substantial for detector look directions along the surface of the spacecraft compared to directions perpendicular to the spacecraft bus.

For a positively charged spacecraft in sunlight, the photoelectron current can dominate if the plasma density is quite low. Once a photoelectron is generated, it can escape the spacecraft potential and contribute to the net outgoing current if it has energy above $q\Phi$. However, low-energy photoelectrons ($E < q\Phi$) cannot escape and instead return to the

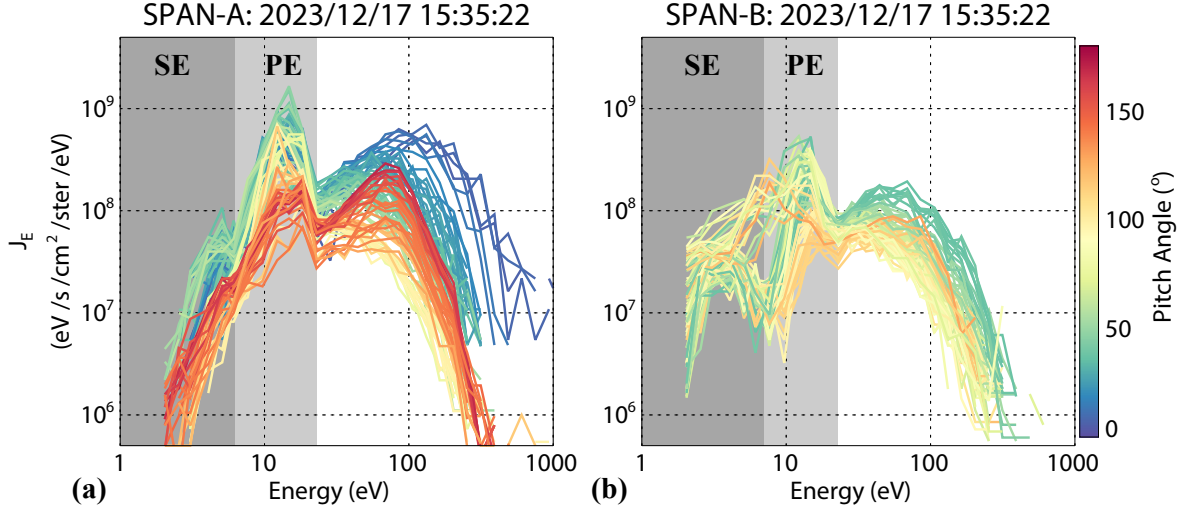


Figure 2.10: Differential energy flux (J_E) as a function of energy, sorted by pitch angle, measured by (a) SPAN-A and (b) SPAN-B on 2023/12/17 at 15:35:22 UTC. The dark grey region represents the majority of the SE population and the light grey region corresponds to the majority of the PE (photoelectron) population, with a break in the spectra near 20 eV.

spacecraft, where they can be measured by an ESA. The onboard detector can measure photoelectrons below $q\Phi$ and ambient electrons above $q\Phi$. Figure 2.10 shows an example of SPAN-A and SPAN-B measuring the solar wind electron distribution with photoelectrons (PE) roughly $91 R_S$ from the Sun on 2023/12/17 at 15:35:22 UTC. A distinct peak can be observed for the SE, PE, core, and suprathermal populations. The break in the energy spectrum between the PE and core populations occurs near 20 eV, signifying a spacecraft potential close to 20 V. In cases where the photoelectron distribution can be distinguished from the SE and core electrons, the SPAN-E instruments can be used to measure Φ . However, when the spacecraft potential is negative or if the photoelectrons cannot be effectively identified (either due to the detector energy range or overlapping electron populations in the energy spectrum with high density), other methods are required to determine Φ .

Another method to measure the spacecraft floating potential Φ for PSP is to utilize the four FIELDS whip antennas (V1 through V4) to measure the difference between Φ and the antenna floating potential Φ_A (Laakso et al. 1995). A negative bias current (I_{bias}) is applied to the antennas from the spacecraft in an effort to match the local plasma potential. The value of I_{bias} is determined for each orbital encounter during bias current sweeps to determine the current-voltage curve of the antenna sheath and minimize the voltage differential from the electric field booms. Liu et al. (2024, *in preparation*) is currently estimating these currents to determine both Φ_A and Φ for encounters 4 through 19. As estimates of the spacecraft potential have not been finalized, our study does not implement Φ in any calculations for later sections. However, we provide a detailed explanation below on implementing the spacecraft

potential to obtain accurate estimates of electron plasma parameters for any spacecraft in general. It is important to note that we only consider uniform surface charging for the spacecraft body, with other studies focusing on more complex models (Diaz-Aguado et al. 2021).

Consider two points in space: one far from the spacecraft in the ambient plasma and the other at the spacecraft near a particle detector that is affected by the spacecraft potential. The ambient energy E_A of a given particle far from the spacecraft is related to the energy E_M measured near the spacecraft based on the spacecraft potential, given by

$$E_A = E_M + q\Phi \quad (2.21)$$

As a simple check, equation 2.21 is valid for an electron with negative charge that is boosted ($E_M > E_A$) toward a positively charged spacecraft by an energy of $q\Phi$. In this scenario, the measured electron flux J_M will be higher than the ambient electron flux J_A , as previously discussed. However, by Liouville's theorem, the phase space density f is conserved along the trajectory of the particle toward the detector on the spacecraft. This condition is true under the assumption that there is no significant diffusion, collisions, or scattering occurring within phase space. Therefore, the distribution function f_M measured by the detector is equal to the ambient plasma distribution f_A far from the spacecraft. In addition, we assume the distribution function is measured by a "point" detector and that Φ is spherically symmetric about the spacecraft body.

Utilizing Liouville's theorem gives us $f_A = f_M$, where we can substitute E_A as a dependent variable in f_A with $E_M + q\Phi$ to calculate the ambient distribution function from the measured particles. This implies that a Maxwellian function, representing the core electrons from equation 2.10, is scaled by a constant factor of $\exp(q\Phi/(k_B T))$. Figure 2.11 shows the effect of the spacecraft potential on a Maxwellian distribution function. A normalized Maxwellian distribution with $\Phi = 0$ V will appear as a linear function in log space. A value of Φ greater than zero volts will only lead to a shift to higher energies with no effect on the slope (Lai et al. 2020). We can convert the normalized distribution function to a quantity proportional to the differential energy flux J_E by multiplying f by E^2 . The difference between the J_E curves with and without a positive spacecraft potential is the height. The peak of the curve remains at the same energy under different values of Φ , but the peak of the curve with non-zero potential is 1.65 times larger than the curve with $\Phi = 0$ V. In addition, ambient electrons are not observed for energies below $q\Phi$, and instead we measure secondaries and photoelectrons for $E < q\Phi$. Therefore, replacing the energy E_A in the core electron distribution function with $E_M + q\Phi$ can increase or decrease the density estimate.

Substituting the measured electron energy with the ambient energy into the core distribution function for applications of least-squares fitting is a simple and straightforward process. Unfortunately, the process is a bit more involved when implementing the spacecraft potential in the plasma moment calculations described by equation 2.11. To obtain the n th moment M_n of the ambient plasma, we start by integrating $f_A(v_A, \theta, \phi) \mathbf{v}_A^n$ over the velocity space element $d^3 \mathbf{v}_A$, described as

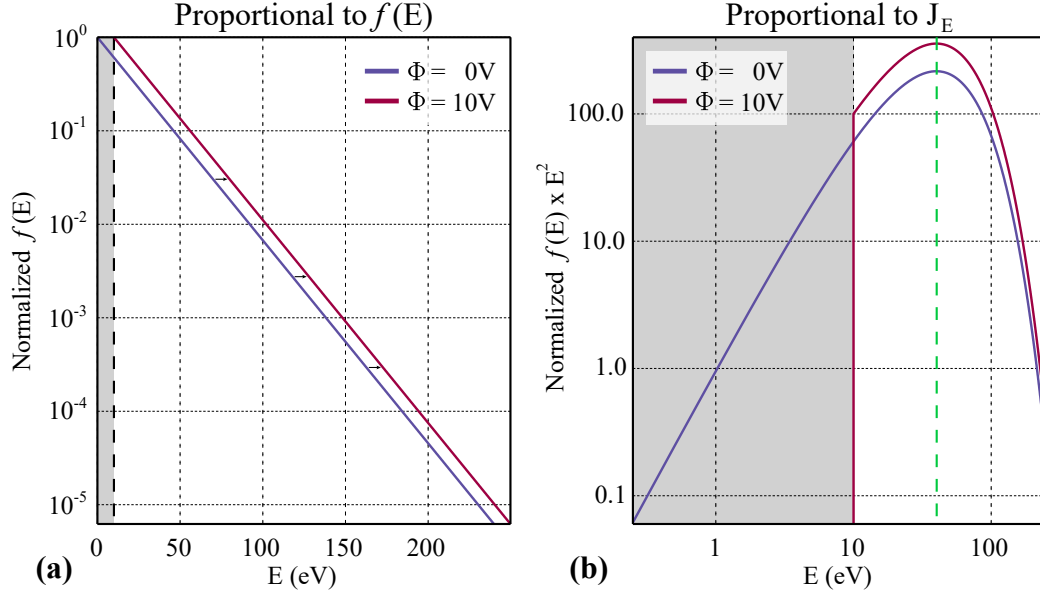


Figure 2.11: **(a)** Normalized Maxwellian distribution function $f(E)$ in log space of ambient electrons measured with a spacecraft potential Φ set to 0 V (purple) and 10 V (red). The vertical black line indicates the shift in energy by $q\Phi = 10$ eV. **(b)** Normalized Maxwellian distribution function $f(E)$ multiplied by E^2 (which is proportional to the differential energy flux J_E) in log space of ambient electrons measured with a spacecraft potential Φ set to 0 V (purple) and 10 V (red). The vertical green line indicates the peak of the curve is the same in energy for both values of Φ . The grey region represents the shift by $q\Phi$, where we expect to observe secondaries and photoelectrons for $E < q\Phi$.

$$\begin{aligned}
 M_n &= \int \int \int f_A(v_A, \theta, \phi) \mathbf{v}_A^n d^3\mathbf{v}_A \\
 &= \int \int \int f_A(v_A, \theta, \phi) \mathbf{v}_A^n v_A^2 dv_A d\Omega \\
 &= \int \int \int f_M(v_M, \theta, \phi) \mathbf{v}_A^n v_A^3 \frac{dv_A}{v_A} d\Omega \\
 &= \int \int \int f_M(E_M, \theta, \phi) \mathbf{v}_A^n \left(\frac{2E_A}{m_e} \right)^{3/2} \left[\frac{dE_A}{2E_A} \right] d\Omega
 \end{aligned} \tag{2.22}$$

where we have used Liouville's theorem as $f_A(v_A, \theta, \phi) = f_M(v_M, \theta, \phi)$, the kinetic energy equation $E_A = \frac{1}{2}mv_A^2$, and the expression $dE_A/2E_A = dv_A/v_A$. Using equation 2.21 and the relation $dE_A = dE_M$, the E_A terms can be replaced with the measured energy. In addition, we can convert the phase space density f_M to differential energy flux J_E in equation 2.22, resulting in

$$\begin{aligned}
 M_n &= \int \int \int \frac{m_e^2}{2E_M^2} J_{EM}(E_M, \theta, \phi) \mathbf{v}_A^n \left(\frac{2}{m_e} \right)^{3/2} \frac{(E_M + q\Phi)^{3/2}}{E_A} \left[\frac{dE_M}{2} \right] d\Omega \\
 &= \int \int \int \frac{\sqrt{2m_e}}{E_M} J_{EM}(E_M, \theta, \phi) \mathbf{v}_A^n (E_M + q\Phi)^{1/2} \left[\frac{dE_M}{2E_M} \right] d\Omega \\
 &= \int \int \int \frac{\sqrt{2m_e}}{E_M} J_{EM}(E_M, \theta, \phi) v_A^n [\hat{\sigma}]^n (E_M + q\Phi)^{1/2} \left[\frac{dE_M}{2E_M} \right] d\Omega
 \end{aligned} \tag{2.23}$$

where we have combined like terms for the energy, and defined the ambient electron velocity as $\mathbf{v}_A^n = v_A^n [\hat{\sigma}]^n$. The scalar v_A^n provides the magnitude of the velocity, while the tensor $[\hat{\sigma}]^n$ provides the directionality of the velocity. For example, the second rank tensor $[\hat{\sigma}]^2$ can be written as

$$[\hat{\sigma}]^2 = \begin{bmatrix} \hat{v}_x \hat{v}_x & \hat{v}_y \hat{v}_x & \hat{v}_z \hat{v}_x \\ \hat{v}_x \hat{v}_y & \hat{v}_y \hat{v}_y & \hat{v}_z \hat{v}_y \\ \hat{v}_x \hat{v}_z & \hat{v}_y \hat{v}_z & \hat{v}_z \hat{v}_z \end{bmatrix} = \begin{bmatrix} \cos^2(\theta) \cos^2(\phi) & \cdots & \cdots \\ \cos^2(\theta) \cos(\phi) \sin(\phi) & \cos^2(\theta) \sin^2(\phi) & \cdots \\ \cos(\theta) \cos(\phi) \sin(\theta) & \cos(\theta) \sin(\phi) \sin(\theta) & \sin^2(\theta) \end{bmatrix} \tag{2.24}$$

where we have converted to spherical coordinates, and the matrix entries with a $\cdots$ represent the symmetric off-diagonal elements of the second rank tensor. By introducing $[\hat{\sigma}]^n$, we can replace v_A^n with the energy measured by the detector to combine the energy terms, given by

$$\begin{aligned}
 M_n &= \int \int \int \frac{\sqrt{2m_e}}{E_M} J_{EM}(E_M, \theta, \phi) \left(\frac{2E_M}{m_e} + \frac{2q\Phi}{m_e} \right)^n [\hat{\sigma}]^n (E_M + q\Phi)^{1/2} \left[\frac{dE_M}{2E_M} \right] d\Omega \\
 &= \int \int \int \frac{m_e}{2E_M} J_{EM}(E_M, \theta, \phi) \left(\frac{2E_M}{m_e} + \frac{2q\Phi}{m_e} \right)^{\frac{n+1}{2}} [\hat{\sigma}]^n \left[\frac{dE_M}{E_M} \right] d\Omega \\
 &= \int \int \int J_{EM}(E_M, \theta, \phi) v_M^{-2} \left(v_M^2 + \frac{2q\Phi}{m_e} \right)^{\frac{n+1}{2}} [\hat{\sigma}]^n \left[\frac{dE_M}{E_M} \right] \cos \theta \, d\theta \, d\phi
 \end{aligned} \tag{2.25}$$

where we have converted the final integral back to the particle velocities measured by the instrument, while continuing to integrate all terms in energy space. To ensure that equation 2.25 is derived correctly, we can take the limit that Φ approaches zero volts. With no spacecraft potential, the measured velocity term becomes v_M^{n-1} , which agrees with equation 2.12 when accounting for $[\hat{\sigma}]^n$. Similar to equation 2.13 with no spacecraft floating potential, we can rewrite equation 2.25 as a discrete sum over the indices i, j, k for the energy (or velocity), elevation, and azimuth steps, respectively, for the measured particles, defined as

$$M_n = \sum_i \sum_j \sum_k J_{E_{ijk}} v_{M_i}^{-2} \left(v_{M_i}^2 + \frac{2q\Phi}{m_e} \right)^{\frac{n+1}{2}} [\hat{\sigma}_{jk}]^n \left[\frac{\Delta E_{M_i}}{E_{M_i}} \right] \cos \theta_j \, \Delta \theta_j \, \Delta \phi_k \tag{2.26}$$

where equation 2.26 reduces to equation 2.13 when $\Phi = 0$ V.

A final remark is that the integration bounds do change when implementing a spacecraft potential in the moments equation. For measuring ambient electrons in a positive spacecraft potential, we must remove all energy steps below $q\Phi$ as these energies are contaminated by photoelectrons and secondary electrons. For a negative potential, the portion of the distribution that is not measured can be modeled using a Maxwellian distribution function in the case of core electrons. Once an accurate estimate of the spacecraft floating potential is determined from the FIELDS whip antennas, or from the photoelectron discontinuity in the energy spectra in the electron distributions, we can implement equation 2.26 to obtain more accurate estimates of the plasma moments for future work.

Before computing plasma moments, we must combine the data from the two separate electron instruments on board PSP. SPAN-A and SPAN-B each cover different regions of the entire 4π steradian sky, with a small amount of overlap. Due to the limited FOV coverage, only partial moments can be calculated when using data from one instrument. To obtain a more complete calculation, we must select a method of combining the information from the two ESAs to determine the total electron moments. One approach is to compute the plasma moments separately for each instrument and then recombine them in a unified coordinate system (typically the spacecraft frame). However, this does not account for both the overlap and missing FOV regions between the two instruments, leading to inaccurate moment estimates. Another method is to first rotate the SPAN-A and SPAN-B measurements into the spacecraft frame, and interpolate the data onto a new FOV map. By resampling on a complete regular grid in elevation and azimuth, this technique addresses the non-uniform FOV coverage regardless of measurement gaps. Before re-gridding our data spatially, we must account for a possible energy shift between the two instruments (discussed in a later section). If SPAN-A is shifted in energy by a small amount compared to SPAN-B (e.g., due to a small voltage shift in an operational amplifier output), we can interpolate the distribution function in energy to ensure the same energy steps are sampled for both instruments.

Once the energy shift correction is applied, we can compare different interpolation methods to obtain a complete FOV of electron measurements. It is important to note that we have preemptively derived equations 2.13 through 2.17 in terms of units of differential energy flux instead of counts. Spatially interpolating the data in units of counts would require a new geometric factor for each new bin in the spherical grid. Instead, resampling the data in terms of differential energy flux intrinsically accounts for the original geometric factor for each instrument during the unit conversion. Figures 2.12a and 2.12b show the SPAN-A and SPAN-B measurements for 314 eV electrons individually rotated to the spacecraft frame. Removing the outer look direction bins for both instruments due to obstruction effects, the overlap between SPAN-A and SPAN-B is reduced. The more critical issue is the missing FOV regions near $\phi \sim \pm 120^\circ$ and the polar regions (e.g., the TPS obstructs the south pole of the FOV). We compare four different interpolation methods to combine the FOV of both instruments on a regular grid with 5° resolution, presented in Figures 2.12c through 2.12f for an electron measurement during orbital encounter 9. The Kriging method (Chilès et al. 2018) performs a weighted average of data points surrounding the query point based on an

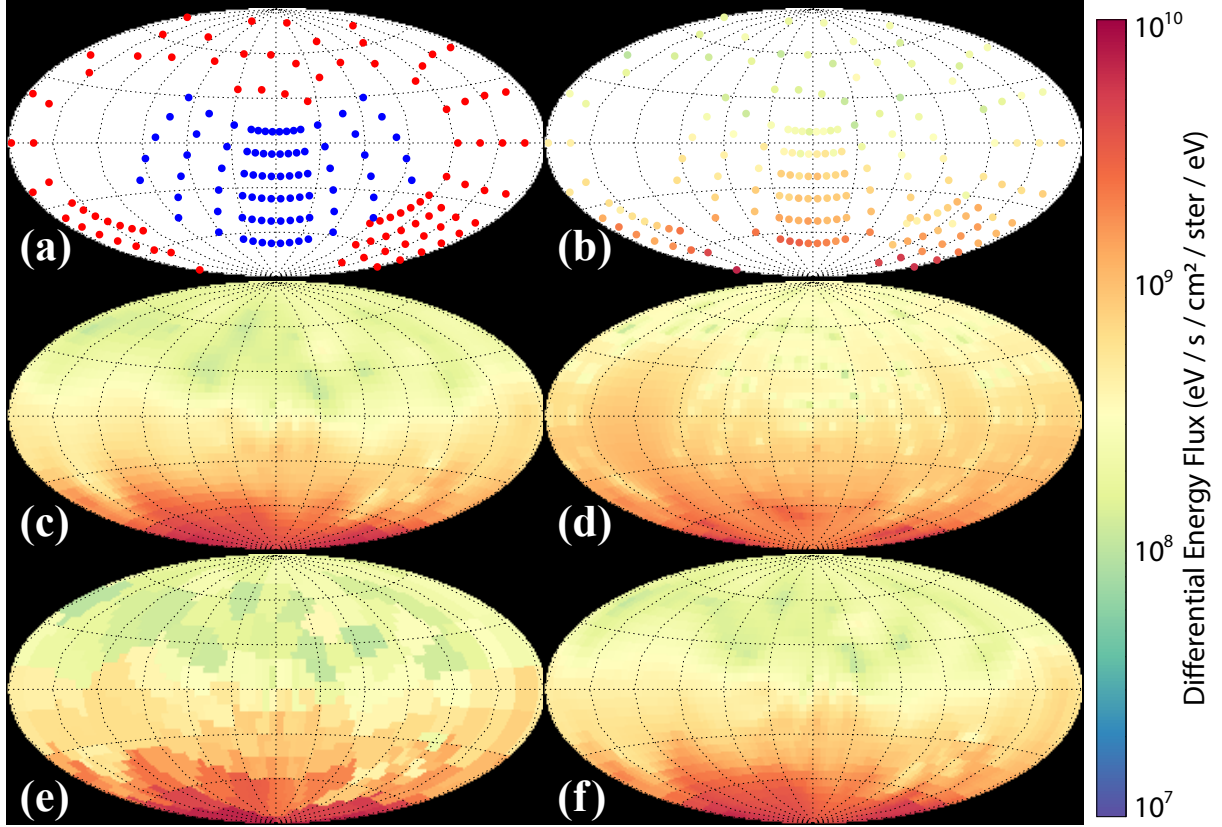


Figure 2.12: FOV maps of SPAN-E in the spacecraft frame for 314 eV electrons on 2021/08/08 at 00:00:05 UT. (a) Center of the look direction bins for SPAN-A (red) and SPAN-B (blue) projected from each instrument frame to the spacecraft frame. (b) SPAN-A and SPAN-B measurements in differential energy flux located at the center of each azimuth and elevation bin. The SPAN-E measurements are interpolated on a sphere using the (c) Kriging, (d) inverse distance, (e) nearest neighbor, and (f) natural neighbor methods.

exponential variogram model that reflects the spatial structure and variability of the data. The inverse distance approach calculates weights for each point that are inversely proportional to the square of its distance to the query point. The nearest neighbor method assigns the value of the query point to the closest data point. Finally, the natural neighbor technique resamples the data based on a linear combination of the three vertices of its enclosing Delaunay triangle and their adjacent vertices. The Kriging and natural neighbor methods provide the smoothest combined FOV map compared to the other two approaches. However, the Kriging method tends to introduce negative values of differential energy flux when zero counts are measured in some look directions. Therefore, we utilize the natural neighbor method to resample electron measurements to estimate the complete plasma moments.

eVDF Hybrid Method

Given the complexity of electron populations observed by SPAN-E, distribution fitting and moment calculations alone are insufficient when applied separately to characterize the entire eVDF. Thus, to obtain accurate electron parameters, we implement a hybrid technique utilizing both methods. We begin by applying the method outlined previously to determine the optimal bi-Maxwellian function that models the core distribution across an energy range that excludes the majority of SE contamination and strahl electrons. The fitted core distribution is then subtracted from the SPAN-E data to isolate the SE and suprathermal distributions in the eVDF. From the residual distribution at energies above E_c , we can calculate the suprathermal moments without requiring a functional form for the strahl distribution. Finally, we can combine the calculated densities of the core and suprathermal components to obtain a total electron density n_e . This approach will be essential in subsequent sections for calibrating both instruments, accounting for factors such as spacecraft potential, anode efficiency, and dead time corrections.

A similar approach can be applied to compute moments for the entire eVDF. We can either restrict the lower energy range to exclude SE contamination or subtract a fitted SE distribution and calculate moments on the residual distribution. While implementing a fitted SE distribution can provide a more complete estimate of the plasma moments, more work is needed to better characterize the low energy population. The SE distribution is observed to resemble a Maxwellian shape in some time intervals, but is dependent on the directionality of higher energy strahl electrons. Future research could model the response of SEs using particle ray-tracing simulations based on electron impacts that occur within the instrument. For the remaining sections, we will calculate combined moments for the entire eVDF only for energies above E_{cut} .

DC Voltage Offset and SE Distributions

The next calibration procedure addresses a slight shift in the measured energy for the two SPAN instruments. The measured “nominal” energy may be shifted by a constant amount ϵ_0 due to a DC voltage offset in the high voltage power supply (HVPS) or the digital to analog converter (DAC) that controls the HVPS. A shift of a few millivolts in the DAC output could lead to ϵ_0 values on the order of a few electron-volts for a large HVPS gain. This deviation would shift the curves in Figures 2.7c and 2.7d either left or right along the x-axis. Moreover, the offset may vary with azimuth angle, as the nested hemispheres of the detectors may not be perfectly symmetric. This asymmetry could result in different gap sizes between hemispheres for each discrete anode, leading to variations in the energy passband. Accurately characterizing this energy offset is essential for refining the resolution of fits to the SE distribution, core electron distribution, and spacecraft potential.

As a preliminary approach to obtain an approximate estimate of the energy offset, we extrapolate the observed particle counts for each instrument toward lower energy. It is expected that the number of counts should drop to zero toward zero electron energy. We

first sort the particle counts as a function of energy for each deflection and anode bin and only consider the 10 lowest energy steps. The particle counts are averaged over all eight deflection bins. Different extrapolation methods are used to determine the energy at which the particle counts drop to zero for each anode separately. These techniques include linear extrapolation, quadratic extrapolation with a three-point neighborhood, and a least squares quadratic fit with a four-point neighborhood. A more accurate estimate of the energy offset is generally obtained for count drop-offs that occur within the energy range of the instrument. At these low energies, most of the counts are theorized to originate from secondary electrons generated inside the instrument, which are not affected by spacecraft floating potential effects.

Before orbital encounter 10, the lowest energy bin recorded was 2 eV. The lowest energy step for PSP encounters 10 through 19 was then set around 10 eV, making it difficult to resolve the energy offset using this method. Thus, we focus our analysis of energy offsets

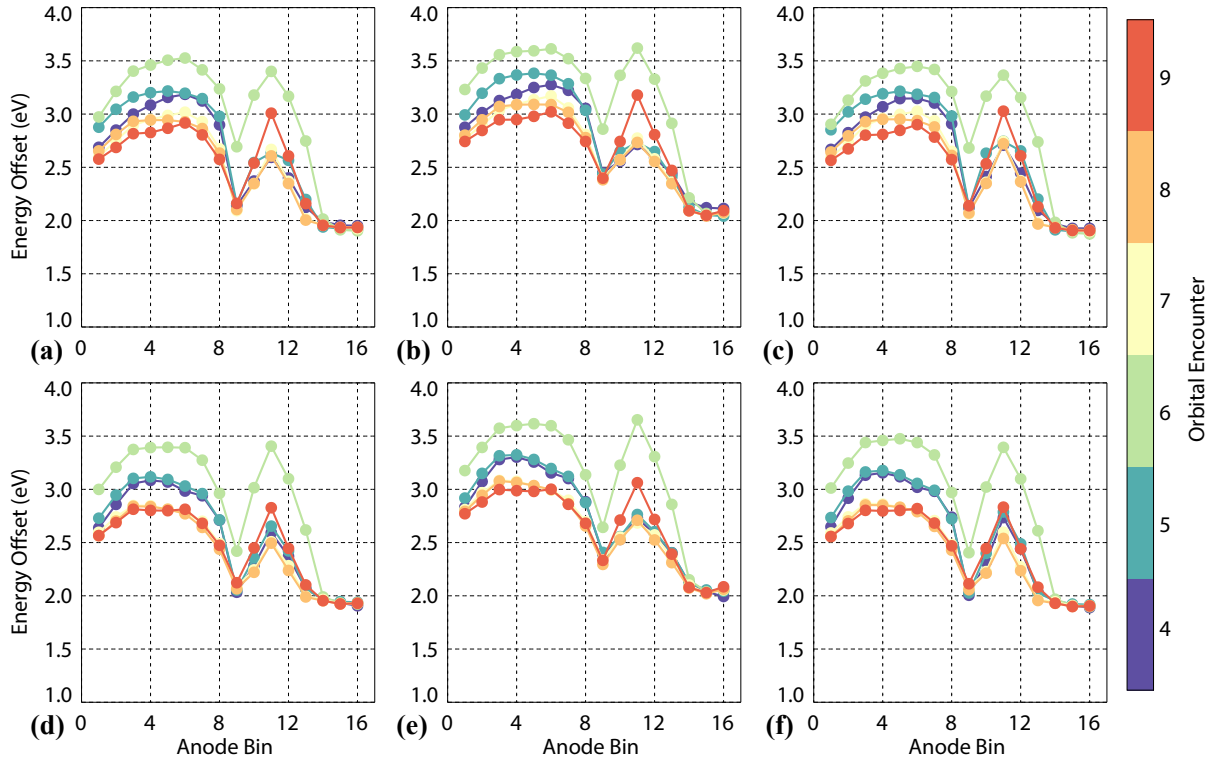


Figure 2.13: SPAN-A instrumental energy offsets for each anode determined by (a) linear extrapolation, (b) least squares quadratic fit, and (c) quadratic extrapolation for measurements with the mechanical attenuator disengaged. For measurements with the attenuator engaged, the energy offsets are also estimated from (d) linear extrapolation, (e) least squares quadratic fit, and (f) quadratic extrapolation. Each curve represents a different PSP orbital encounter.

for orbital encounters 4 through 9. A distribution of calculated energy offsets is computed for each anode, encounter, attenuator state, and extrapolation method. The mean energy offsets of these distributions are shown in Figures 2.13 and 2.14 for SPAN-A and SPAN-B, respectively. The energy offset is plotted as a function of anode bin number for all 16 anodes. The top panels are derived from measurements when the attenuator is disengaged, and the bottom panels are computed for cases with the attenuator engaged. The mean energy offsets determined by linear extrapolation, least squares quadratic fitting, and quadratic extrapolation are shown in the left, middle, and right panels, respectively. The energy offsets do not vary monotonically with time for both instruments.

The different extrapolation techniques in Figure 2.13 result in similar energy offsets for SPAN-A. In contrast, the energy offsets for SPAN-B in Figure 2.14 are highly dependent on the extrapolation method for the outer anodes. For SPAN-A, most of the counts reduce to

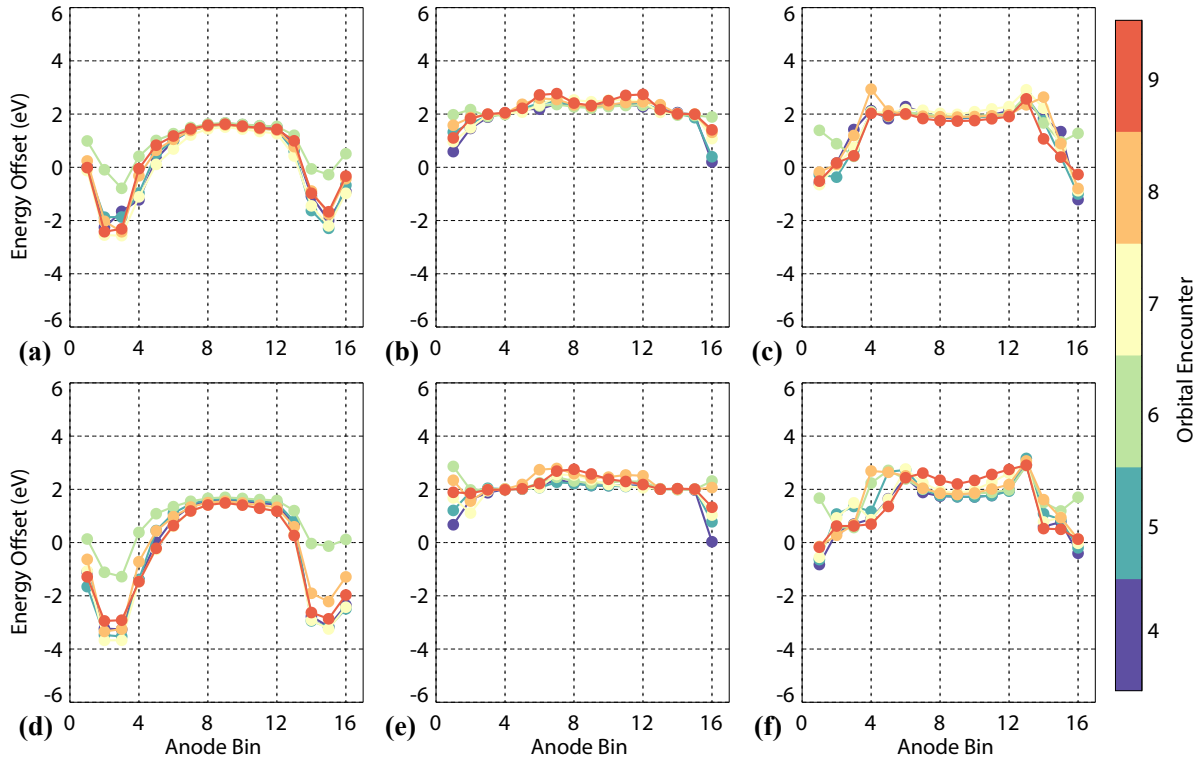


Figure 2.14: SPAN-B instrumental energy offsets for each anode determined by (a) linear extrapolation, (b) least squares quadratic fit, and (c) quadratic extrapolation for measurements with the mechanical attenuator disengaged. For measurements with the attenuator engaged, the energy offsets are also estimated from (d) linear extrapolation, (e) least squares quadratic fit, and (f) quadratic extrapolation. Each curve represents a different PSP orbital encounter.

zero above 2 eV, which allows for simple identification of the offset energy as it lies within the measured energy range. For SPAN-B, the peak of the SE distribution is near the lower energy limit of the instrument, making it difficult to determine the offset energy by extrapolation. For the inner 8 anodes of SPAN-B, all the extrapolation techniques estimate an offset near 2 eV. However, the energy offsets of the outer anodes range from -3 to 3 eV. The linear extrapolation would overestimate the slope of the SE distribution since the derivative of the counts decreases from 0 eV to E_{SE} . Thus, this type of extrapolation provides a lower bound on the energy offset. The quadratic methods instead show mostly positive energy offsets. We use the quadratic extrapolation to determine the energy offsets, as fitting to the SE distribution requires a more complex function.

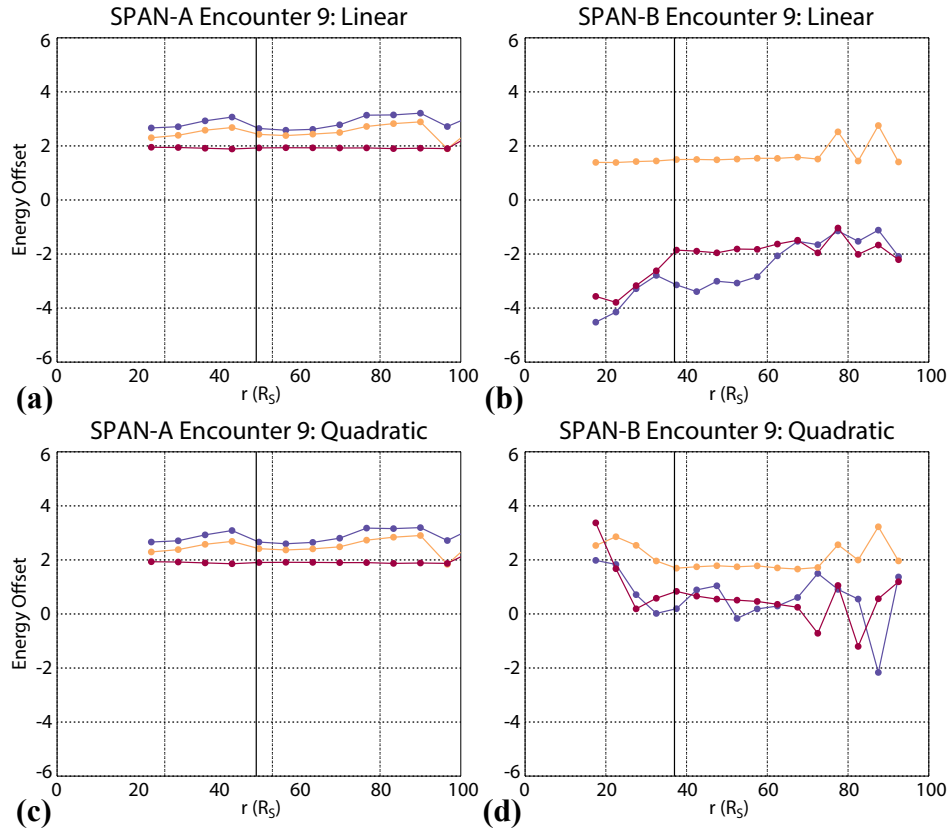


Figure 2.15: Energy offsets as a function of heliocentric distance r for orbital encounter 9 determined by linear extrapolation for (a) SPAN-A and (b) SPAN-B. Energy offsets are also shown based on quadratic extrapolation for (c) SPAN-A and (d) SPAN-B. The different curves represent the average offsets for anode bins 3 (purple), 10 (orange), and 15 (red), where the black vertical line denotes at which heliocentric distance the attenuator is engaged as PSP travels closer to the Sun.

We can compare the energy offset for the linear and quadratic extrapolations to provide insight as to how the SE distribution changes under different conditions. The offset can be analyzed as a function of heliocentric distance r since the particle flux will highly depend on proximity to the Sun. Figure 2.15 portrays the energy offset as a function of heliocentric distance for SPAN-A and SPAN-B during orbital encounter 9 for anode bins 3, 10, and 15 using either the linear or quadratic extrapolation. The black vertical line denotes when the attenuator is engaged for smaller distances. The energy offsets for SPAN-A are roughly constant for all extrapolation methods during the entire encounter. This is also mainly true for SPAN-B when the attenuator is disengaged and PSP is far from the Sun. However, as the spacecraft approaches the Sun (especially when the attenuator is engaged), SPAN-B linear offsets decrease while SPAN-B quadratic offsets increase. This may correspond to changes in the SE distribution for each solar encounter rather than instrumental offsets. Another method of determining the energy offset ϵ_0 is to fit the SE population to an empirical distribution function. A bi-Maxwellian model does not consistently provide an accurate fit to the SE component over an entire encounter. A slightly more complicated function is the Probability Density Function (PDF) of the Weibull distribution (equation 8.1 in Hallinan (1993)), defined for positive energies E with an energy shift ϵ_0 as

$$f(E) = \frac{b}{a} \left(\frac{E + \epsilon_0}{a} \right)^{b-1} \exp \left[- \left(\frac{E + \epsilon_0}{a} \right)^b \right] \quad (2.27)$$

where a and b are referred to as the scale and shape parameters, respectively. This function can take the form of an exponential, right-skewed, or Maxwellian distribution based on the parameter values. The SE distribution tends to be right-skewed with energy, but varies with particle flux of high energy electrons. Figure 2.16 shows SPAN-A and SPAN-B electron count distributions measured far and near the Sun with superimposed Weibull distributions. The solid and dashed black curves represent the Weibull distributions corresponding to the larger anodes and smaller anodes, respectively. Figure 2.16a does not include a solid black curve as the small anodes do not observe a significant SE population. The distributions exhibit a more Maxwellian profile at greater heliocentric distances when the attenuator is disengaged and shift toward a right-skewed distribution when the attenuator is engaged. It is important to note that a Weibull distribution does not perfectly model the data measured close to the Sun for energies above 15 eV. Future work could develop an empirical model to fit the SE population to better constrain the energy offset for each instrument separately.

Using the extrapolation methods as our current best estimate for the energy offsets, the SPAN-A and SPAN-B offsets are averaged over all encounters with the quadratic extrapolation method, as shown in Figure 2.17. The averaged energy offsets range from roughly 1.9 to 3.1 eV for SPAN-A and from -0.1 to 3 eV for SPAN-B. For SPAN-A, values of ϵ_0 are slightly smaller for measurements with the attenuator engaged. Separating the energy offsets based on the resolution of the anodes reveals two distinct unimodal trends. This may allude to the possibility of asymmetrical hemispheres (including the spoiler section) along the azimuthal direction. For SPAN-B, the energy offsets between different attenuator states are less consis-

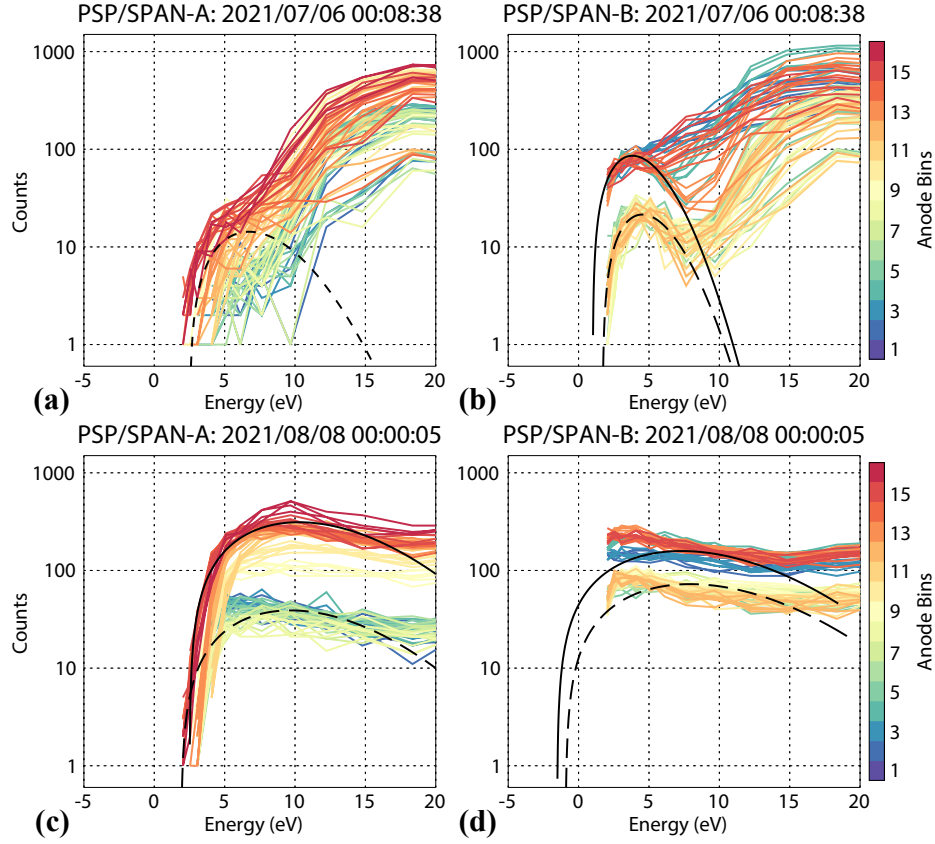


Figure 2.16: SE counts as a function of energy for each anode look direction measured at a time period far from the Sun during orbital encounter 9 for (a) SPAN-A and (b) SPAN-B. The SE counts are also shown for a time interval measured near the Sun during orbital encounter 9 for (c) SPAN-A and (d) SPAN-B. Superimposed Weibull PDF distributions with an energy offset are plotted for the large and small anodes as black solid and dashed curves in each plot. The parameter b is set to 2 for all black curves, with a set to 6, 4, 11, 12.5 for panels (a), (b), (c), and (d), respectively. The energy offset ϵ_0 is set to $[0, 1, 2.5, -1.5]$ and $[2.5, 1.7, 1.9, -1.0]$ for the solid and dashed black curves, respectively. A separate scaling factor was also applied to match the density of the SE population for each case.

tent and show a larger range dependent on the resolution of the anodes. We estimate ϵ_0 for SPAN-A as 1.9 eV from Figure 2.17a since the extrapolation methods converge to a single minimum value of the energy offset. We utilize one final method to confirm the energy shift for SPAN-B using a least-squares fitting technique to the core electron population.

After shifting the measured energies by 1.9 eV for SPAN-A, we can perform a combined fit to the core electrons described in a previous section using both SPAN-A and SPAN-B data. The single alteration we make to equation 2.10 is introduce another fitting parameter

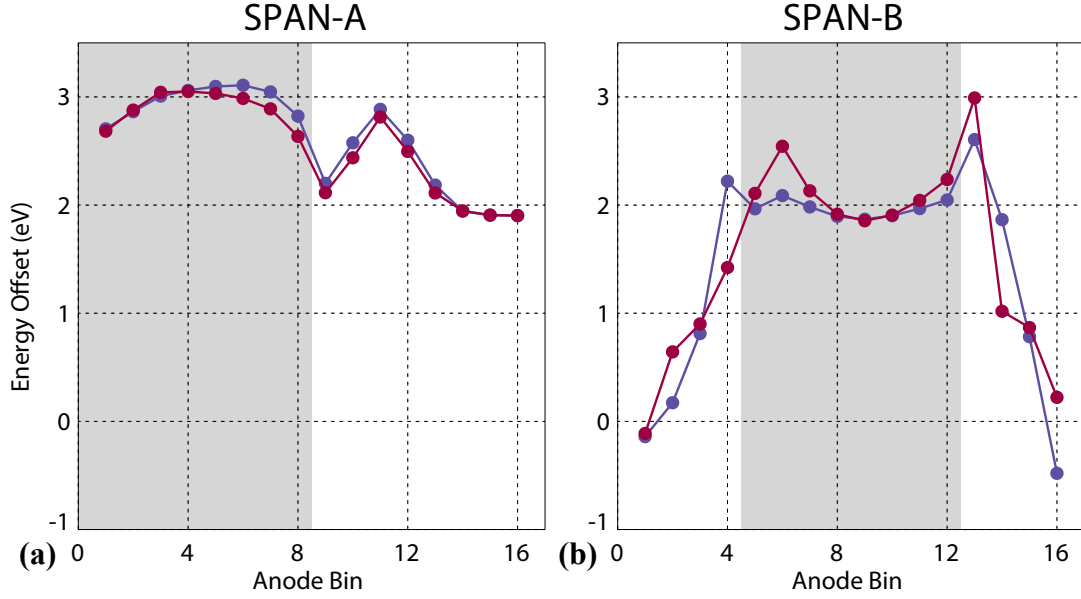


Figure 2.17: Average instrumental energy offsets for each anode determined for (a) SPAN-A and (b) SPAN-B using linear extrapolation, averaged over all orbital encounters. The curves in each panel are computed from measurements when the attenuator was disengaged (purple) and engaged (red). The grey shaded area denotes the eight anodes with 6° resolution.

$\Delta\epsilon$, which represents the relative energy offset between SPAN-A and SPAN-B. The average $\Delta\epsilon$ from the fit to all measurements between encounters 4 through 9 can then be considered as the value of ϵ_0 for SPAN-B. After removing fits with a value of reduced χ^2 less than 1.5, and only considering cases where R_T is near 1 within 5%, we obtain a distribution of SPAN-B energy offsets in Figure 2.18. The weighted average of the histogram is roughly 0.5, depicted by the orange vertical line. Therefore, we estimate ϵ_0 for SPAN-B as 0.5 eV, which is within the range of extrapolated energy offsets in Figure 2.17b.

Anode Efficiencies

After adjusting the energy steps of SPAN-A and SPAN-B, we can determine the absolute anode detector efficiencies for both instruments. Typically, the relative efficiencies of all anodes for an ESA can be calculated based on the measurement of the core distribution in all look directions. The core electrons are expected to be nearly isotropic ($R_T \sim 1$), meaning the peak core differential energy flux should be roughly equal for all look directions in each instrument FOV. Therefore, the relative efficiencies can be computed by matching the differential energy fluxes to an average of all look directions in cases of isotropic distributions. This method does not rely on another external instrument (SPAN-I or FIELDS) to calibrate the efficiencies.

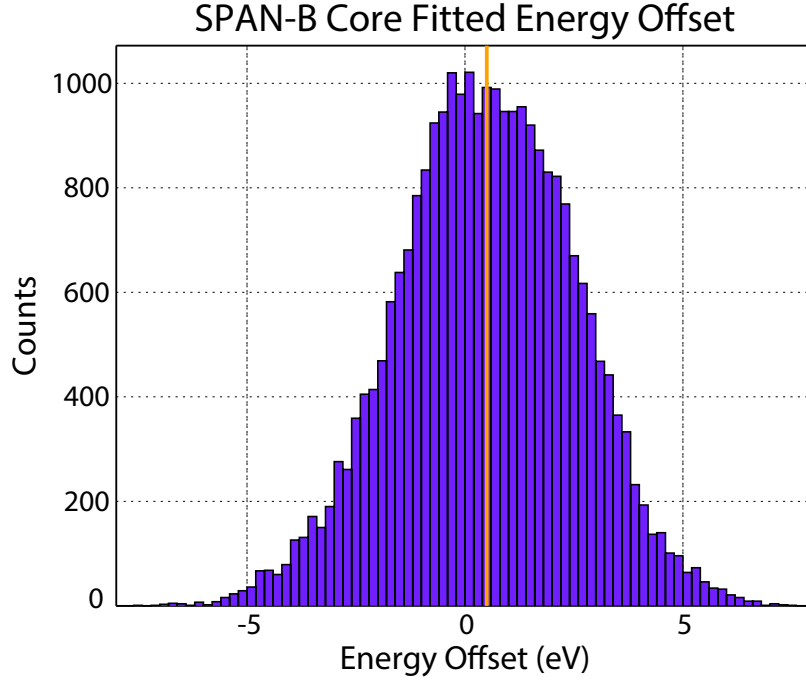


Figure 2.18: Histogram of SPAN-B energy offsets from Maxwellian distribution fitting to the core electron population using both SPAN-A and SPAN-B measurements. The SPAN-A energy offset was set to 1.9 eV, where the average energy offset (orange vertical line) for SPAN-B was determined as 0.5 eV.

However, once the relative anode efficiencies are determined, the ESA may still measure an absolute electron density that does not match the true density in the solar wind. In this case, we can determine absolute anode efficiencies based on solar wind parameters measured from other instruments on board PSP. To determine the absolute value, we first fit the core distribution using the method described earlier using the electron density n_{QTN} derived from Quasi-Thermal Noise theory (QTN; Moncuquet et al. 2020) and the solar wind bulk velocity from SPAN-I. We limit our fits to measurements where $n_{QTN} < 1000 \text{ cm}^{-3}$ (due to dead time effects - described in the next section), $|R_T - 1| < 0.2$ and $T_c < 120 \text{ eV}$ (for colder, isotropic core electrons), and well-measured velocities from SPAN-I (described in Appendix 2.8). The resulting fit parameters are $T_{c||}$, R_T , and the absolute anode efficiencies.

Once the core distribution is fitted, we implement the eVDF hybrid method by computing the plasma moments of the suprathermal component from the residual between the data and core electron model. The density n_s can then be determined solely for the suprathermal population. It is important to note that the density derived from QTN analysis of FIELDS data is the total electron density, not the core density. Thus, we must determine the scale factor s_a for the anode efficiencies based on the relation $n_{QTN} = s_a(n_c + n_s)$. Once the scale

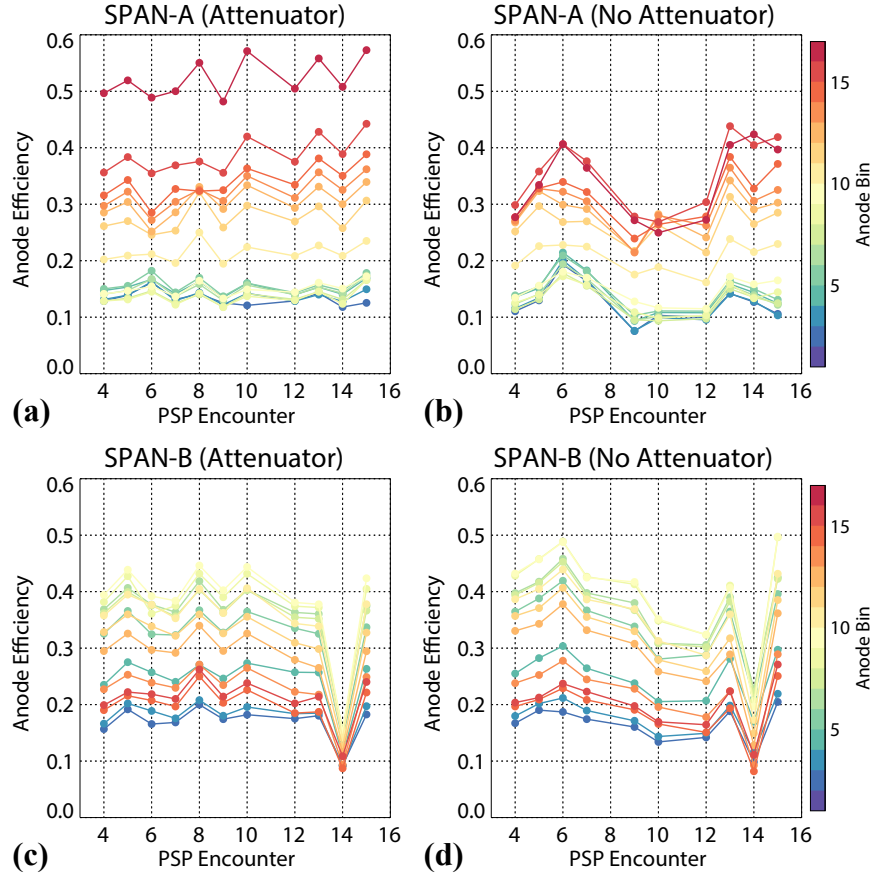


Figure 2.19: SPAN-A anode efficiencies as a function of orbital encounters for 15 anodes during intervals when the attenuator is (a) engaged or (b) disengaged. SPAN-B anode efficiencies as a function of orbital encounters for 14 anodes during intervals when the attenuator is (c) engaged or (d) disengaged.

factor is calculated for each measurement, we can apply this value to the fitted efficiencies to obtain the final absolute anode efficiencies. We limit our fits to cases where the reduced χ^2 is less than 2.0, and avoid in-flight calibration tests, such as spoiler or MCP testing intervals.

Figure 2.19 shows the absolute anode efficiencies for SPAN-A and SPAN-B as a function of PSP orbital encounters during intervals when the attenuator is engaged or disengaged. The outer anode efficiencies are set to one due to issues with the outer FOV look directions. In general, the detectors do not appear to significantly degrade over time. A sharp decrease in efficiency is observed for SPAN-B during encounter 14 due to an accidental change in the MCP gain. For SPAN-A, the range of efficiency estimates is larger when the attenuator is disengaged. Most anode efficiencies are between 0.1 and 0.6, making this a crucial step for accurately determining the electron density and other parameters.

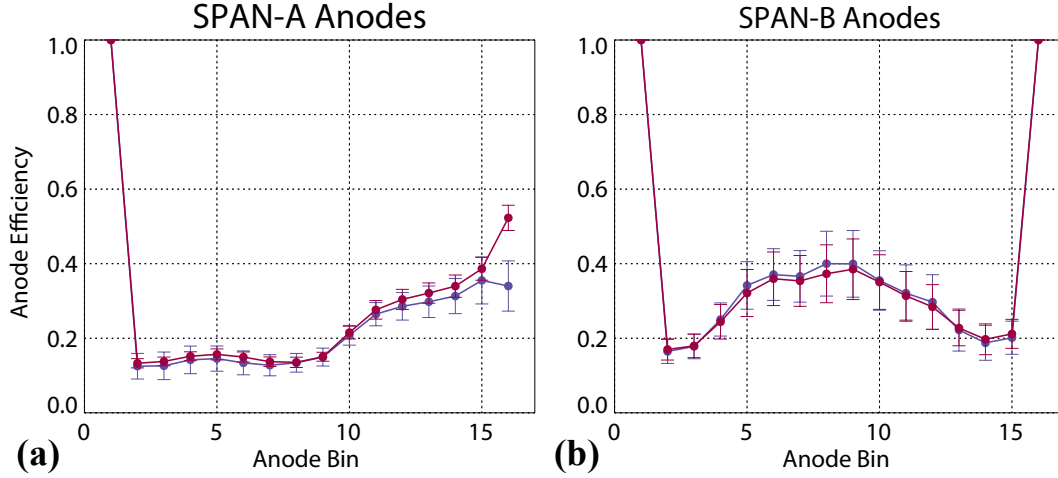


Figure 2.20: (a) SPAN-A anode efficiencies when the attenuator is engaged (red) and disengaged (purple). (b) SPAN-B anode efficiencies when the attenuator is engaged (red) and disengaged (purple). The vertical error bars are the standard deviation based on the average efficiency across all encounters.

The anode efficiencies are then averaged over all encounters for each instrument and attenuator state, shown in Figure 2.20. The efficiencies tend to agree for both states of the attenuator, except for anode 16 of SPAN-A. In addition, the error bars in calculating the average efficiencies are larger when the attenuator is disengaged. After implementing these values in the geometric factor when converting from units of counts to differential energy flux or distribution function, we can repeat the core distribution fits to determine a total electron density from the eVDF hybrid method. Figure 2.21 shows a 2D histogram of the total electron density determined by the eVDF hybrid method (n_e) and QTN (n_{QTN}) for PSP orbital encounter 4. For densities below 1000 cm^{-3} , the majority of the data points lie along the $n_e = n_{QTN}$ line. However, for larger densities, SPAN-E is observed to underestimate the density compared to the QTN method. One possible explanation may be related to the dead time of the ESAs on board PSP, addressed in the next section.

Dead Time Corrections

The final calibration step discussed in this chapter describes the dead time correction when an ESA measures a certain number of counts within a given time interval. The electronics of the instrument is limited to detecting one particle event at a time. Thus, the intrinsic ability for the instrument to detect new particle events due to the processing time must be corrected to ensure an accurate estimate of particle counts, especially in cases of high flux environments. If uncorrected, dead time causes an underestimation of particle flux, as the detector fails to register some incoming events.

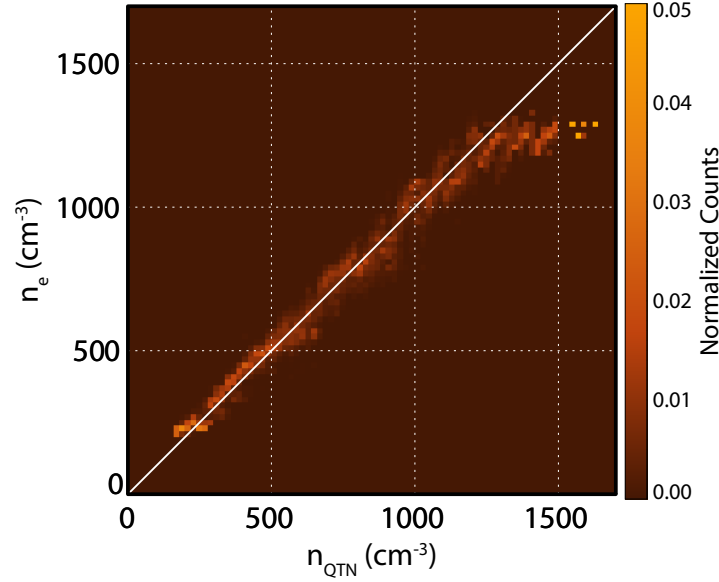


Figure 2.21: 2D histogram of the total electron density determined by the eVDF hybrid method (n_e) and QTN (n_{QTN}) for two days near perihelion in orbital encounter 4, where the counts are column normalized. The white line portrays where $n_e = n_{QTN}$.

For a Poisson process, the interval density $f(t)$ describes the probability density of the time t between successive counts, following an exponential distribution (Müller 1991) given by

$$f(t) = \rho \exp[-\rho t], \quad t > 0 \quad (2.28)$$

where ρ is the count rate, related to the counts C per a sample time Δt by the expression $\rho = C/\Delta t$. In the case of dead times for detectors, there are generally two types: non-paralyzable (non-extending) and paralyzable (extending) dead times. Non-paralyzable detectors have a fixed dead time interval when counts are lost between measured samples, leading to instrument saturation for high count rates. Instead, paralyzable detectors can reset the dead time interval with each count (both measured and undetected samples), extending the overall dead time interval and possibly leading to a cascade of delayed measurements for high count rates. It is possible for detectors to exhibit some combination of both types of dead time, but we will only discuss them individually to quantify their effect on the measured counts.

We can estimate the dead time τ for a non-paralyzable and paralyzable detector to calculate the measured number of counts C_M observed in the solar wind from the actual number of counts C for an integration time Δt (Müller 1991) using the equations

$$C_M = \begin{cases} \frac{C}{1 + C(\tau/\Delta t)} & , \quad \text{Non-paralyzable} \\ C \exp[-C(\tau/\Delta t)] & , \quad \text{Paralyzable} \end{cases} \quad (2.29)$$

However, the ESA instrument provides the number the measured counts C_M with some dead time τ . Re-writing equation 2.29 to compute the actual number of counts C (using a series expansion on the exponential term for the paralyzable case), we obtain

$$C = \begin{cases} \frac{C_M}{1 - C_M(\tau/\Delta t)} & , \quad \text{Non-paralyzable} \\ C_M \left[1 - C_M(\tau/\Delta t) \sum_{k=0}^{\infty} \frac{[k C_M(\tau/\Delta t)]^k}{(k+1)!} \right]^{-1} & , \quad \text{Paralyzable} \end{cases} \quad (2.30)$$

We apply the eVDF hybrid method to determine τ in the non-paralyzable case by first fitting to the core distribution with the density and velocity estimate from QTN and SPAN-I, respectively. The free parameters are the core temperature variables and the dead time τ , where τ is implemented when converting the distribution function into number of counts. We limit the measurements to cases where the density is above 1000 cm^{-3} and the computed suprathermal component is less than 5% of the total electron density. After applying the hybrid method to a few case studies in encounter 4, we determine an initial estimate of the dead time correction near $0.25 \text{ } \mu\text{s}$.

2.4 Electron Strahl Fitting

After the above calibration steps are implemented, we can characterize the strahl direction and angular width measured by SPAN-A and SPAN-B. To determine these properties, we fit a 2D quantized Gaussian function to electron FOV distributions for each energy above 100 eV, as shown in Figure 2.22a. The quantization accounts for the number of counts collected within a certain bin size in θ and ϕ . In addition, we fit to the full FOV in elevation and azimuth instead of pitch angle distributions to avoid data smoothing and lower resolution when converting the 2D FOV into pitch angle. A 2D Gaussian function in terms of θ and ϕ can be described by

$$f(\theta, \phi) = H \exp \left[- \left(\frac{\theta - \theta_0}{\sigma_\theta} \right)^2 \right] \exp \left[- \left(\frac{\phi - \phi_0}{\sigma_\phi} \right)^2 \right] \quad (2.31)$$

where H is the height of the function, θ_0 and ϕ_0 are the mean values, and σ_θ and σ_ϕ are the widths of the Gaussian along the elevation and azimuth. To account for the bin sizes $\Delta\theta$ and $\Delta\phi$ in each look direction, we can define the bin edges θ_1 , θ_2 , ϕ_1 , and ϕ_2 as

$$\begin{aligned} \phi_1 &= \phi - \Delta\phi & \theta_1 &= \theta - \Delta\theta \\ \phi_2 &= \phi + \Delta\phi & \theta_2 &= \theta + \Delta\theta \end{aligned} \quad (2.32)$$

where $\theta_1 < \theta_2$ and $\phi_1 < \phi_2$. To quantize equation 2.31, we can integrate over each bin to obtain $f_Q(\theta, \phi)$ as

$$f_Q(\theta, \phi) = \frac{H}{\Delta\theta\Delta\phi} \int_{\theta_1}^{\theta_2} \exp \left[- \left(\frac{\theta - \theta_0}{\sigma_\theta} \right)^2 \right] d\theta \int_{\phi_1}^{\phi_2} \exp \left[- \left(\frac{\phi - \phi_0}{\sigma_\phi} \right)^2 \right] d\phi \quad (2.33)$$

We then perform a change of variables on equation 2.33 given by

$$f_Q(\Theta, \Phi) = \frac{H\sigma_\theta\sigma_\phi}{\Delta\theta\Delta\phi} \int_{\Theta_1}^{\Theta_2} \exp [-\Theta^2] d\Theta \int_{\Phi_1}^{\Phi_2} \exp [-\Phi^2] d\Phi \quad (2.34)$$

where

$$\begin{aligned} \Phi &= \frac{\phi - \phi_0}{\sigma_\phi} & \Theta &= \frac{\theta - \theta_0}{\sigma_\theta} \\ d\Phi &= \frac{d\phi}{\sigma_\phi} & d\Theta &= \frac{d\theta}{\sigma_\theta} \\ \Phi_1 &= \frac{\phi_1 - \phi_0}{\sigma_\phi} & \Theta_1 &= \frac{\theta - \Delta\theta}{\sigma_\theta} \\ \Phi_2 &= \frac{\phi_2 + \phi_0}{\sigma_\phi} & \Theta_2 &= \frac{\theta + \Delta\theta}{\sigma_\theta} \end{aligned} \quad (2.35)$$

we can next define the $\text{erfc}(x)$ function as the complementary error function, defined by

$$\text{erfc}(x) = 1 - \text{erf}(x) = \frac{2}{\sqrt{\pi}} \int_x^\infty \exp [-s^2] ds \quad (2.36)$$

In addition, we can use the bounds property for definite integrals

$$\int_a^\infty f(x)dx = \int_a^b f(x)dx + \int_b^\infty f(x)dx \quad (2.37)$$

where $a < b < \infty$, or in our case: $\Theta_1 < \Theta_2 < \infty$ and $\Phi_1 < \Phi_2 < \infty$. Utilizing equations 2.36 and 2.37, we can rewrite the 2D quantized Gaussian function as

$$\begin{aligned} f_Q &= \mathcal{H} \left[\int_{\Theta_1}^\infty \exp [-\Theta^2] d\Theta - \int_{\Theta_2}^\infty \exp [-\Theta^2] d\Theta \right] \left[\int_{\Phi_1}^\infty \exp [-\Phi^2] d\Phi - \int_{\Phi_2}^\infty \exp [-\Phi^2] d\Phi \right] \\ &= \frac{\mathcal{H}\pi}{4} [\text{erfc}(\Theta_1) - \text{erfc}(\Theta_2)] [\text{erfc}(\Phi_1) - \text{erfc}(\Phi_2)] \end{aligned} \quad (2.38)$$

where we define the new fitted height $\mathcal{H}$ of the 2D quantized Gaussian function as

$$\mathcal{H} = \frac{H\sigma_\theta\sigma_\phi}{\Delta\theta\Delta\phi} \quad (2.39)$$

It is important to note that Φ in this context does not represent the spacecraft potential, but rather a change of variables (unfortunately, there is only so much symbolic juggling we can do before it turns into a circus act). From equation 2.38, we can find the strahl beam direction $\hat{\mathbf{s}} = (\theta_0, \phi_0)$ by locating the peak of the function in the θ_0 and ϕ_0 directions. Furthermore, the strahl beam width can be determined from the widths σ_θ and σ_ϕ of the quantized function along the elevation and azimuth. A combined width σ can be calculated by averaging the two widths in each direction.

For each measured electron energy, the strahl FOV distributions are fitted separately for SPAN-A and SPAN-B. The strahl tends to be confined primarily to one of the instrument FOVs at a given time. Figure 2.22b presents a fitted 2D quantized Gaussian function overlaid on a FOV distribution measured at 314.5 eV by SPAN-A. For each measurement, we select the final strahl parameters from either the SPAN-A or SPAN-B FOV fit. We first determine whether the strahl direction $\hat{\mathbf{s}}$ is aligned either parallel or anti-parallel to the magnetic field $\mathbf{B}$ within 15° . Both the strahl direction and preferred magnetic field direction ($\mathbf{B}$ or $-\mathbf{B}$) must then be oriented within either the FOV of SPAN-A or SPAN-B. If the magnetic field direction is outside both instrument FOVs, then the fit is disregarded. The 4 Hz magnetic field data from FIELDS is averaged over the SPAN-E cadence, where time intervals with large $\mathbf{B}$ fluctuations (>25 nT) within each SPAN-E measurement are excluded to ensure an accurate estimate of $\mathbf{B}$ and the strahl fit. We then select either the SPAN-A or SPAN-B fits based on the above criteria, to analyze the strahl in relation to other solar wind conditions.

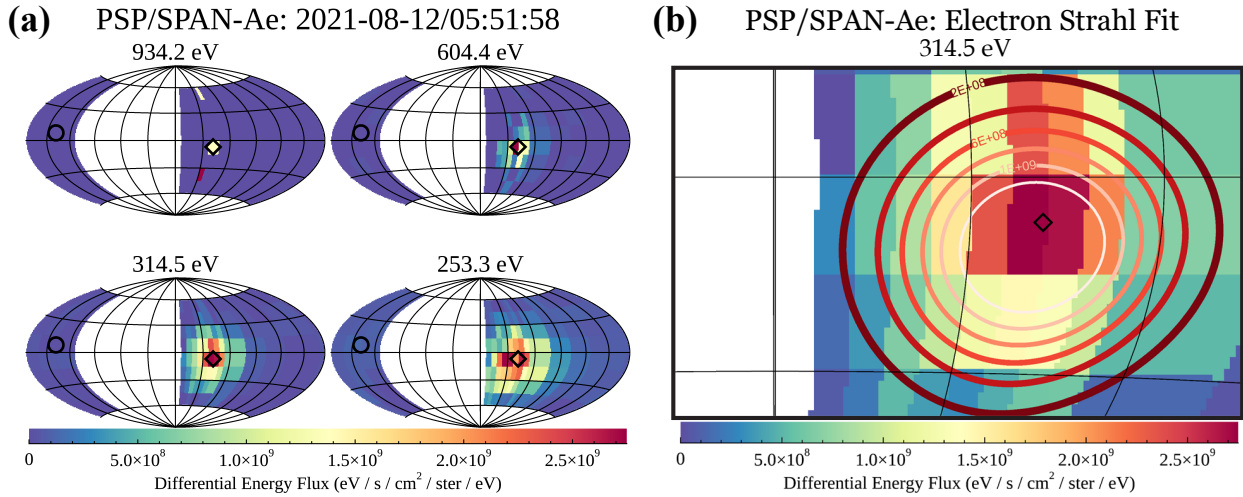


Figure 2.22: (a) SPAN-A differential energy flux measurements collected on 2021/08/12 at 05:51:58 UT for electron energies near 934.2, 604.4, 314.5, and 253.3 eV. The black diamond and circle represent the parallel and anti-parallel directions to the magnetic field vector. (b) SPAN-A FOV map for 314.5 eV electrons with a fitted strahl beam distribution (contour lines), directed toward the magnetic field direction.

2.5 Results

Figure 2.23 presents a time series of PSP data during orbital encounter 12, when the spacecraft reached a perihelion distance of $13.3 R_S$ from the Sun. The second panel shows the angular distance r_{HCS} from the HCS in degrees, indicating a HCS crossing on 02 June 2022. The HCS position is calculated based on potential field source surface (PFSS) modeling (Badman et al. 2020). When the spacecraft crosses this boundary, the magnetic field direction reverses, typically accompanied by a simultaneous reversal in the ePADs. The differential energy flux for each ePAD is normalized to the average value in the third panel. The combined strahl width is observed to increase near the HCS, but is not mutually exclusive to these events. We have also computed the electron collisional age A_e and electron plasma beta $\beta_e = 4.027 \times 10^{-11} n_e T_e / (10^{-10} B^2)$, where the magnetic field is in units of nanotesla. We can also compare the magnitude of the fitted electron bulk velocity with the SPAN-I proton bulk velocity and the Alfvén speed ($v_A = 10^{-10} B / \sqrt{4\pi n m}$) in units of km s^{-1} .

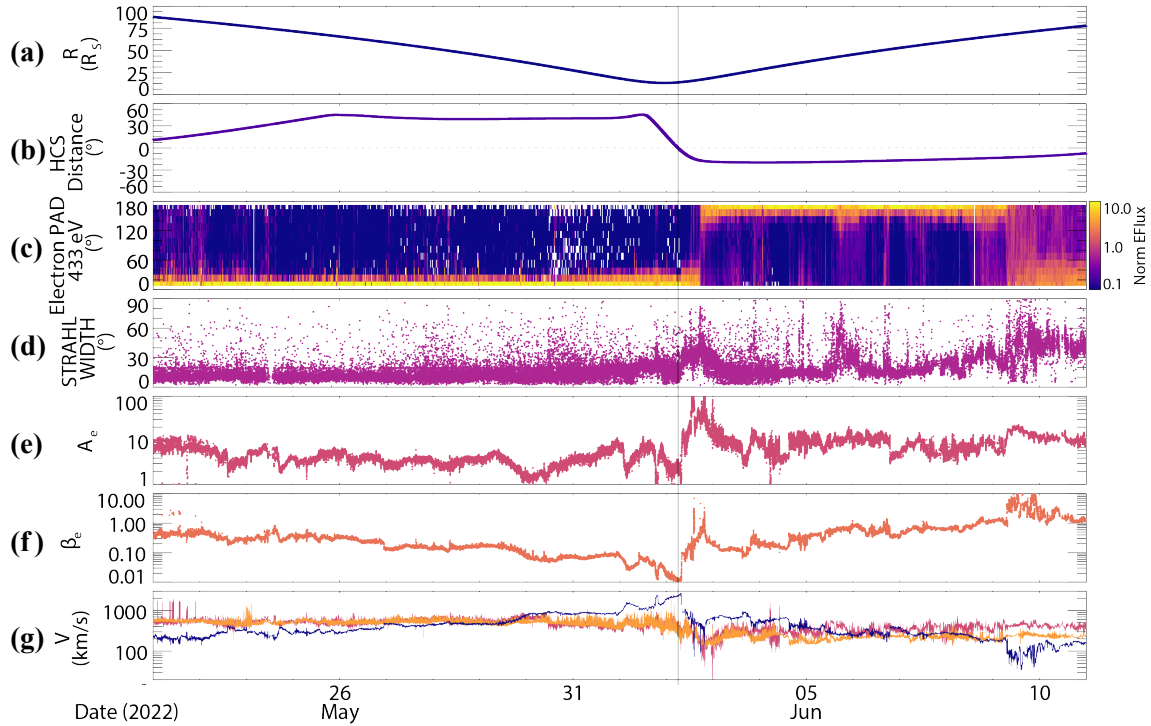


Figure 2.23: Time series of PSP orbital encounter 12. (a) heliocentric distance in solar radii, (b) r_{HCS} computed from PFSS, (c) ePADs measured at 433 eV with normalized differential energy flux, (d) averaged strahl angular width, (e) electron collisional age, (f) electron plasma beta, and (g) SPAN-E bulk velocity (pink), SPAN-I velocity (orange), and Alfvén velocity (blue). The vertical line denotes the PSP HCS crossing, according to PFSS.

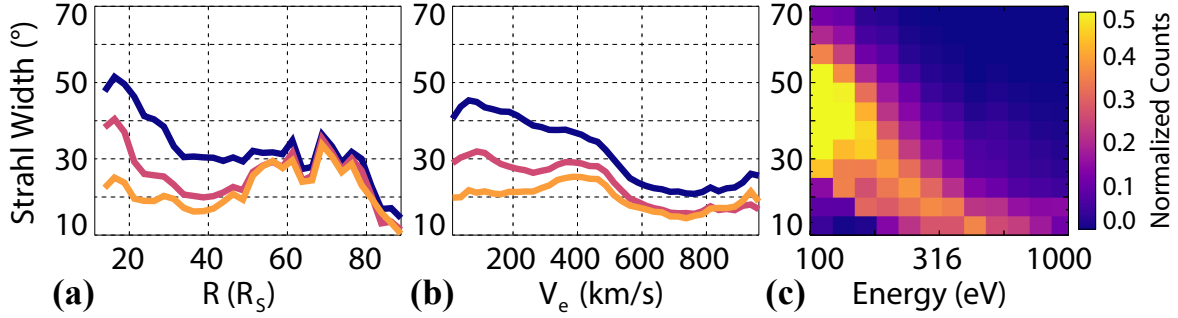


Figure 2.24: Average strahl width for all PSP orbital encounters 4 through 12 as a function of (a) heliocentric distance, (b) SPAN-E bulk velocity and (c) measured electron energy. The three curves in (a) and (b) denote energies near 200 (blue), 300 (pink), 450 (orange) eV. The counts in (c) are normalized by each energy bin. The dataset is composed of samples ± 10 days from each encounter perihelion.

In addition, we can plot the average strahl width as a function of other solar wind parameters, portrayed in Figure 2.24. As the energy tables of the ESAs may change from encounter to encounter, we compute an average strahl width for energies near 200 (blue), 300 (pink), 450 (orange) ± 50 eV. The strahl width is shown to decrease with radial distance from the Sun until roughly $35 R_S$, after which the distribution begins to broaden. The measurements beyond $80 R_S$ appear to trend toward narrow beams, but the sampling size greatly decreases at these heliocentric distances. For later encounters, PSP begins its cruise phase with limited SPAN-E measurements at these distances. In addition, the strahl width tends to decrease with faster solar wind speeds measured by SPAN-E, where a cutoff speed forms near the boundary between the slow and fast solar wind types. Figure 2.24c depicts a 2D histogram of the strahl width as a function of electron energy, with column normalized counts. The general trend is that the strahl beam narrows with increasing energy from 100 to 1000 eV. Furthermore, the strahl beam tends to broaden with increasing total electron density, temperature, and magnetic field strength (not shown).

As strahl scattering could be influenced by multiple mechanisms and solar wind conditions, we analyze the average strahl width in terms of two other solar wind parameters. Figure 2.25a depicts the strahl width organized by collisional age and electron plasma beta, where wider strahl beams are observed for $A_e > 10$. For smaller collisional ages, the strahl width is minimized for measurements with $0.1 < \beta_e < 1.0$. Figure 2.25b shows the electron bulk speed as a function of plasma beta, with a strahl width minimum region for $0.1 < \beta_e < 1.0$ and fast speeds. The strahl beam tends to widen for speeds slower than 500 km s^{-1} for all values of β_e . Figure 2.25c plots the electron bulk speed as a function of Alfvén speed, where the strahl width measured during sub-Alfvénic intervals ($v_e < v_A$) tend to remain near 45° . When $v_e > v_A$, the strahl width is observed to vary from 10 to 60° . For super-Alfvénic intervals, the strahl distribution broadens for smaller values of v_A . The final

panel portrays the strahl width as a function of electron bulk speed and collisional age. The strahl width is measured to increase with larger collisional ages and slower bulk speeds.

2.6 Discussion

SPAN-A and SPAN-B on board PSP offer valuable measurements for studying electron dynamics near the Sun, but require thorough calibration to ensure accuracy of solar wind properties. The constant energy shift ϵ_0 determined for each ESA was found to depend on anode. This may be a result of other asymmetries in the nested hemispheres, including the spoiler section. Another possibility is that a multiplicative factor κ to the energy steps is needed to align the measured electron energies between the two instruments. Near perihelion, this small energy shift does not greatly affect the estimate of the core distribution temperatures as T_c is typically greater than 50 eV. However, the core electrons can be colder

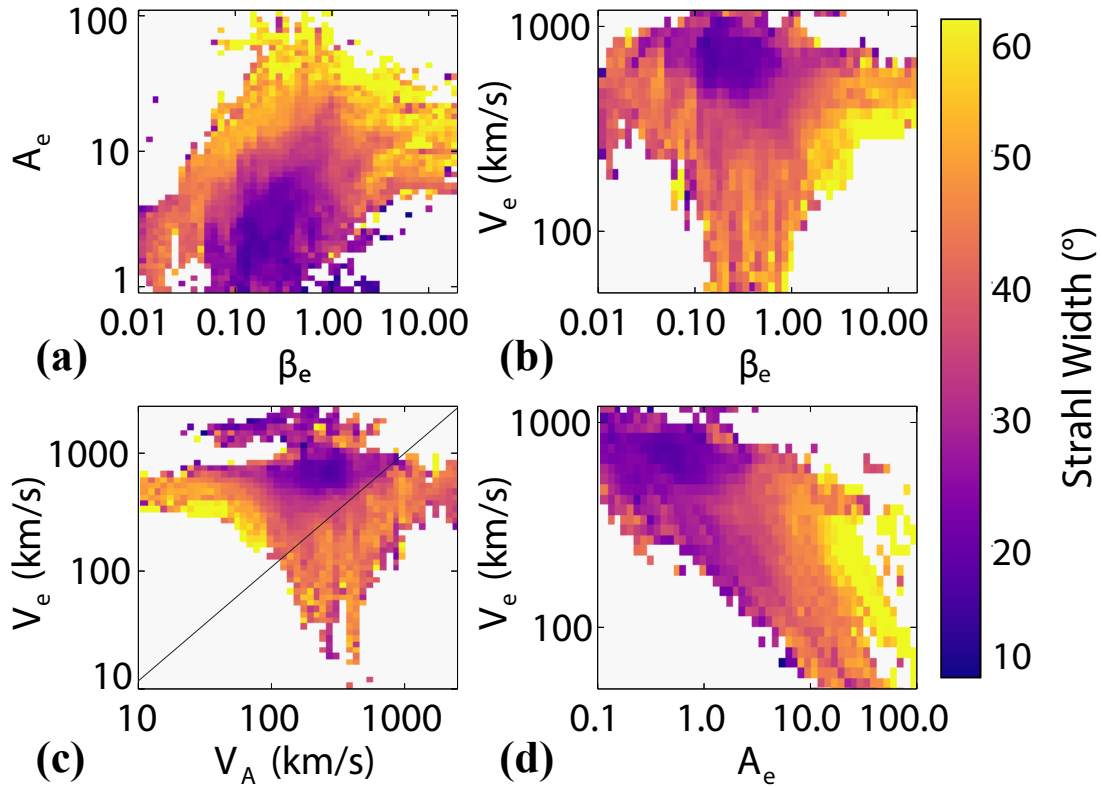


Figure 2.25: Averaged strahl angular width (color bar) for PSP orbital encounters 4 through 12 as a function of (a) collisional electron age versus electron plasma beta, (b) electron bulk velocity versus electron plasma beta, (c) electron bulk velocity versus Alfvén velocity, and (d) electron bulk velocity versus collisional electron age for electrons measured near 300 eV.

for intervals outside of perihelion, and require an accurate estimate of the offset in energy. This is especially true for analyzing electron distributions during Venus gravity assists when the core temperature is around 10 eV or less. Future work could implement both ϵ_0 and κ as a function of anode to combine electron measurements from both instruments.

Furthermore, we determined the relative and absolute anode efficiencies to estimate the total electron density measured by SPAN-A and SPAN-B. The absolute efficiencies were found to vary between 10% and 60%, leading to an underestimation of the electron density without this consideration. It is important to note that the spacecraft floating potential (Diaz-Aguado et al. 2021) was not included in these calculations. Liu et al. (2024, *in preparation*) is working to determine an estimate of the calibrated spacecraft floating potential based on the antenna voltages from the FIELDS suite and photoelectron current. A simple estimate of the spacecraft potential can be obtained from averaging the negative of the biased antenna floating potential. However, the photoelectron current must be included to ensure all currents in the solar wind are accurately balanced. Future work could refine the analysis presented in this chapter using the spacecraft potential to improve the agreement between the electron density from SPAN-E and QTN.

We are also interested in better understanding the SE response for each instrument. Currently, we fit a Weibull PDF distribution to the SE population separately for each instrument. The presence of more strahl electrons at higher energies lead to more anisotropy in the SE distribution, where the look direction with the largest SE flux is correlated with the flux in the corresponding pitch angle direction for the strahl. However, the density of the SE population is still estimated near n_c even when the strahl density is less than 5% of the total electron density. Future work could model this distribution of low electrons through ray-tracing simulations, assuming a certain yield response of SE creation from impacting electrons within each instrument. Once the SE electrons are modeled, we can remove the distribution from the measured counts to obtain a more accurate estimate of the plasma moments across all energies.

After SPAN-A and SPAN-B calibration, we computed the combined strahl angular width as a proxy for strahl beam scattering. The minimum heliocentric distance where the strahl width broadens further from the Sun suggests the dominating mechanisms change below $35 R_S$. Numerous studies with PSP and Solar Orbiter data (e.g., Cattell et al. 2021; Colombari et al. 2024; Halekas et al. 2021; Jagarlamudi et al. 2021) have correlated the occurrence and properties of whistler waves to strahl scattering. However, the absence of these waves observed by Cattell et al. (2022) within roughly $28 R_S$ indicates another mechanism must scatter the strahl electrons. Coulomb collisions could be a viable candidate within $35 R_S$, where higher densities are observed as the density scales roughly with the square of the heliocentric distance. Higher densities do correspond to wider strahl distributions, which could be related to an increase in collisional frequency. It is possible that the driving mechanisms within 0.15 AU could primarily be Coulomb collisions and adiabatic focusing by the magnetic field. The magnetic focusing of the strahl beam would overcome the collisions with distance from the Sun until about $\sim 35 R_S$ where particle instabilities scatter the strahl. Once PSP travels within $10 R_S$ of the Sun, we can determine whether this trend continues closer to the

source of the strahl population.

We also observe an inverse relationship of strahl angular width with solar wind velocity and measured electron energy. Previous studies have seen narrow strahl widths for high velocity solar wind streams (Anderson et al. 2012), as well as narrow strahl widths for higher electron energies (Berčič et al. 2020). The fast wind is typically less collisional compared to the slow wind, as the cross sectional area of a collision scales as $\sigma_{CR} \propto v^{-4}$. For higher velocities, particles have to travel closer to the target particles for significant deflection due to the smaller cross sectional area. Therefore, Coulomb collisions may be less likely for high speed streams and result in less scattering of the strahl.

In addition to slower speeds, the strahl distribution is widened for intervals of higher collisional age, supporting our claim for a strong dependence of Coulomb collisions on strahl scattering. The strahl width is well organized in Figure 2.25d for electron bulk velocity and collisional age. For sub-Alfvénic streams, we find the angular width is roughly constant with solar wind speed. These intervals are typically dominated by the magnetic field instead of kinetic effects. We also observe broadened strahl distributions for $\beta_e > 1$ for all observed collisional ages, when the plasma is primarily influenced by thermal effects. For $\beta_e < 1$ and low collisional ages, the strahl is primarily focused by the magnetic field, reaching a minimum of 10° in angular width. For most super-Alfvénic streams, the strahl is broadened for slower solar wind speeds, providing more evidence toward Coulomb collisions as the primary scattering mechanism near the Sun. However, this study did not examine other possible particle instability candidates that may scatter the strahl below $35 R_S$. While the adiabatic focusing could be the likely dominant process within this distance, another mechanism could effectively scatter the strahl besides collisions. Future work could relate the occurrence and properties of solitary waves or ion acoustic waves (Cattell et al. 2022) to strahl scattering to compare its influence versus Coulomb collisions.

2.7 Conclusion

After performing multiple calibration procedures, we are able to relate the amount of strahl electron scattering close to the Sun to various solar wind parameters. Due to the absence of observed narrow band whistler waves within $28 R_S$ of the Sun, our study suggests that Coulomb collisions could be the primary mechanism to scatter the Sun within this heliocentric distance. We observed a minimum of the strahl beam width as a function of radial distance, where magnetic focusing could begin to prevail over Coulomb collisions closer to the Sun. The strahl distribution broadens for higher densities, slower solar wind speeds, and lower electron energies. After PSP completes its 22nd orbital encounter within $10 R_S$ of the Sun, we can better understand the degree of Coulomb collision effects closer to the origin of the strahl.

Furthermore, future work can refine the method of calculating the strahl angular width measured by SPAN-A and SPAN-B. This study fitted the strahl beam distribution separately for each instrument, but a combined fit between both ESAs could provide a better estimate of

strahl scattering in extreme cases. In addition, we hope to improve these fits to independently determine the magnetic field direction solely from the strahl direction. During time intervals of a narrow strahl beam measured by the smaller resolution anodes, the SPAN-E estimate of the magnetic field unit vector could provide intercalibration with *FIELDS*. Together with *FIELDS* data, the ePADs of strahl electrons can help determine magnetic topology in the heliosphere. The ePADs can allow us to identify whether the field lines connected to the Sun are open or closed, based on the bidirectionality of the strahl, which is particularly useful in cases of HCS crossings (Szabo 1994) and CME flux ropes (Gosling et al. 1987). Thus, improving our method to combine and fit the instrument data can provide valuable insights to other applications in the solar wind.

2.8 Appendix: SPAN-I FOV Coverage

To calibrate the SPAN-E instrument, we utilize solar wind density and velocity derived from SPAN-I measurements. However, the SPAN-I solar wind parameters are only reliable when the majority of the proton distribution is within the FOV of the instrument. Therefore, we must find a proxy to measure the FOV coverage of the protons for SPAN-I to properly calibrate the electron instrument on board PSP. As a basic approach, we sum the particle counts measured by SPAN-I across all the front anode bins to obtain counts as a function of ϕ . Similarly, we sum the counts across all deflection bins to obtain counts as a function of θ . Figure 2.26 shows the converted differential energy flux as a function of ϕ and θ measured by SPAN-I on 2021/08/08 at 23:59:59 UT.

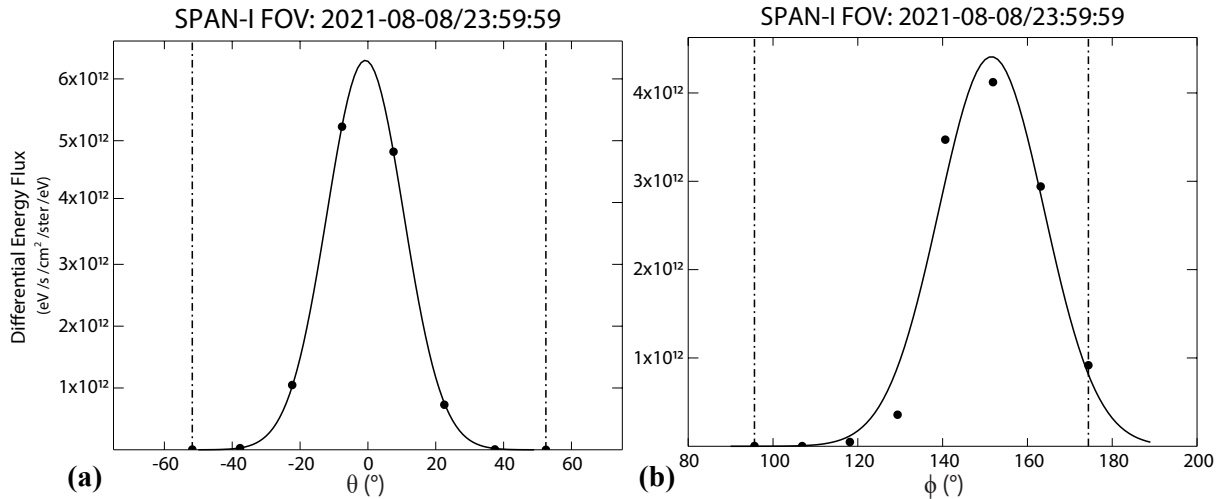


Figure 2.26: SPAN-I differential energy flux averaged over the (a) θ and (b) ϕ angles on 2021/08/08 at 23:59:59 UT. A 1D Gaussian distribution (black curve) is fitted to the SPAN-I data (black points), where the vertical dashed lines indicate the limit of the instrument FOV.

We then fit the summed data in units of differential energy flux J_E as a function of each angle to a 1D Gaussian function. Figure 2.26 depicts a fitted Gaussian function for each angle within the limits of the instrument FOV, described by

$$f(x) = h \exp \left[- \left(\frac{x - x_0}{w} \right)^2 \right] \quad (2.40)$$

where h is the height of the distribution, w is the width, and x_0 is the center of the Gaussian. From the fit, we can determine the amount of FOV coverage based on the area of the distribution within the FOV range. We can integrate the fitted curve using the cumulative distribution function CDF at the FOV boundaries θ_1 , θ_2 , ϕ_1 , and ϕ_2 , where $\theta_1 < \theta_2$ and $\phi_1 < \phi_2$. The CDF can be calculated using the equation

$$\text{CDF}(z) = \frac{1}{2} \left[1 + \text{erf} \left(\frac{z}{\sqrt{2}} \right) \right] \quad (2.41)$$

where $\text{erf}(x)$ is the error function and z is defined for θ and ϕ as

$$\begin{aligned} z_{\phi_1} &= \frac{\phi_1 - \phi_0}{w_\phi} & z_{\theta_1} &= \frac{\theta_1 - \theta_0}{w_\theta} \\ z_{\phi_2} &= \frac{\phi_2 - \phi_0}{w_\phi} & z_{\theta_2} &= \frac{\theta_2 - \theta_0}{w_\theta} \end{aligned} \quad (2.42)$$

based on the Gaussian fit parameters. The integrated area beneath the fitted curves in ϕ and θ is computed as $A_\phi = \text{CDF}(z_{\phi_2}) - \text{CDF}(z_{\phi_1})$ and $A_\theta = \text{CDF}(z_{\theta_2}) - \text{CDF}(z_{\theta_1})$. The values of A_ϕ and A_θ are referred to as the amount of FOV coverage for each direction, and are multiplied together to determine a combined FOV coverage A_{FOV} . We calculate these parameters for SPAN-I proton measurements between PSP orbital encounters 4 through 19.

As a further constraint, we limit our fits to cases where both the fitted Gaussian peak and the maximum measured value of J_E are within $106^\circ < \phi < 164^\circ$ and $-40^\circ < \theta < 40^\circ$. This ensures that the peak of the Gaussian curve is well within the measured range to estimate the FOV coverage. Previous studies have limited reliable SPAN-I measurements to intervals only when the measured peak of J_E in ϕ is less than about 160° . In most cases, the upper bound on the measured peak of J_E in ϕ is most important when filtering the data. However, we check the entire range of ϕ and θ as we are computing a combined FOV coverage estimate. To determine the limit of the FOV coverage that solar wind parameters can be accurately measured by SPAN-I for instrument calibration, we compare the proton density with the density determined by Quasi-Thermal Noise theory (QTN; Moncuquet et al. 2020). The QTN method provides accurate density estimates above 100 cm^{-3} due to the sensitivity and range of the FIELDS Radio Frequency Spectrometer (RFS) instrument. One method to quantify the accuracy of SPAN-I fits and moments is dependent on the density ratio n_i/n_{QTN} as it approaches near 1.0, where $n_i = n_p + 2n_\alpha$ for proton n_p and α -particle n_α densities derived from moment calculations.

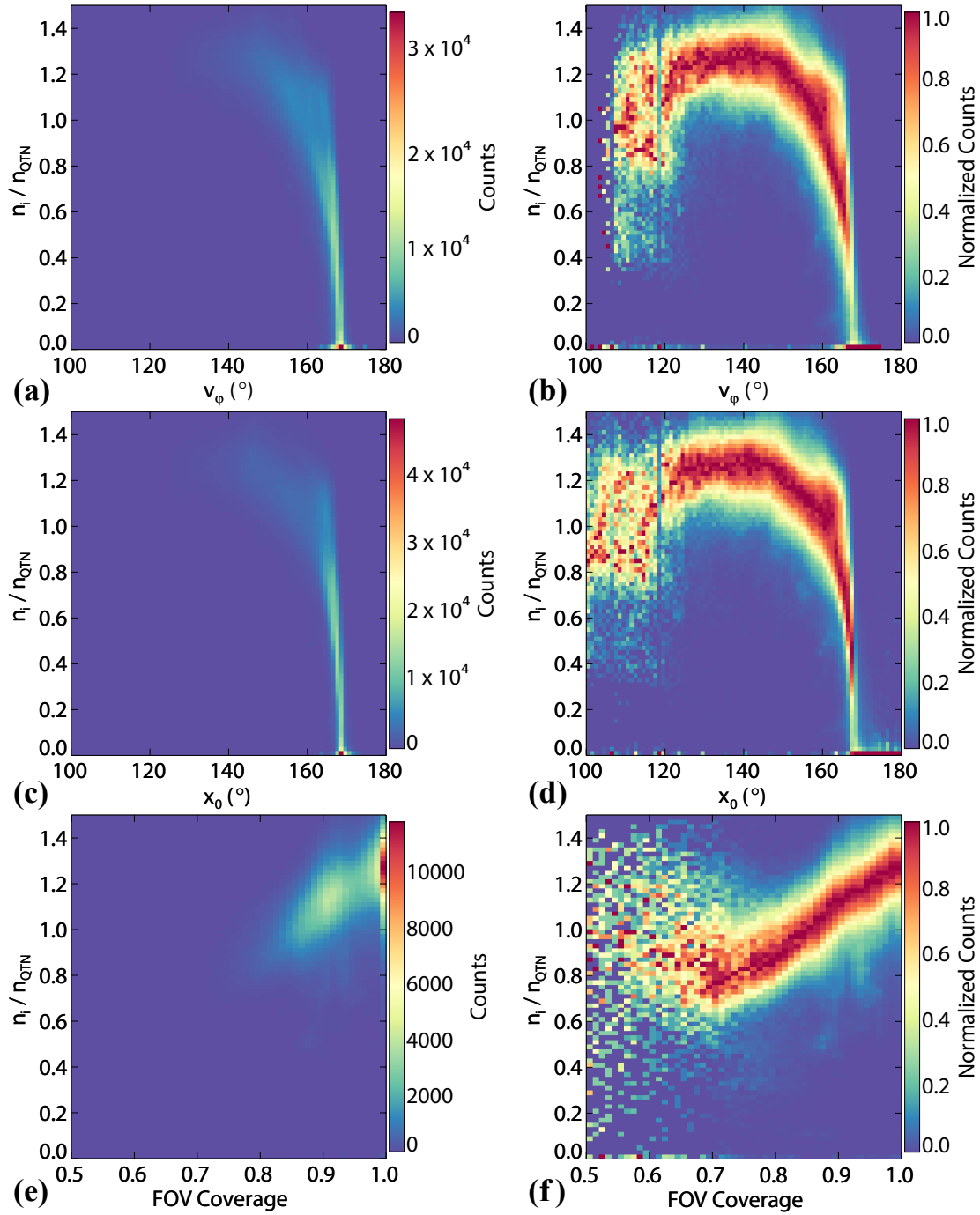


Figure 2.27: 2D histogram of the SPAN-I density to QTN density ratio as a function of the (a) SPAN-I velocity in ϕ (counts), (b) SPAN-I velocity in ϕ (column normalized counts), (c) fitted Gaussian center (counts), (d) fitted Gaussian center (column normalized counts), (e) FOV coverage (counts), (f) FOV coverage (column normalized counts).

We first analyze FOV effects with the SPAN-I fitted velocity direction in ϕ from Livi et al. (2022). The ϕ angle tends to be more important than θ in determining FOV effects on the computed ion distribution fits and moments. Figure 2.27a portrays a 2D histogram of the ratio between SPAN-I proton density and QTN electron density as a function of the SPAN-I fitted velocity v_ϕ . The majority of SPAN-I measurements are collected with v_ϕ near 170° . However, once the histogram is column normalized, shown in Figure 2.27b, the SPAN-I density is within 15% of the QTN density for $v_\phi < 160^\circ$. A similar trend emerges when we instead use the fitted Gaussian center x_0 in Figures 2.27c and 2.27d. The agreement in both trends using v_ϕ or x_0 supports the claim that a simple Gaussian function can accurately model the data to estimate the FOV coverage.

Figure 2.27e portrays a similar 2D histogram, but as a function of the combined FOV coverage A_{FOV} , where the data are clustered near $A_{FOV} \sim 0.9$ and $A_{FOV} \sim 1.0$. The SPAN-I measurements with $n_i/n_{QTN} > 1.0$ for almost complete FOV coverage may be a result of calibrating the SPAN-I instruments to the QTN density for FOV coverage close to 0.9. Figure 2.27f shows a linear trend for $A_{FOV} > 0.75$ with the density ratio when we column normalize the counts. As expected, the agreement between the proton and QTN electron density improves as the FOV coverage increases. This linear trend can be fitted to determine an empirical relationship between the combined FOV coverage and the density ratio, shown in Figure 2.28. The vertical error bars are computed from the column spread in Figure 2.27f. The final fitted equation obtained from Figure 2.28 is given by $n_i/n_{QTN} = 2.83A_{FOV} - 1.53$. Therefore, for SPAN-I densities within 15% of the QTN density, we filter out any samples with a combined FOV coverage A_{FOV} less than 0.865 to calibrate the SPAN-E instrument.

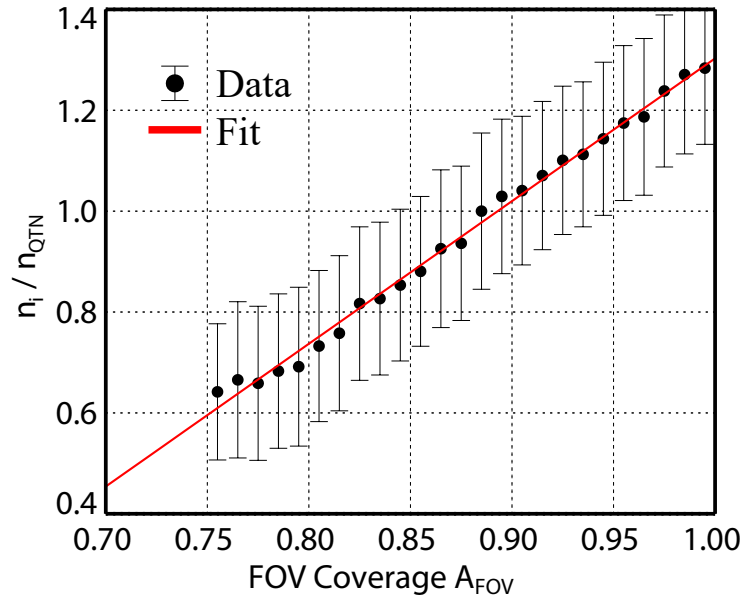


Figure 2.28: The SPAN-I density to QTN density ratio as a function of the combined FOV coverage. The error bars for the data (black) are based on the spread in each column of Figure 2.27f. A linear fit (red) is performed to the data to obtain the empirical relationship $n_i/n_{QTN} = 2.83A_{FOV} - 1.53$.

Chapter 3

Near-Sun In situ and Remote Sensing Observations of a CME and its Effect on the Heliospheric Current Sheet

The published version (Romeo et al. 2023b) of this chapter was coauthored by the following individuals, and is included in this thesis with their express permission: Carlos Braga, Samuel Badman, Davin Larson, Michael Stevens, Jia Huang, Tai Phan, Ali Rahmati, Roberto Livi, Samer Alnussirat, Phyllis Whittlesey, Adam Szabo, Kristopher Klein, Tatiana Niembro-Hernandez, Kristoff Paulson, Jaye Verniero, David Lario, Nour Raouafi, Tamar Ervin, Justin Kasper, Marc Pulupa, Stuart Bale, and Mark Linton.

3.1 Abstract

During the thirteenth encounter of the Parker Solar Probe (PSP) mission, the spacecraft traveled through a topologically complex Interplanetary Coronal Mass Ejection (ICME) beginning on 05 September 2022. PSP traversed through the flank and wake of the ICME while observing the event for nearly two days. The Solar Probe ANalyzer (SPAN) and FIELDS instruments collected in situ measurements of the plasma particles and magnetic field at $\sim 13.3 R_S$ from the Sun. We observe classical ICME signatures, such as a fast-forward shock, bidirectional electrons, low proton temperatures, low plasma β , and high alpha particle to proton number density ratios. In addition, PSP traveled through two magnetic inversion lines, a magnetic reconnection exhaust, and multiple sub-Alfvénic regions. We compare these in situ measurements to remote sensing observations from the Wide-field Imager for Solar PRobe Plus (WISPR) instrument on board PSP and the Sun Earth Connection Coronal and Heliospheric Investigation (SECCHI) on STEREO. Based on white-light coronagraphs, two CMEs are forward modeled to best fit the extent of the event. Furthermore, ADAPT-GONG magnetograms and Potential Field Source Surface (PFSS) modeling portray a global reconfiguration of the Heliospheric Current Sheet (HCS) after the ICME event, suggesting these eruptions play a significant role in the evolution of the HCS.

3.2 Introduction

Interplanetary Coronal Mass Ejections (ICMEs) are macro-scale structures of coronal plasma material and magnetic fields that drive space weather events and shape the heliosphere. ICMEs are one of the most energetic phenomena in our solar system, and originate from the expulsion of Coronal Mass Ejections (CMEs) from the Sun. There have been many proposed trigger mechanisms that eject CMEs from closed field lines on the Sun, such as magnetic reconnection (Gosling 1990). The occurrence rate of these eruptions have been shown to vary with the solar cycle (St. Cyr et al. 2000; Webb et al. 1994), from less than one per day during solar minima, to a few events per day during solar maxima. Although white-light CME observations occur relatively frequently, the number of events that have been detected using in situ measurements is comparatively low. However, recently launched solar missions, such as Parker Solar Probe (PSP; Fox et al. 2016; Raouafi et al. 2023), provide in situ data much closer to the Sun to characterize the structure of CMEs at the early stages of their evolution.

The first spacecraft observations of CMEs can be traced back to the early 1970s, when white-light coronagraphs on board the Orbiting Solar Observatory (OSO) and the Skylab space station revealed the presence of coronal transients (Tousey 1973). With the launch of the Solar Maximum Mission (SMM), the classical three-part CME structure was confirmed to have a bright frontal loop, dark cavity, and a prominence core (Crifo et al. 1983; Illing et al. 1986). However, this structure is not universal, as some CMEs can be more complex. Narrow CME structures have been observed (Howard et al. 1985), either as narrow jets originating from coronal holes or narrow blobs from coronal streamers (Webb et al. 2012). Other structures, such as Halo CMEs (Howard et al. 1982), appear large over a wide angular range as they primarily expand along the Sun-observer line.

In addition to remote sensing observations, a number of in situ signatures have been identified for the detection of ICMEs, as compiled by Zurbuchen et al. (2006). It is important to note that none of these characteristics are necessarily guaranteed to be observed by a single spacecraft traveling through the event, owing to the inherent variability of each ICME (Gosling 1993). A certain group of ICMEs, known as magnetic clouds (MCs; Burlaga et al. 1981), are categorized by regions of smoothly rotating magnetic fields ($\gg 30^\circ$) with a magnitude enhancement of at least 10 nT and low plasma β at 1 AU (Klein et al. 1982). MCs tend to have low proton temperatures T_p , satisfying the condition $T_p < 0.5 T_{exp}$, where T_{exp} is the expected proton temperature based on the solar wind (SW) speed relationship (Gosling et al. 1973; Richardson et al. 1995). A significant difference between the expected and measured proton temperature may indicate solar wind plasma that has been perturbed by an ICME event or has experienced local compression or rarefaction. Cane et al. (2003) defined a quality index for MC identification, such that an index of two is assigned to ICME events that exhibit all typical features of MCs. An index of one indicates evidence of a MC solely based on magnetic field rotation, while an index of zero corresponds to events with no MC features. MCs are composed of a cylindrical bundle of twisted magnetic field lines with a helical structure, known as a flux rope (Burlaga et al. 1981; Klein et al. 1982). Some

MCs have been modeled with an elliptical cross section for the magnetic flux rope topology (Nieves-Chinchilla et al. 2018). While some of these characteristics have been observed separately for other macro-scale structures in the solar wind, such as interplanetary sector boundaries (e.g., Heliospheric Current Sheet (HCS) crossings; Bothmer et al. 1992; Schulz 1973), they tend to simultaneously occur for MCs (Bothmer et al. 1997). Other ICMEs may contain “complex ejecta” without flux ropes, and typically have weaker magnetic fields and higher proton temperatures than MCs (Burlaga et al. 2001). However, MCs may appear as non-rope configurations in the case of a partial observation of the entire magnetic field structure (e.g., a spacecraft traveling through the legs of a flux rope).

Other in situ ICME signatures include directional discontinuities in the magnetic field around ICME boundaries (Janoo et al. 1998), fast-forward interplanetary shocks for fast ICMEs (Manchester et al. 2012; Reames 1999), and less magnetic field variation compared to ICME sheath regions (Klein et al. 1982). At 1 AU, periods of enhanced alpha particle (fully ionized Helium) to proton density ratios ($> 8\%$) (Hirshberg et al. 1972) and low plasma density ($< 1 \text{ cm}^{-3}$) (Richardson et al. 2000) have occurred for ICME events. A monotonic decrease in solar wind velocity has also been reported after a sharp increase near the structure front (Klein et al. 1982). The most common signature, valid for both slow and fast ICMEs, is the observation of bidirectional electrons above 80 eV (Gosling et al. 1987; Nieves-Chinchilla et al. 2008). The electron pitch angle distributions (ePADs) typically show a single strahl beam directed along the magnetic field (Pilipp et al. 1987). The presence of bidirectional electrons (i.e., electrons streaming both parallel and anti-parallel to the magnetic field) may indicate connection to both footpoints on the Sun from closed field line loops within ICMEs.

The scope of this study encompasses a time interval of nearly two days, from 05 September to 07 September 2022, as PSP traversed an ICME at ~ 13.3 solar radii. This event provides an unprecedented opportunity to study plasma and field structures of an ICME in detail at an early stage of its evolution. PSP observed the flank and wake regions of the ICME, providing novel observations compared to measurements near 1 AU. We focus primarily on PSP in situ measurements – the mission does offer remote sensing capabilities that are discussed in greater detail elsewhere (e.g., Vourlidas et al. 2023, *in preparation*). The event included a combination of several interesting features related to classical ICME in situ signatures (described above) that have not previously been observed by PSP. Section 3.3 provides a summary of the in situ data, highlighting important features and time intervals related to the ICME. Section 3.4 presents the CME model from Braga et al. (2023, *in preparation*) along with white-light coronagraph observations. A reference to an animation of the model is provided to synchronize the remote sensing observations, in situ measurements, and modeled CME propagation over time, providing a comprehensive visualization of the event. Section 3.5 uses Potential Field Source Surface (PFSS) modeling and in situ data to investigate the source region and large scale coronal structure before and after the event. Finally, Section 3.6 compares PSP observations with CME modeling and coronal magnetograms to offer insight on the ICME structure and evolution.

3.3 PSP In situ Measurements

Instrumentation

The PSP mission was launched on 12 August 2018 to collect unprecedented in situ measurements of the solar wind close to the Sun (Fox et al. 2016; Raouafi et al. 2023). The main science goals are to study the Sun’s lower coronal structure, solar wind heating and acceleration, and energetic particle acceleration. The spacecraft orbits the Sun in a highly elliptical trajectory, using Venus gravity assists to gradually decrease its perihelion distance from 35.7 to 9.9 solar radii (R_S) over the course of the prime mission. During each encounter phase, the instruments on board PSP collect measurements of solar wind plasma ions and electrons, magnetic and electric fields, energetic particles, and white-light images off-disk. The orbital perihelion for the thirteenth PSP encounter, which is the focus of this chapter, was reached on 06 September 2022, with a closest approach of 13.3 R_S .

For this study, we use data from the Solar Wind Electrons, Alphas, and Protons (SWEAP; Kasper et al. 2016) suite. SWEAP consists of two electron electrostatic analyzers collectively known as Solar Probe ANalyzer for Electrons (SPAN-E; Whittlesey et al. 2020), one ion electrostatic analyzer referred to as the Solar Probe ANalyzer for Ions (SPAN-I; Livi et al. 2022), and a Faraday Cup called the Solar Probe Cup (SPC; Case et al. 2020). SPC was powered off for most of the period of interest due to temperature constraints. Thus, the proton and alpha particle parameters in this study, i.e., number density n , temperature T , and velocity $\mathbf{v}$, are estimated exclusively from the SPAN-I moments. To obtain electron plasma parameters, we follow a similar method from Halekas et al. (2020) with a few minor differences. An anisotropic bi-Maxwellian distribution is fitted to the core electron population along all pitch angles using SPAN-E three dimensional electron velocity distributions. We have re-computed the anode efficiencies for both SPAN-E instruments to match the plasma density calculated from other onboard instruments, discussed below. Below ~ 25 eV, the SPAN-E instruments are contaminated by secondary electrons and photoelectrons. We determine the local minimum between the secondary and core electron distributions, and fit the core from this minimum to the peak of the distribution.

We also utilize Level 2 magnetic field data from the fluxgate magnetometer in Radial-Tangential-Normal (RTN) coordinates at ~ 4 samples per second from the FIELDS suite (Bale et al. 2016). The RTN coordinate system is defined where $\mathbf{R}$ is directed anti-sunward from the spacecraft, $\mathbf{T}$ is tangential to the spacecraft orbit, and $\mathbf{N}$ is normal to the RT plane (Fränz et al. 2002). FIELDS measurements are used to study the background magnetic field structure and calculate ePADs. For comparison with the SPAN derived plasma densities, we calculate the Quasi-Thermal Noise (QTN) total electron density using the Radio Frequency Spectrometer (RFS) from the FIELDS suite, specifically the Low Frequency Receiver (LFR) data from the V1-V2 and V3-V4 antenna dipole pairs (Pulupa et al. 2017). We implement a new simplified heuristic algorithm to calculate the QTN electron density, described in Appendix 3.8. We find the density derived from this method agrees with the QTN electron density from Moncuquet et al. (2020) within 5 – 10%.

Overview of In situ Observations

Figure 3.1 presents an overview of the encounter 13 ICME event between 05 September and 07 September 2022. The first panel shows three different plasma densities n computed from QTN (black), SPAN-E core electron fits (red) and SPAN-I proton moments (blue). The second panel displays the bulk SW speed $|\mathbf{v}|$ from SPAN-I moments for protons (blue) and alpha particles (green). Core electron (red), proton (blue), and alpha particle (green) temperatures are shown in the third panel. We compute an expected proton temperature (magenta) to compare with T_p as measured by SPAN-I (see Appendix 3.9 for details). The next panel shows the alpha particle to proton number density ratio n_a/n_p , with a dashed horizontal line indicating the typical fast solar wind value of 4%. Figures 3.1e and 3.1f display the Alfvén Mach Number ($M_A = v/(B/\sqrt{4\pi nm})$), and electron (red) and proton (blue) plasma β ($\beta_s = 8\pi n_s kT_s/B^2$ for species s), where the horizontal dashed lines are set to 1. The QTN electron density measurement is used for both β_e and β_p calculations, as the absolute accuracy of this method is superior and quasi-neutrality of the plasma may be safely assumed. While n_p is replaced by the QTN electron density for β_p (which includes a small alpha particle density portion), the proton temperature is used to compute β_p . The next panel shows the differential energy flux (Eflux) spectrogram, with units of $\text{eV s}^{-1} \text{cm}^{-2} \text{ster eV}^{-1}$, for protons summed over all Field-of-View (FOV) angles measured by SPAN-I. Figure 3.1h depicts the magnetic field data in RTN coordinates from FIELDS. The bottom panel shows ePADs measured at 433 eV from SPAN-E, normalized to the differential energy flux average for each sample. This energy band is selected as the peak of the strahl distribution tends to be within $\sim 300\text{--}500$ eV. The vertical dashed lines represent different time intervals when the plasma parameters display significant changes or interesting features. These intervals are summarized in Table 3.1, with a lettered label for each subinterval. Remote sensing observations from PSP and other spacecraft, described in Section 3.4, indicate the CME eruption became visible near time t_0 (05 September 2022 at 16:10 UT).

The SPAN-I sensor is mounted behind the thermal protection shield of PSP and therefore does not have a direct view of the Sun line. The gap in the SPAN-I FOV is mostly apparent at far heliocentric distances, where the spacecraft velocity is slow (relative to perihelion) and the solar wind appears more radial in the spacecraft frame. As the spacecraft approaches perihelion, the aberration between the SW velocity and spacecraft velocity becomes large enough to measure the full three dimensional velocity distribution function (VDF). The transition between regimes is variable and depends on local solar wind conditions. The light and dark orange shaded regions in Figure 3.1 depict intervals when either ions or alpha particles are outside of the SPAN-I FOV or energy range. During this event, the proton VDF moved out of the SPAN-I FOV and briefly out of the instrument energy range above 20 keV.

Interval	Start Time (UTC)	In situ Signatures	ICME Feature
t_0	22-09-05/16:10:00	Remote Sensing Observations	CME Eruption Visible
t_1	22-09-05/17:27:20	n, v, T , & B increase, 90° ePADs	ICME Shock
t_{2a}	22-09-05/17:33:04	Diffuse ePADs, Smooth $\mathbf{B}$	End of Sheath
t_{2b}	22-09-05/18:02:26	$T_p < 0.5 T_{exp}$	ICME Flank
t_{3a}	22-09-05/19:08:46	Bidirectional ePADs, $\mathbf{B}$ rotations	Closed Field Lines
t_{3b}	22-09-05/19:49:23	$T_p < 0.5 T_{exp}$, $n < 10^3 \text{ cm}^{-3}$	
t_{4a}	22-09-05/20:09:05	$n \gg 10^3 \text{ cm}^{-3}$, $\mathbf{B}$ rotation	First Inversion Line
t_{4b}	22-09-05/21:19:06	$T_p < 0.5 T_{exp}$, ePAD inversion	
t_{5a}	22-09-06/01:14:17	v & T decrease	Exiting ICME Plasma
t_{5b}	22-09-06/03:26:00	Stable plasma conditions	Typical SW
t_6	22-09-06/08:39:49	$M_A < 1$, Diffuse ePADs	Sub-Alfvénic Region
t_{7a}	22-09-06/17:27:28	$M_A > 1$, ePAD inversion	CS Leading Edge
t_{7b}	22-09-06/17:40:00	$M_A \ll 1$	CS Trailing Edge
t_8	22-09-07/06:11:28	M_A increase	
t_9	22-09-07/11:30:16	Pre-event SW conditions	End of Event

Table 3.1: Selected time intervals (UTC) during the encounter 13 ICME event. Some intervals are divided into subintervals, denoted by a lettered subscript. The plasma signatures described in the table represent the most significant in situ observations for each interval and correspond to an ICME feature.

ICME Shock Analysis

Figure 3.2 shows a magnified time series from time intervals t_1 (17:27:20 UTC) to t_4 (20:09:05 UTC) with similar panels to Figure 3.1. In addition, Figure 3.2 includes normalized ePADs measured at 1152 eV, and the magnetic field elevation angle θ_B and azimuthal angle ϕ_B . The magnetic field angles have been smoothed over an averaging window of 5.5 minutes to focus on large-scale rotations in $\mathbf{B}$. At time t_1 on 05 September 2022, the electron number density n_e is measured to abruptly increase from about 1880 to 4000 cm^{-3} , followed by a gradual decrease below 1000 cm^{-3} . The bulk speed $|\mathbf{v}|$ sharply increases from ~ 340 to 650 km s^{-1} , and continues to grow until 1350 km s^{-1} . The proton temperature T_p surges from 65 to 520 eV with a subsequent increase to 2 keV. The magnetic field magnitude B abruptly strengthens from approximately 625 to 1500 nT, and continues to about 1650 nT. These values of n , v , T , and B downstream of the discontinuity are not typical at these solar distances based on previous PSP encounters. This may indicate that PSP has encountered a compression region. In addition, PSP observes the magnetic field to become more turbulent after t_1 , along with an increase in higher energy protons ($E > 3500 \text{ eV}$, see Figure 3.1g) and electrons ($E > 350 \text{ eV}$, not shown).

To estimate the orientation of the leading discontinuity at time t_1 , we perform a shock geometry analysis using measurements from PSP. We implement a standard single-spacecraft analysis of the Magnetohydrodynamic (MHD) Rankine-Hugoniot (RH) jump conditions

across the discontinuity to derive the most probable shock plane propagation vector using the methods described in Viñas et al. (1986), Szabo (1994), and Koval et al. (2008).

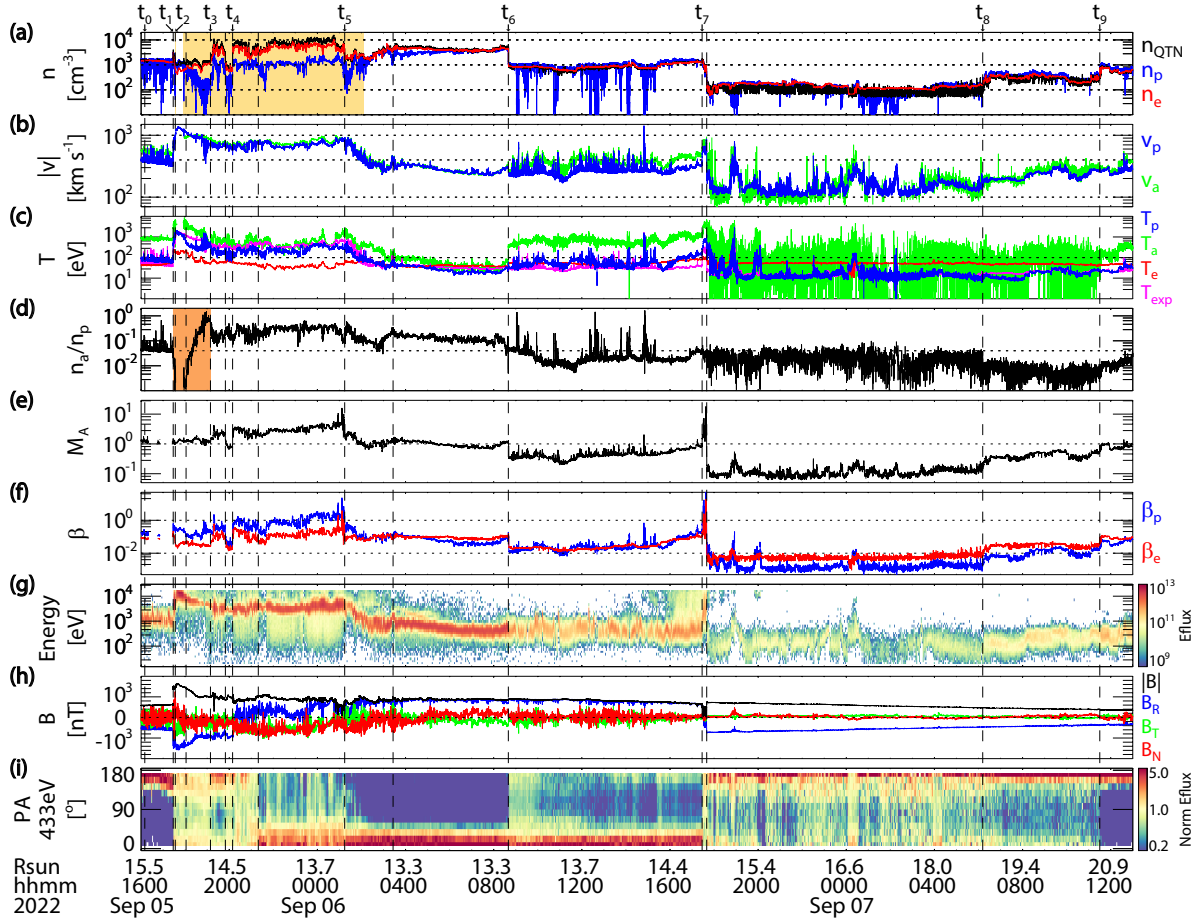


Figure 3.1: Time series of the encounter 13 ICME event observed by PSP. **(a)** QTN electron (black), proton (blue), and core electron (red) density, **(b)** proton (blue) and alpha particle (green) bulk speed, **(c)** proton (blue), alpha particle (green), core electron (red), and expected proton (magenta) temperature, **(d)** alpha particle to proton number density ratio, **(e)** Alfvén Mach number, **(f)** proton (blue) and electron (red) plasma β , **(g)** proton differential energy flux spectrum with units of $\text{eV s}^{-1} \text{cm}^{-2} \text{ster eV}^{-1}$, **(h)** RTN magnetic field with magnitude (black), **(i)** normalized electron pitch angle distribution at 433 eV. The heliocentric distance of PSP in R_S is shown at the bottom along the time axis. The light and dark orange shaded regions depict intervals when ions or alpha particles are outside of the SPAN-I FOV or energy range, respectively. The vertical dashed lines represent different time intervals when the plasma parameters display significant changes or interesting features, corresponding to the intervals in Table 3.1. The subinterval labels are not shown.

The sampled upstream and downstream in situ measurements from PSP represent the slab-symmetric asymptotic conditions. Using the pre-shock and post-shock conditions, the shock orientation and speed are determined based on the optimal solution that satisfies Maxwell's equations and the conservation of mass, momentum, and energy across the discontinuity. The RH equations will yield a unique and robust solution if the bootstrap analysis consis-

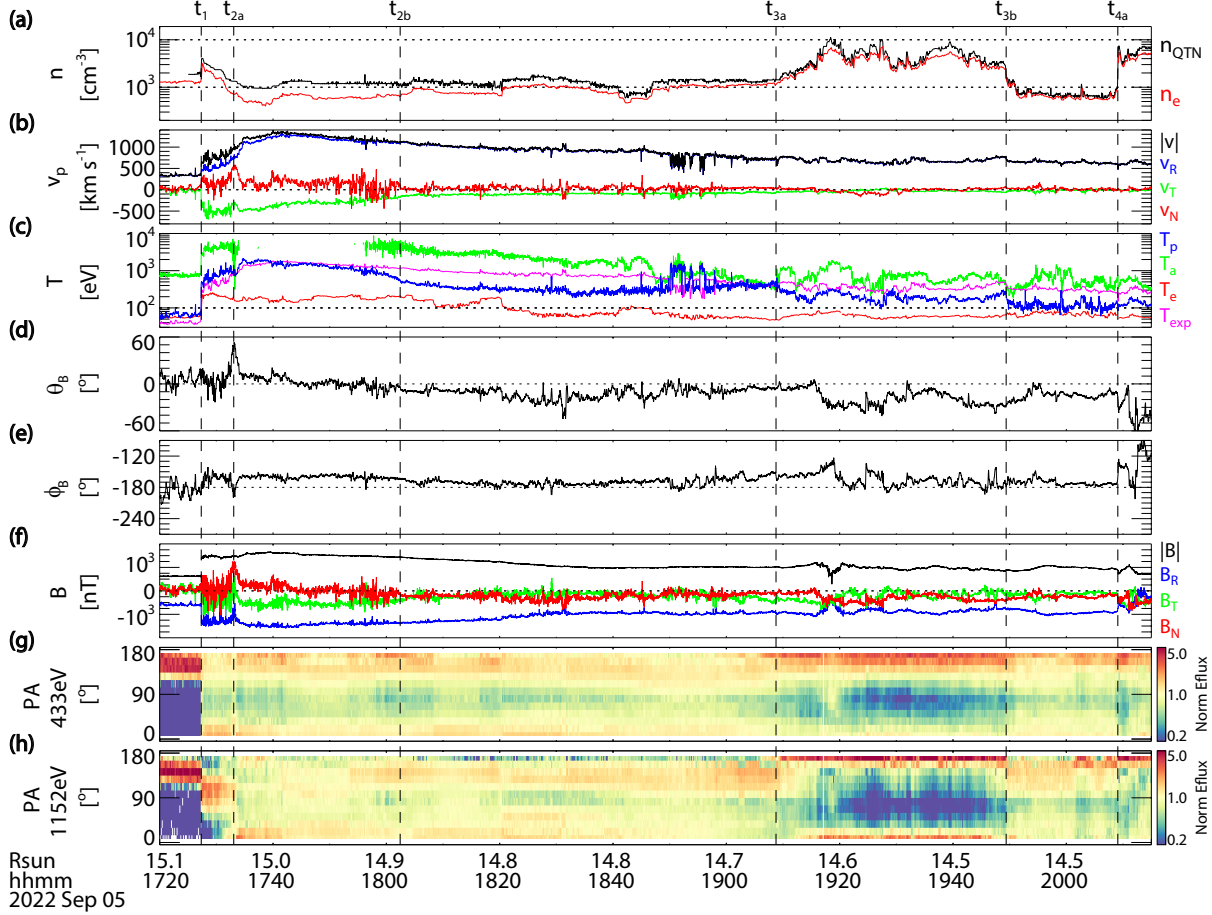


Figure 3.2: Time series from 17:20 to 20:15 UTC on 05 September 2022 as observed by PSP. (a) QTN (black) and core (red) electron density, (b) RTN proton velocity with magnitude (black), (c) proton (blue), alpha particle (green), core electron (red), and expected proton (magenta) temperature, (d) magnetic field elevation angle, (e) magnetic field azimuthal angle, (f) RTN magnetic field with magnitude (black), and normalized ePADs measured at (g) 433 eV and (h) 1152 eV. The magnetic field angles have been smoothed using a moving average of 5.5 minutes. The heliocentric distance of PSP in R_S is shown at the bottom along the time axis. The vertical dashed lines represent different time intervals, corresponding to the intervals in Table 3.1, from t_1 to t_{4a} .

tently approaches the same solution. The model assumes ideal MHD, thermal isotropy, and equal partitioning of pressures between ions and electrons. As none of these are strictly true, we expect only a rough estimate for the shock solution from this exercise.

The upstream and downstream plasma parameters near time t_1 are described in Table 3.2, averaged ~ 60 and 30 seconds before and after the discontinuity, respectively. These include the RTN velocity $\mathbf{v}_{RTN}$, RTN magnetic field $\mathbf{B}_{RTN}$, proton scalar thermal speed ω_p , and electron number density n_e . The associated uncertainties are based on the ensemble of outcomes from sampling the upstream and downstream conditions. We find a fast-forward shock solution of cross-boundary flow with a Fast Mode Mach number transition from 2.4 ± 0.1 on the upstream side to 0.65 ± 0.03 on the downstream side. The most probable shock orientation is found to be significantly non-radial, with a propagation vector directed $-51 \pm 4^\circ$ from +R in the RT-plane and $23 \pm 3^\circ$ from +R in the RN-plane. The most probable propagation speed is estimated as $1002 \pm 30 \text{ km s}^{-1}$ in the SW frame. Figure 3.3 illustrates this inferred shock geometry with the orbital trajectory of PSP for encounter 13 in Heliocentric Earth Equatorial (HEEQ) coordinates. The portion of the orbit from times t_1 to t_9 listed in Table 3.1 is highlighted in orange. The HEEQ coordinate system is defined with the x axis directed toward the intersection between solar equator and solar central meridian of date and the z axis toward the solar rotation axis (Fränz et al. 2002). PSP is located on its orbit at time t_1 , with the projected shock plane and normal in red. The black dashed line in Figure 3.3 represents the eruption axis of the event, aligned with active region AR 13102 (previously named AR 13088), discussed in Section 3.5. In addition, Figure 3.3 includes the orbital positions of Earth, Solar Orbiter (SOLO; Müller, D. et al. 2020), and the Solar Terrestrial Relations Observatory (STEREO; Kaiser et al. 2008), specifically STEREO-A (STA). This shock geometry is consistent with encountering the flank of an ICME as it expands into interplanetary space (see Figure 3.9a below).

Plasma Parameter	Upstream	Downstream
$v_R \text{ (km s}^{-1}\text{)}$	333 ± 10	455 ± 16
$v_T \text{ (km s}^{-1}\text{)}$	8 ± 17	-358 ± 46
$v_N \text{ (km s}^{-1}\text{)}$	60 ± 32	251 ± 41
$B_R \text{ (nT)}$	-602 ± 20	-1349 ± 43
$B_T \text{ (nT)}$	86 ± 59	142 ± 196
$B_N \text{ (nT)}$	102 ± 75	317 ± 222
$\omega_p \text{ (km s}^{-1}\text{)}$	113 ± 4	294 ± 13
$n_e \text{ (cm}^{-3}\text{)}$	1962 ± 74	3989 ± 60

Table 3.2: Upstream and downstream conditions from the leading discontinuity at time t_1 for the RTN velocity $\mathbf{v}_{RTN}$, RTN magnetic field $\mathbf{B}_{RTN}$, proton scalar thermal speed ω_p , and electron number density n_e . The uncertainties are based on the ensemble of outcomes from sampling the upstream and downstream conditions.

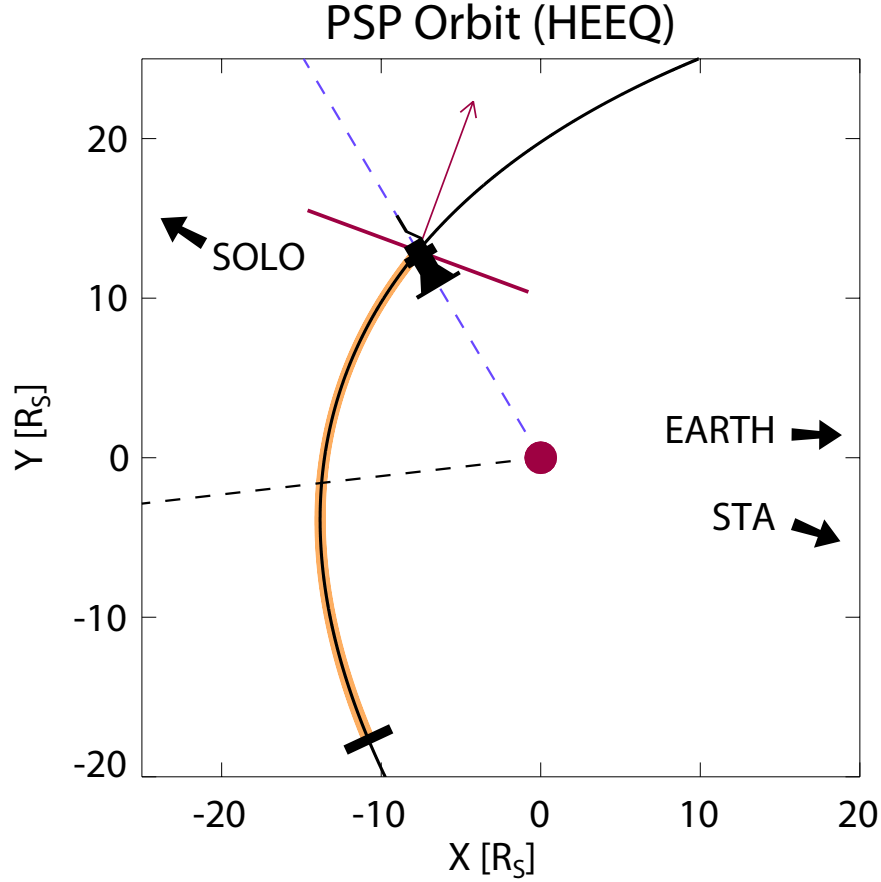


Figure 3.3: Polar view of the encounter 13 PSP orbit in HEEQ coordinates, where the trajectory follows a counterclockwise direction around the Sun (red) with time. The position of PSP at time t_1 is shown with the spacecraft directed toward the Sun, where the purple dashed line represents the Sun-PSP line. The portion of the orbit highlighted in orange signifies the ICME event duration from time t_1 to t_9 . The eruption axis of the event is shown by the black dashed line, aligned with active region AR 13102. The shock plane and normal are shown with a red line segment and arrow, respectively. In addition, the positions of Solar Orbiter (SOLO), STEREO-A (STA) and Earth are indicated by black arrows.

ICME Sheath and Flank

Before encountering the region influenced by the ICME at time t_1 , the electron strahl is directed 180° from the magnetic field direction (Figures 3.2g-h). We can examine the evolution of the ePADs to study the magnetic topology of the event. After the discontinuity at t_1 , the electron distributions evolve differently depending on energy, shown in Figure 3.4 from 17:25 to 17:45 UTC. Each panel displays a different energy between ~ 200 eV and 2 keV,

spanning most of the instrumental range. The ePADs with energies exceeding 750 eV exhibit a strong peak near 90° , most likely associated with the ICME shock. As the energy decreases, the peak in the distribution slowly shifts towards higher pitch angles. Eventually, the ePADs appear diffuse with no strong indication of bidirectional electrons. However, ePADs below 300 eV retain a strong, distinct strahl flowing anti-parallel to the magnetic field. The duration of the 90° ePADs is dependent on energy, but primarily concludes around time t_{2a} (17:33:04 UTC).

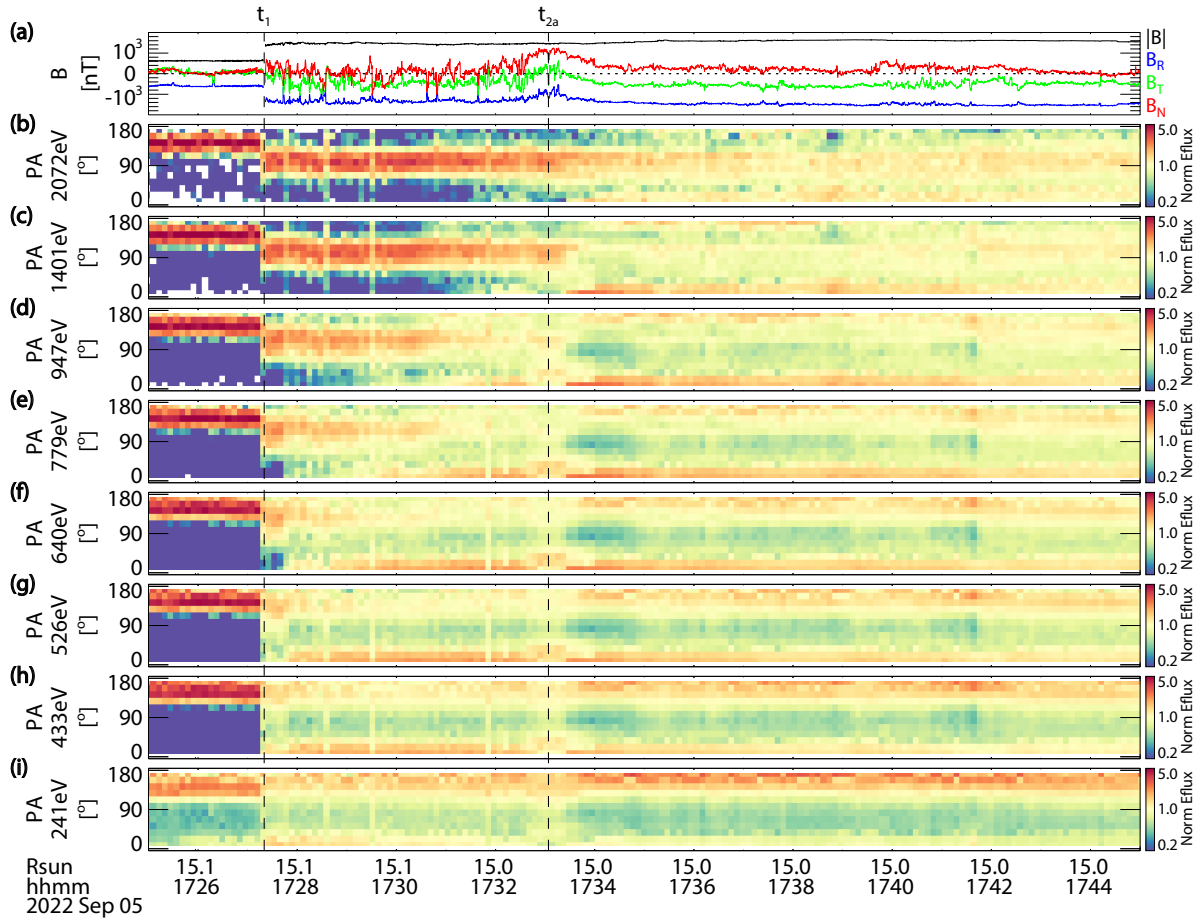


Figure 3.4: Time series from 17:25 to 17:45 UTC on 05 September 2022, showing the (a) RTN magnetic field vector with magnitude (black). The normalized ePADs are shown for energies (b) 2072 eV, (c) 1401 eV, (d) 947 eV, (e) 779 eV, (f) 640 eV, (g) 526 eV, (h) 433 eV, and (i) 241 eV. The time intervals t_1 and t_{2a} are designated by the vertical black dashed lines. The ePADs above 750 eV exhibit a strong peak near 90° after time t_1 until roughly time t_{2a} , while electrons below 300 eV show a strong, distinct strahl flowing anti-parallel to the magnetic field.

While the strahl measured below 300 eV by SPAN-E are predominantly unidirectional following time t_{2a} , the ePADs above 300 eV remain diffuse. A small rotation in θ_B of about 50° is centered around this t_{2a} from 17:32:33 to 17:34:37 UTC (Figure 3.2d). After this time, the magnitude of the fluctuations in $\mathbf{B}$ diminishes, suggesting the spacecraft exits the ICME sheath region. Furthermore, v , T_p , and B reach maximum values a few minutes after t_{2a} of approximately 1375 km s^{-1} , 2300 eV , and 1640 nT , before gradually decreasing in Figure 3.2. Typical evidence of an MC is identified at time t_{2b} (18:02:26 UTC), when $T_p < 0.5T_{exp}$ and the fluctuations in $\mathbf{B}$ diminish further. At time t_{3a} (19:08:46 UTC), SPAN-E observes dominantly bidirectional electrons for energies above 300 eV, accompanied by a period of high density ($\sim 10^4 \text{ cm}^{-3}$). This interval of bidirectionality is one of the clearest examples ever observed by PSP during an encounter. This interval is strongly suggestive of closed magnetic field lines driven out into interplanetary space by the ICME. Subsequent to time t_{3a} , the alpha particles return to the energy range measured by SPAN-I with $n_a/n_p > 8\%$ (Figure 3.1d). This interval of high n_a/n_p is not typical of solar wind plasma, and may indicate the ICME flank. Furthermore, PSP detects two rotations in θ_B of about 40° starting at time 19:15:30 UTC, with a period of ~ 20 minutes. FIELDS also observes a small rotation in ϕ_B of more than 30° over 5 minutes, centered around 19:18:30 UTC. Following time t_{3b} (19:49:23 UTC), the plasma density drops below 1000 cm^{-3} and $T_p < 0.5T_{exp}$. This interval is also marked by diffuse ePADs, suggesting PSP no longer observes closed magnetic field lines.

The low density period continues until time t_{4a} (20:09:05 UTC), when n abruptly increases to $\sim 7300 \text{ cm}^{-3}$, shown in Figure 3.5. The radial component of $\mathbf{B}$ intermittently flips polarity between times t_{4a} and t_{4b} (21:19:06 UTC), along with the dominant direction of the strahl. This indicates the first magnetic inversion line that PSP travels through during the event. Within this interval, the values of M_A and β_p simultaneously increase to ~ 3 and ~ 1 , respectively. In addition, PSP measures a sharp rotation in θ_B of more than 50° . The rotation in ϕ_B is instead more gradual over time, rotating 180° until time t_{5a} (22-09-06/01:14:17 UTC). The sign of B_R is consistently positive after time t_{4b} , when the strahl is predominantly flowing parallel to the magnetic field direction, suggesting PSP is no longer grazing the inversion line. Moreover, $T_p < 0.5T_{exp}$ is observed for a brief interval of ~ 22 minutes after t_{4b} . At time t_{5a} , the strahl distribution becomes more consistent and unidirectional, indicating open field lines. PSP observes v and T_p to decline gradually from ~ 850 to 330 km s^{-1} and from ~ 400 to 40 eV , respectively. This change is paired with M_A and plasma β decreasing to ~ 1 and ~ 0.1 , respectively. This shift towards typical solar wind conditions may indicate that PSP is no longer traveling through ICME plasma. As PSP exits this region, the spacecraft observes a θ_B rotation from -75° to 0° over more than one hour. After time t_{5b} (03:26:00 UTC), the observed plasma parameters presented in Figure 3.5 become less variable.

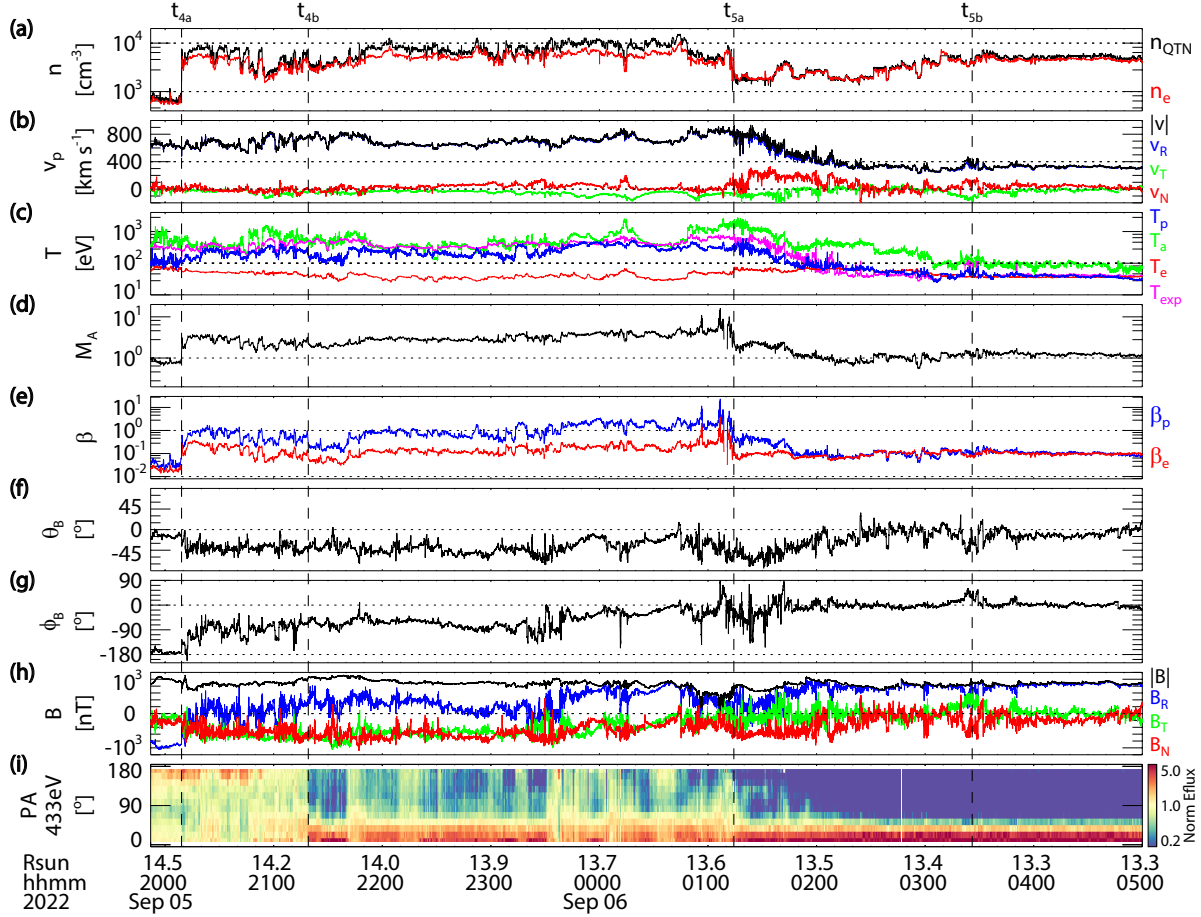


Figure 3.5: Time series from 22-09-05/19:52 to 22-09-06/05:00 UTC as observed by PSP. (a) QTN (black) and core (red) electron density, (b) RTN proton velocity with magnitude (black), (c) proton (blue), alpha particle (green), core electron (red), and expected proton (magenta) temperature, (d) Alfvén Mach number, (e) proton (blue) and electron (red) plasma β , (f) magnetic field elevation angle, (g) magnetic field azimuthal angle, (h) RTN magnetic field with magnitude (black), and (i) normalized ePADs at 433 eV. The magnetic field angles have been smoothed using a moving average of 5.5 minutes. The heliocentric distance of PSP in R_S is shown at the bottom along the time axis. The vertical dashed lines represent different time intervals, corresponding to the intervals in Table 3.1.

ICME Wake

After time t_6 (08:39:49 UTC), the ePADs diffuse across the range of pitch angles while maintaining a preference toward 0° , illustrated in Figure 3.6. In addition, the values of n_e , M_A , and plasma β suddenly decrease below 1000 cm^{-3} , 0.02 and 0.4, respectively. The alpha

particle temperature T_a abruptly increases above 400 eV, while n_a/n_p drops below 4%. This interval is characterized by a sub-Alfvénic region behind the ICME unique to the first half of the event.

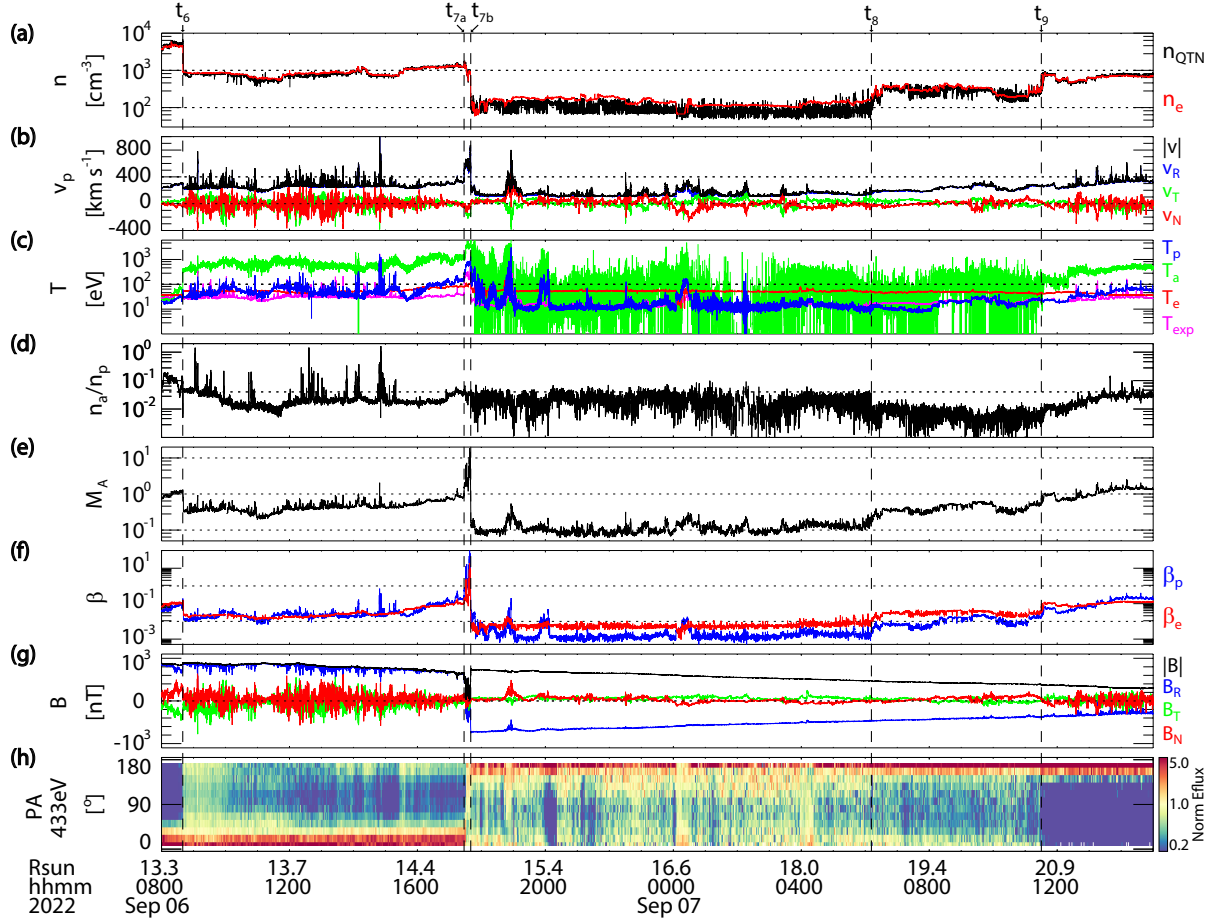


Figure 3.6: Time series from 22-09-06/08:00 to 22-09-07/15:00 UTC as observed by PSP. (a) QTN (black) and core (red) electron density, (b) RTN proton velocity with magnitude (black), (c) proton (blue), alpha particle (green), core electron (red), and expected proton (magenta) temperature, (d) alpha particle to proton number density ratio, (e) Alfvén Mach number, (f) proton (blue) and electron (red) plasma β , (g) RTN magnetic field with magnitude (black), and (h) normalized ePADs at 433 eV. The heliocentric distance of PSP in R_S is shown at the bottom along the time axis. The vertical dashed lines represent different time intervals, corresponding to the intervals in Table 3.1, from t_6 to t_9 .

Between times t_{7a} (17:27:28 UTC) and t_{7b} (17:40:00 UTC) shown in Figure 3.7, both M_A and β_p become greater than 1, while v and T_p increase above 600 km s⁻¹ and 700 eV, respectively. At the beginning of this interval, ϕ_B rotates from 0° to 180° in under 2 minutes,

while θ_B fluctuates between -40° and 40° (not shown). The spacecraft passes through a second magnetic inversion line during this period as evidenced by the strahl polarity changing from 0° to 180° . The plasma conditions before t_{7a} and after t_{7b} are highly asymmetric, most likely attributed to different sides of a current sheet (CS). A reconnection exhaust is observed during this crossing, characterized by proton flow speed (Figure 3.7b) and temperature (Figure 3.7c) enhancements, and a significant drop in magnetic intensity (Figure 3.7g).

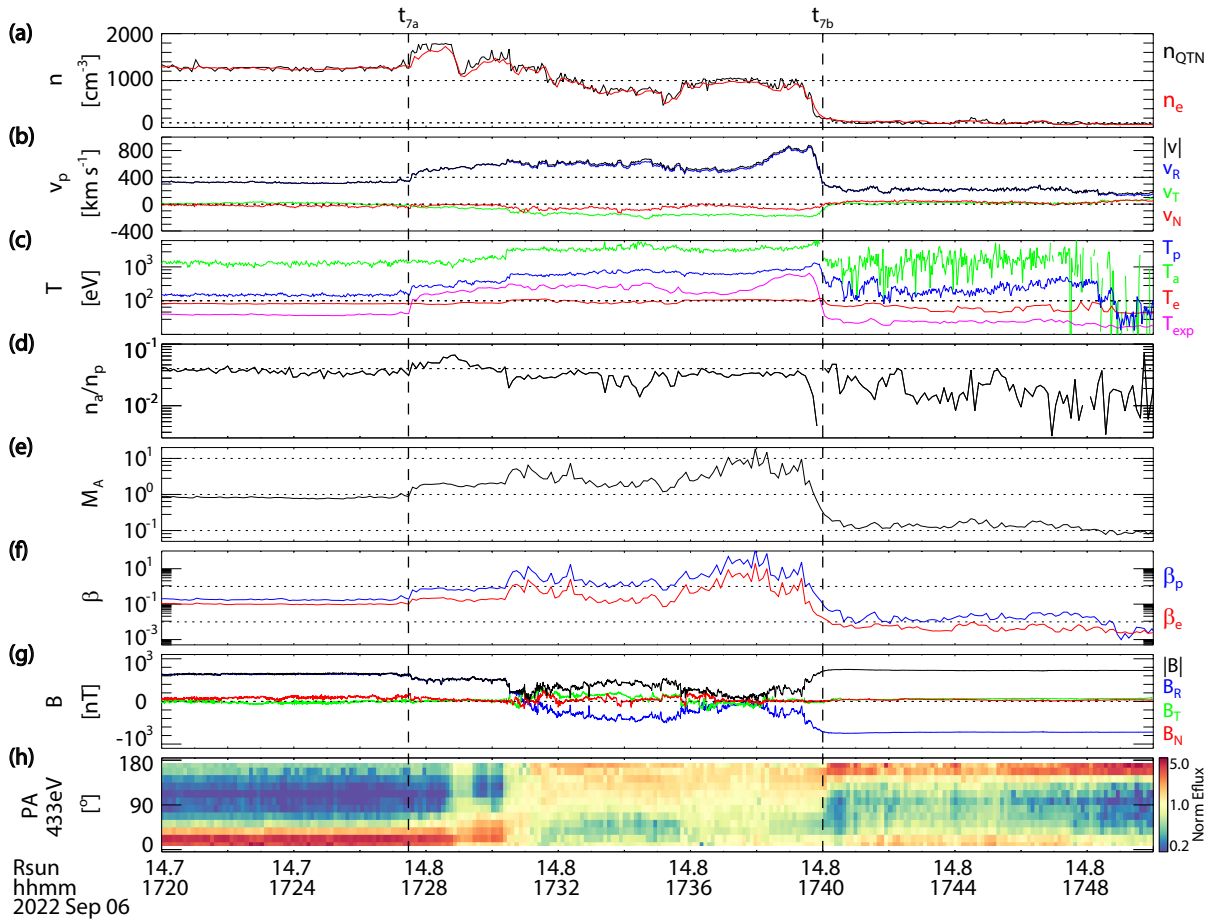


Figure 3.7: Time series from 22-09-06/17:20 to 22-09-06/17:50 UTC as observed by PSP during the second magnetic inversion line crossing. **(a)** QTN (black) and core (red) electron density, **(b)** RTN proton velocity with magnitude (black), **(c)** proton (blue), alpha particle (green), core electron (red), and expected proton (magenta) temperature, **(d)** alpha particle to proton number density ratio, **(e)** Alfvén Mach number, **(f)** proton (blue) and electron (red) plasma β , **(g)** RTN magnetic field with magnitude (black), and **(h)** normalized ePADs at 433 eV. The heliocentric distance of PSP in R_S is shown at the bottom along the time axis. The vertical dashed lines represent t_{7a} and t_{7b} , corresponding to Table 3.1.

After time t_{7b} , n and v decrease below 100 cm^{-3} and 300 km s^{-1} , respectively. This interval is followed by a dominantly sub-Alfvénic region of space with $M_A \sim 0.1$. In addition, SPAN-I observes a considerable alpha particle depletion that does not appear to be an instrumental artifact (not shown). This low density region may be a result of PSP traveling through the rear region of the ICME. After 17:48:30 UTC, v and T_p decrease below 150 km s^{-1} and 100 eV respectively. These conditions persist for more than 12 hours until time t_8 (07 September 2022 at 06:11:28 UTC), when n and M_A begin to rise, while n_a/n_p slightly decreases to 0.01. The increase in n_e may indicate that ambient solar wind is flowing through the region behind the ICME. Eventually, at time t_9 (11:30:16 UTC), PSP exits the sub-Alfvénic region and observes solar wind parameters similar to pre-event conditions, indicating the end of the event.

3.4 CME Reconstruction From Remote Sensing

Besides in situ instrumentation, PSP includes the Wide-field Imager for Solar PRobe Plus (WISPR; Vourlidas et al. 2016). WISPR is situated on the ram side of PSP and is comprised of two telescopes with a combined 95° radial by 58° transverse FOV and spatial resolution of $6.4'$. The range of the inner (WISPR-I) and outer (WISPR-O) telescopes span from 13.5° to 53° and 50.5° to 108.5° , respectively. A snapshot of the combined WISPR FOV on 05 September 2022 at 17:30 UT is shown in Figure 3.8. This image has been processed via the L3 Algorithm developed by the WISPR team and detailed in the Appendix of Liewer et al. (2023). High particle density is associated with brighter regions in the white-light image. WISPR, along with other remote-sensing instruments discussed below, observes the beginning of the CME eruption on 05 September 2022 at $\sim 16:10$ UT. At time t_1 when PSP detects a fast-forward shock, WISPR images a clear, high density spherical front. The leading edge observed by WISPR continues to propagate beyond PSP toward larger heliocentric longitudinal values after t_1 .

Additional imagers on board other spacecraft observed the ICME event, such as the Sun Earth Connection Coronal and Heliospheric Investigation (SECCHI; Howard et al. 2008) instrument suite on STEREO-A (STA). SECCHI comprises five telescopes that cover a wide FOV from the Sun to interplanetary space. One of these telescopes, called COR2, is a visible light Lyot coronagraph with a FOV extent of 2 to $15 R_S$ from the Sun. Figure 3.8a displays a running difference image from the COR2 telescope on 05 September 2022 at 17:26 UT. STEREO-A observes a wide angled CME halo event during this time. As the expelled material expands with time, it appears to envelop the Sun. This suggests that at least a portion of the plasma propagates along the Sun-observer line. Both the COR2 and WISPR snapshots in Figure 3.8 represent the images observed nearest to time t_1 , when PSP encounters an ICME shock. During this time, COR2 observes a spherical front, outlined by a red dashed curve in Figure 3.8a, which is considered a shock candidate. While this front appears to intercept PSP from the STEREO-A images, we note that the coronagraph observations represent line-of-sight integrated emission, and are not sufficient to determine

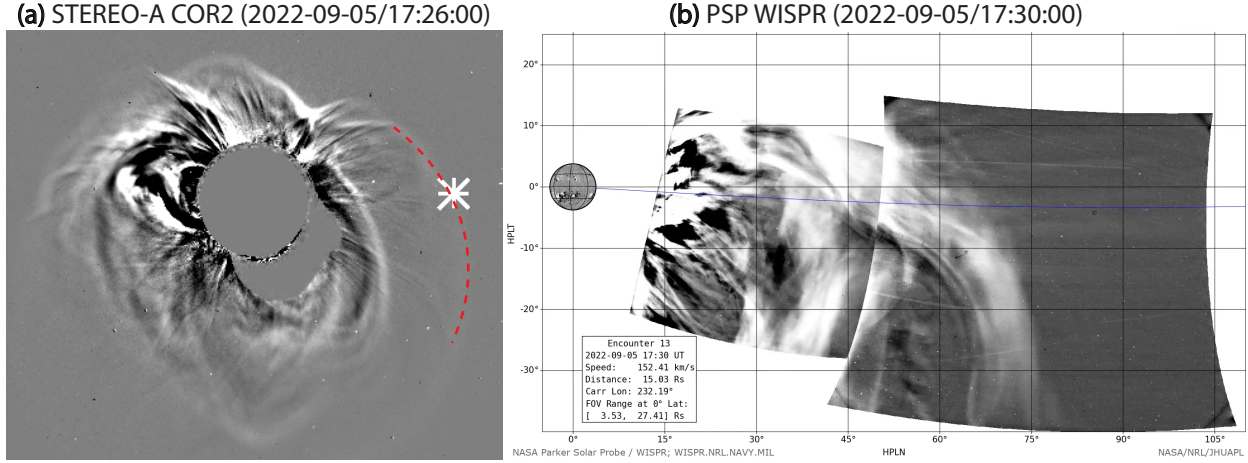


Figure 3.8: Remote sensing observations of the CME on 05 September 2022 near time t_1 when PSP encountered a fast-forward shock. **(a)** Running difference COR2 image from STEREO-A at 17:26 UT, where the location of PSP is represented by the white asterisk. An expanding front observed by COR2 is outlined by a red dashed curve. **(b)** The combined FOV of WISPR-I (left) and WISPR-O (right) at 17:30 UT, processed using the L3 algorithm and overlaid on a grid of heliocentric longitude and latitude. The Sun is shown to scale at the origin, with an overlaid SDO/HMI Carrington map. The orbital plane of PSP is represented by the blue line. The corresponding spacecraft speed, distance from the Sun, Carrington longitude, and FOV range at 0° latitude for this time are reported near the bottom left.

if this front reaches PSP. Therefore, it is possible that PSP could be behind or in front of the CME shock.

To better evaluate the relative position between the ICME and PSP, we reconstruct the CME using forward-modeling, as in Braga et al. (2023, *in preparation*). As CMEs are optically thin, coronagraphs only measure the integrated brightness over many layers along the line-of-sight. While these observations provide a limited perspective, we can implement CME reconstruction to improve our overall understanding. The forward modeling from Braga et al. (2023, *in preparation*) implements a graduated cylindrical shell (GCS) model (Braga et al. 2022; Thernisien 2011; Thernisien et al. 2006) of a flux rope structure. This empirical model is created from two conical legs connected to a curved tubular front with a cross section that increases with height from the solar surface. This shape is sometimes referred to as a “hollow croissant.” The model is governed by three positional parameters (Stonyhurst latitude θ , Stonyhurst longitude ϕ , and tilt angle of the source region neutral line γ), two angular widths (angular width $2\alpha_w$ between the legs and half-angle δ of the conical leg), and the CME height h . The half-angle δ is sometimes rewritten in terms of the aspect ratio $\kappa = \sin(\delta)$, which sets the expansion rate versus the CME height. In addition, we assume that the CME expands self-similarly, i.e., all model parameters are kept constant,

except the height h . These free parameters are determined based on observations from the SECCHI and WISPR instruments, along with the Large Angle Spectroscopic Coronagraph (LASCO; Brueckner et al. 1995) on board the Solar and Heliospheric Observatory (SOHO; Domingo et al. 1995).

To reach agreement with coronagraph observations, this event is modeled by two separate CMEs, denoted as the primary and secondary portions of the CME event. Figure 3.9 shows the primary CME (CME1) and secondary CME (CME2) at various stages of their expansion. CME1 erupted on the backside of the Sun with respect to Earth, propagating away from the planet. The model parameters for CME1 are $\theta = 168^\circ$, $\phi = -20^\circ$, $\gamma = 66^\circ$, $\alpha_w = 45^\circ$, and $\delta = 43^\circ$. Figure 3.9a-b shows CME1 (blue) at the time that PSP observes a fast-forward shock (t_1), with a height of $h = 14 R_S$. The CME is projected in HEEQ coordinates for both the polar and equatorial planes. The dark blue outline represents the cross section of the structure with the PSP orbital plane. The spacecraft is predicted to enter the flank of CME1 for a brief period of time from $\sim 19:40$ UT to $20:20$ UT, skimming the outer edge of the tubular front depicted in Figure 3.9c-d. About $02:30$ UT on 06 September, PSP travels through a leg of CME1, where the height has expanded to $h = 104 R_S$, shown in Figure 3.9e-f. The parameters for CME2 are $\theta = 241^\circ$, $\phi = 5^\circ$, $\gamma = 0^\circ$, $\alpha_w = 20^\circ$, and $\delta = 23^\circ$. This CME is modeled to propagate quasi-perpendicular to the Sun-Earth line. The spacecraft enters CME2 on 06 September around $16:40$ UT, when $h = 202 R_S$, shown in Figure 3.9g-h. This time closely corresponds to t_7 , when PSP crosses the second magnetic inversion line.

It is important to note that the CME model will only reconstruct the large-scale morphology, and not necessarily small-scale structure observed by PSP. In addition, self-similar propagation may not hold true for the entire duration of this event. This event is only observable to COR2 and LASCO C3 between $\sim 16:30$ and $\sim 18:30$ UT. After this period, the CME reaches the outer edge of these coronagraphs and we do not have sufficient observational constraints to evaluate whether self-similarity is maintained. Moreover, the CME may experience some deflection, distortion, or rotation as it propagates from the Sun. Even with three observational points of view, we expect an error of CME propagation up to 1-2 hours (McComas et al. 2023). A more detailed description of the CME models and their agreement with coronagraph observations can be found in Braga et al. (2023, *in preparation*).

We now have a model that is constrained by remote sensing estimates and geometrical assumptions for the expansion rate of the CME. To compare the model with in situ data, Figure 3.10 shows a snapshot of the event at $19:10$ UT on 05 September 2022 near time t_{3a} . This time is denoted by closed field lines possibly related to the flank of the structure. Due to the limited cadence of the imaging observations, the closest COR2 image to this time is $19:11$ UT, illustrated in Figure 3.10a. Similarly, Figure 3.10b shows the WISPR image at $19:07$ UT, observed near time t_{3a} . Figure 3.10c portrays the PSP orbital position at $19:10$ UT, along with CME1. Finally, a subset of the in situ data is shown in Figure 3.10d, spanning from $15:10$ to $23:10$ UT. The black vertical line is centered at $19:10$ UT indicating the time of the current observation. An animation of this figure is available (Romeo et al. 2023a), spanning from 22-09-05/13:00 to 22-09-07/16:00 UT.

Based on the animation from Figure 3.10, COR2 and WISPR begin to observe a per-

turbation in the white-light images around 16:10 UT. According to the GCS model, CME1 erupts around 16:10 UT on 05 September 2022, while CME2 is modeled to erupt around 20:30 UT. The in situ instruments on board PSP observe CME perturbed plasma after t_1 , more than an hour after the eruption. By the end of the event at time t_9 , both in situ and remote sensing measurements (shown in the animation) return to pre-CME conditions.

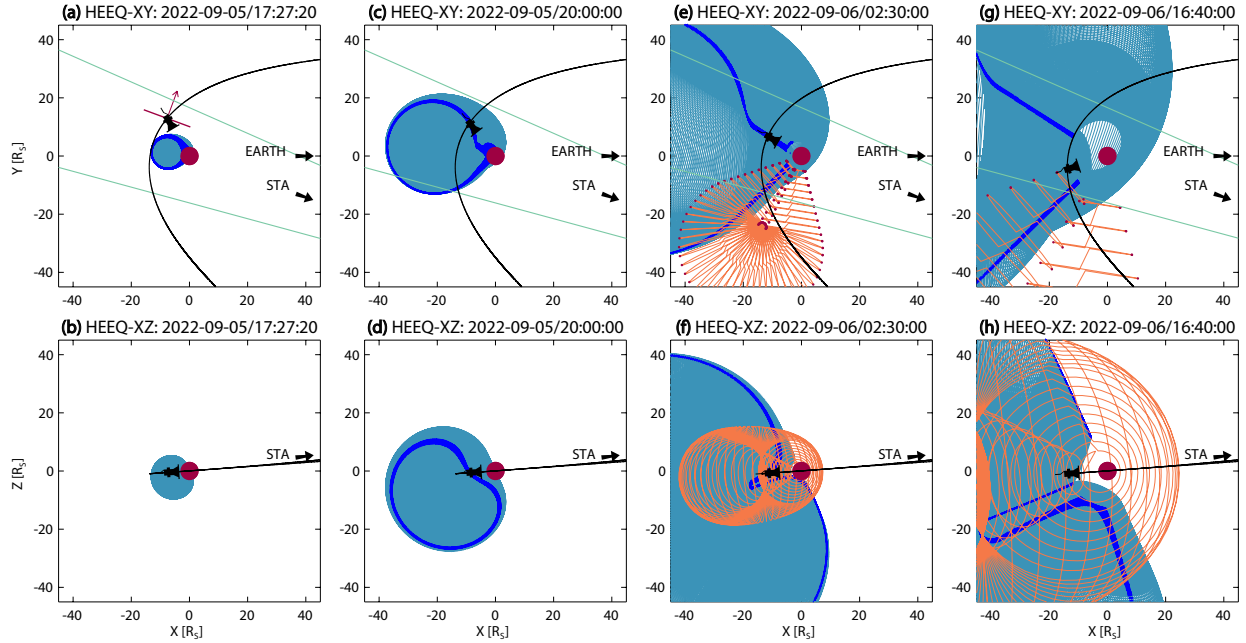


Figure 3.9: Encounter 13 PSP orbit in HEEQ coordinates, where the trajectory follows a counterclockwise motion around the Sun (red, not to scale) with time. The CME1 (blue region) model at time 22-09-05/17:27:20 UT, when PSP observes a fast-forward shock, is shown in the (a) XY plane (solar equatorial plane) and (b) XZ plane. The dark blue outline corresponds to the CME cross section within the orbital plane of PSP. The shock plane and normal are represented by a red line segment and arrow, respectively. At time 22-09-05/20:00 UT when PSP briefly enters the flank, CME1 is illustrated in the (c) XY plane and (d) XZ plane. Both CME1 and CME2 (orange region) at time 22-09-06/02:30 UT, when PSP enters the leg of CME1, are shown in the (e) XY plane and (f) XZ plane. The dark red outline along the wire frame of CME2 corresponds to the cross sectional area within the orbital plane of PSP. At time 22-09-06/16:40 UT, when PSP passes the second magnetic inversion line, CME1 and CME2 are modeled in the (g) XY plane and (h) XZ plane. In addition, the orbital position of STEREO-A (STA) and Earth are depicted by black arrows, with the COR2 FOV extent portrayed by two green lines in the solar equatorial plane.

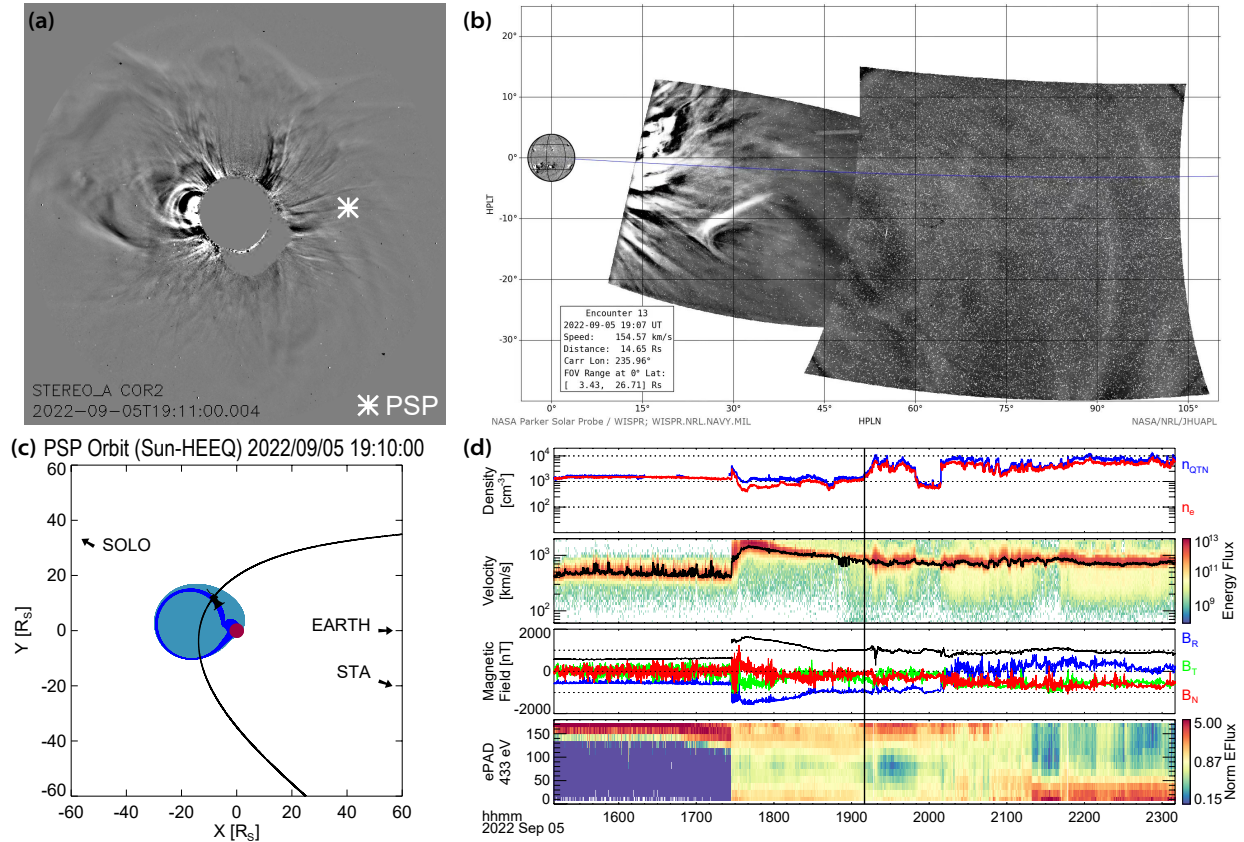


Figure 3.10: Combination of in situ data, remote sensing observations, and CME modeling for the ICME on 05 September 2022 near time t_{3a} , when PSP observes closed magnetic field lines based on bidirectional electrons. **(a)** Running difference COR2 image from STEREO-A at 19:11 UT, where the location of PSP is represented by the white asterisk. **(b)** The combined WISPR FOV at 19:07 UT, overlaid on a grid of heliocentric longitude and latitude, with an SDO/HMI Carrington map overlaid on a scale representation of the Sun. **(c)** PSP orbital trajectory in HEEQ coordinates at 19:10 UT, with CME1 in blue. The orbital location of Solar Orbiter (SOLO), STEREO-A (STA), and Earth are shown by black arrows. **(d)** In situ measurements observed by PSP. From top to bottom, the electron density from QTN (blue) and SPAN-E (red), solar wind velocity spectra from SPAN-I, RTN magnetic field vector with magnitude (black), and normalized ePADs measured at 433 eV. The time series is between 15:10 and 23:10 UT, with a vertical black bar indicating the time 19:10 UT. An animation of this figure is available in the online Journal, spanning from 22-09-05/16:00 to 22-09-07/12:00 UT, syncing the different cadence rates of the coronagraph imaging, in situ measurements, and CME modeling. A high resolution version of the animation is provided by Romeo et al. (2023a).

3.5 Large Scale Reorganization of the Coronal Field

The in situ data observed by PSP and SOLO can be used in combination with Potential Field Source Surface (PFSS) models (Altschuler et al. 1969; Schatten et al. 1969) to investigate the effect of the eruption on the large scale configuration of the solar corona. While the CME reconstruction detailed in Section 3.4 models the large-scale morphology of the density structure carried by the CME, it does not constrain the magnetic structure and how it interacts with the coronal field, nor inform about the solar origin of the event. We follow the framework described in Badman et al. (2020) to project in situ magnetic polarity measurements taken by PSP and SOLO (Horbury et al. 2020) down to spatial coordinates in the outer corona before and after the CME eruption. We compare these measurements with a PFSS extrapolated model of the HCS before and after the eruption, examining the global magnetic topology, and how it relates to the source active region of the CME as seen in photospheric magnetograms.

In this study, we implement Air Force Data Assimilative Flux Transport (ADAPT; Arge et al. 2010) magnetograms modeled from Global Oscillation Network Group (GONG; Harvey et al. 1996) observations. Referred to as ADAPT-GONG magnetograms, these synoptic maps provide the lower boundary condition of the magnetic field at the surface of the photosphere. The PFSS model assumes a magnetostatic corona with zero current between the photosphere and an equipotential source surface region with radius $R_{SS} = 2.5 R_S$. At the source surface, the magnetic field lines are assumed to be radial, with solar wind plasma flowing away from the Sun along open field lines. With these assumptions, the coronal field is represented as a scalar potential allowing the PFSS model to generate the three dimensional magnetic field structure within the modeling domain.

Figure 3.11 shows the ADAPT-GONG magnetograms in Carrington coordinates observed roughly one week before and after the CME event. The red and blue shading on the synoptic maps indicate positive and negative radial photospheric magnetic fields, respectively. The nominal HCS is calculated from the PFSS model and shown in each magnetogram as a black curve. To connect these results with PSP in situ measurements above $13 R_S$ from the Sun, the spacecraft orbit is ballistically propagated toward the source surface (Nolte et al. 1973). For any given point in space, the position of PSP can be connected back toward the Sun along its corresponding Parker spiral field line based on the measured solar wind speed. The orbital trajectories for PSP and SOLO are ballistically mapped down to the source surface in Figure 3.11, where the black and white squares indicate the initial and final positions of each orbit, respectively. The orbits are colorized according to the measured polarity of the magnetic field at each spacecraft. The orbits span about one month for each magnetogram, and include a few data gaps of the magnetic field. By including the magnetic field observed by SOLO, we can compare the polarity measured by PSP before and after the event.

In Figure 3.11a, the magnetogram was computed with observations during a full Carrington rotation until 01 September 2022, prior to the ICME event. The modeled HCS structure agrees well with the polarity flip measured by PSP and SOLO. Between 01 August and 02 September, both PSP and SOLO traverse the entire longitudinal range of the Sun.

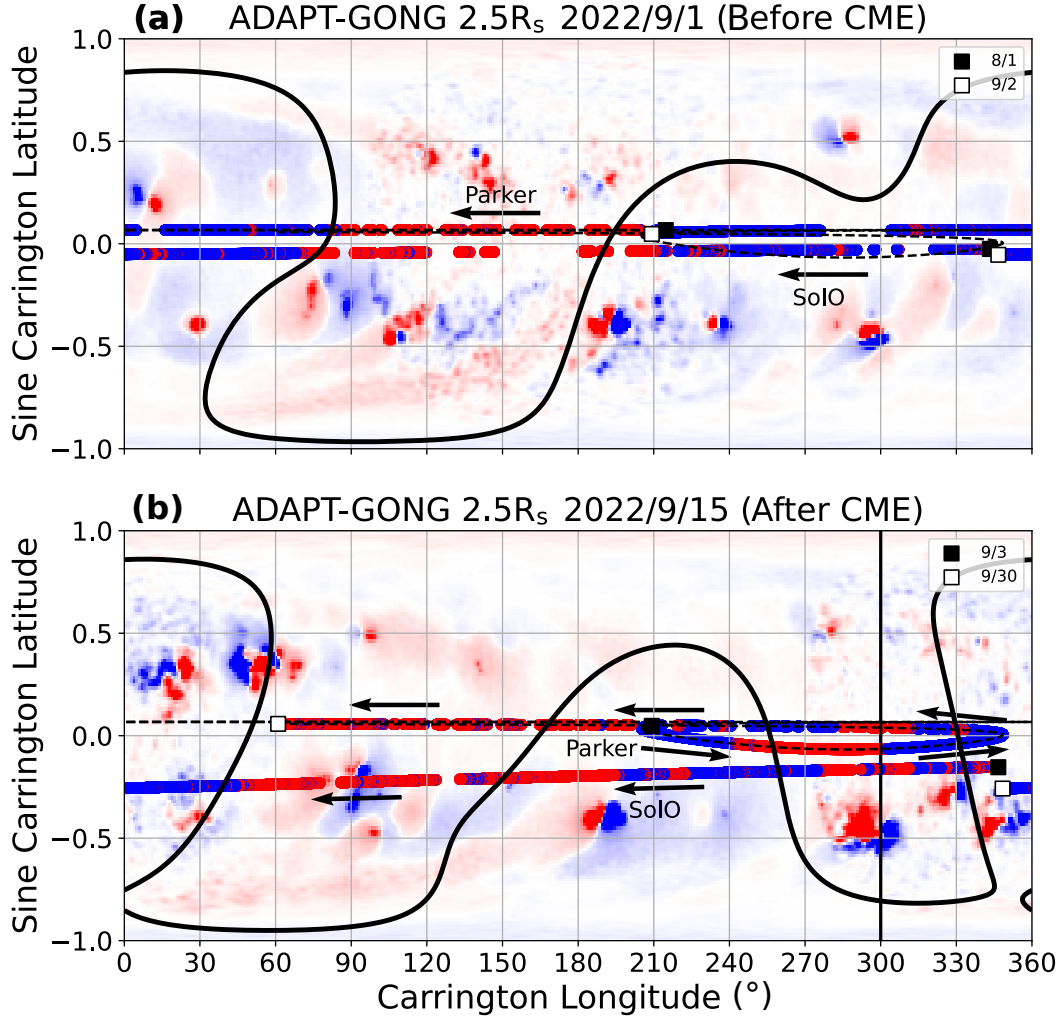


Figure 3.11: ADAPT-GONG magnetograms observed on (a) 01 September 2022 (prior to the CME eruption) and (b) 15 September 2022 (after the CME and rotation of the source active region into view). The red and blue map shading indicates positive and negative radial photospheric magnetic fields, respectively. The black curve represents the nominal HCS obtained through a potential field source surface extrapolation for a given magnetogram. The trajectories of PSP and SOLO are ballistically mapped down to the same altitude as the HCS and colorized according to the measured polarity in space. The black and white squares indicate the initial and final positions of each orbit, respectively, with arrows depicting the orbital direction. The in situ data in panel (a) show a two sector structure of the magnetic field over the course of the Carrington rotation prior to the CME. Four polarity sectors emerge after the CME, where one of the positive sectors is aligned with the location of the source active region (AR 13102) near $(\sin(\theta_{Carr}), \phi_{Carr}) \cong (-0.5, 300^\circ)$, indicating a large-scale reconfiguration of the HCS.

The magnetic field has a two sector structure over the course of the Carrington rotation prior to the ICME. We define a sector as a region of a dominant polarity of the radial field, bounded by the HCS and the solar equator. After the event, the magnetic field substantially changes globally, shown in Figure 3.11b with four sectors. A new sector of positive polarity emerged above 180° in Carrington longitude that is aligned with an active region (AR 13102) at $(\sin(\theta_{Carr}), \phi_{Carr}) \cong (-0.5, 300^\circ)$. The vertical black line in Figure 3.11b depicts the location of PSP at time t_{7a} when the spacecraft crosses the second magnetic inversion line. This would suggest a strong coronal reconfiguration occurred during the CME event, and a new large scale warp in the HCS was introduced by the CME advecting the neutral line of the source active region out into space.

3.6 Discussion

Numerous studies have compared remote sensing and in situ measurements of ICMEs near Earth (e.g., Bothmer et al. 2017; Möstl et al. 2009; Rouillard et al. 2010). PSP provides a new viewpoint near the Sun, which can help constrain ICME structure and dynamics earlier in its evolution. The spacecraft has observed a few ICMEs earlier in its mission (e.g., Lario et al. 2020; McComas et al. 2023; Nieves-Chinchilla et al. 2022), but we expect to observe more events as the Sun approaches solar maximum. This specific event is unique in that PSP traversed through the flank and wake of the structure for nearly two days (a novel situation enabled by the uniquely low perihelion and high prograde angular velocity of PSP), resulting in an interesting combination of in situ measurements which are difficult to fully explain with simple modeling approaches. Based on the animation of Figure 3.10, some aspects of the event clearly match between the in situ measurements, remote sensing observations, and CME modeling. Both the coronagraphs and modeling report an eruption time of the event around 16:10 UT. COR2 and WISPR also observe a large quasi-spherical front passing PSP near time t_1 , when the spacecraft was $\sim 15.05 R_S$ from the Sun. After performing a shock geometry analysis using SPAN and FIELDS data, we find a non-radial shock normal near PSP at time t_1 . As PSP is closer to the flank than the apex of the ICME, we would not expect the shock normal to be radial. Between times t_1 and t_{2a} , the magnetic field fluctuations increase due to the turbulence of the sheath region. We also observe 90° ePADs for energies primarily above 750 eV. These electrons may be heated due to adiabatic compression from the shock (Manchester et al. 2012). Liu et al. (2022b) has suggested the enhancement of perpendicular suprathermal electrons from interplanetary shocks could be attributed to whistler heat flux instabilities through pitch-angle scattering along with normal betatron acceleration of electrons. Furthermore, PSP measures enhanced wave activity in the AC spectra (not shown) from FIELDS during this interval. Thus, the period from t_1 to t_{2a} is denoted as the ICME sheath region, characterized by a transition to more turbulent magnetic field measurements.

After the sheath region, PSP observes some classical signatures of MCs, such as magnetic field enhancements, small rotations, low plasma β , and low T_p . However, not all of these

features are measured simultaneously, or persist for the same duration. Moreover, we do not see a large rotation in the magnetic field typically associated with flux ropes. One reason that PSP does not observe a MC could be due to the spacecraft's distance from the center of the ICME (Cane et al. 1997; Kilpua et al. 2011), or due to gradual erosion of magnetic flux from reconnection at the MC frontal edge (Dasso et al. 2007; Ruffenach et al. 2012). This may explain why the dominant polarity prior to the first polarity inversion (negative) is inconsistent with the polarity of the leading edge of the source active region (positive). In addition, it is possible that PSP is within another region of the ICME outside of the MC. Kilpua et al. (2013) has suggested additional regions exist within an ICME, defined as the faint MC front region (between the ICME and MC front boundaries) and the broader MC rear region (between the ICME and MC tail boundaries). Statistically, front regions have been shown to have higher n , v , T , and B values than other ICME regions, while the rear region is categorized by much lower n , v , and B . The MC is recorded to have the lowest T , with higher n_a/n_p than the front region.

In addition, we do not observe bidirectional strahl during the majority of this interval, but rather diffuse ePADs. One possible explanation could be that the magnetic field lines have reconnected in the CME legs with open field lines (Gosling et al. 1995), or that electron flows are typically stronger in the direction that corresponds to the local flow of the nearest footpoint (Pilipp et al. 1987). The only interval to show strong bidirectionality is from t_{3a} to t_{3b} , and corresponds to PSP slightly grazing the flank of the CME1 model within an error of 1-2 hours. The small rotations in $\mathbf{B}$, low plasma β , and low T_p may indicate the presence of a MC at time t_{3a} , but the measured density is much lower before this interval. Eventually, PSP crosses a magnetic inversion line at time t_{4a} , entering a region of more turbulent and dense plasma. This interval may be associated with turbulence seen in the flank and wake regions of ICMEs (Fan et al. 2018). After t_{5a} , v and T decrease until the plasma is observed to return to typical values for ambient solar wind at these heliocentric distances. The spacecraft may be within a region where new solar wind has begun to fill the ICME wake. However, PSP appears to observe a new transient with sub-Alfvénic flow at time t_6 . PSP crosses a second magnetic inversion line at time t_{7a} . The two portions of the ICME may be separated by this current sheet, where one side is observed to be highly sub-Alfvénic. However, one must be careful when interpreting CME2 as it does not have an associated active source region based on these observations. This could suggest that CME1 and CME2 are produced by the same active region, erupting at different times. CME2 could also be a stealth CME, which is an eruption with no low coronal signatures (LCSs) to identify a source region (Ma et al. 2010; Robbrecht et al. 2009). However, Alzate et al. (2017) has associated these type of events to limitations regarding remote sensing observations and image processing. The LCSs related to stealth CMEs may form higher in the corona or in low density regions (Alzate et al. 2017). Observations from SOLO may provide more insight on CME2 and any associated LCSs as the spacecraft observed the backside of the Sun during the event (Long et al. 2023, *submitted*).

Figure 3.12a illustrates a simple interpretation of the event informed by the ADAPT-GONG magnetograms, PFSS model, and PSP in situ data. We consider a single CME,

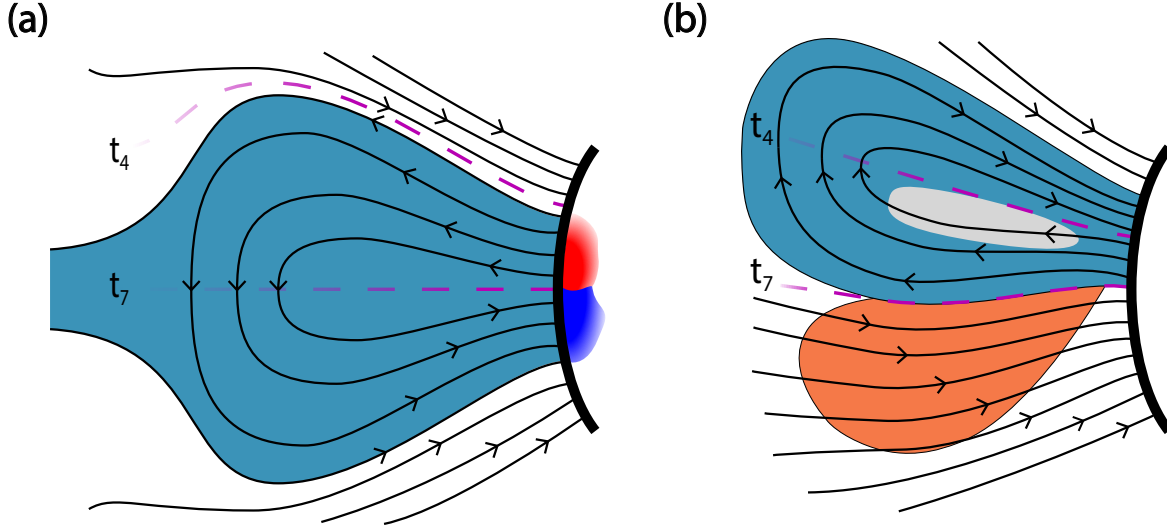


Figure 3.12: Illustrations of **(a)** a single ICME model based on coronal magnetograms and in situ data, and **(b)** a simplified two-part ICME model based on white-light coronagraphs and in situ observations. PSP travels along a circular orbit around the Sun from top to bottom in each figure. The magnetic field lines are represented by the black curves connected to the solar surface, with arrows depicting their respective polarity. The purple dashed lines depict the magnetic inversion lines observed by PSP at times t_{4a} and t_{7a} . The single ICME in (a) is portrayed by the blue region, while CME1 and CME2 in (b) are shown by the blue and orange regions, respectively. The active region AR 13102 is shown on the solar surface in (a) with positive (red) and negative (blue) polarity to agree with the ADAPT-GONG magnetograms. The grey area in (b) represents the region of space outside of CME1 that PSP travels through between times t_{4a} and t_6 . CME1 could be concave in three dimensions, allowing PSP to exit and re-enter the structure.

where PSP travels along a circular orbit from top to bottom in the figure. We refrain from including the orbit in the illustration as the structure evolves over the course of two days. After exiting the sheath region, PSP travels through plasma perturbed by the ICME until crossing the first ICME leg at time t_{4a} . The magnetic field lines are directed away from the Sun in the first leg (blue region) after PSP crosses the first magnetic inversion line (purple dashed line). PSP is within the ICME leg until t_{7a} , when the spacecraft crosses another magnetic inversion line. During this interval, PSP is located directly above the active source region described in Section 3.5. The two polarity regions (red and blue) of AR 13102 may correspond to each leg of the ICME, indicating that PSP travels from one leg to the other at time t_7 . This inversion line would then correspond to the current sheet of the ICME (Lin et al. 2000) in its wake connected back towards the Sun. The sub-Alfvénic region after t_{7b} may be due to an asymmetry in the flow directions during magnetic reconnection (Murphy

et al. 2012). The interval between t_{7a} and t_{7b} could be related to a reconnection exhaust as T_p and plasma β are higher inside this period, while $|\mathbf{B}|$ is smaller (Gosling et al. 2005; Phan et al. 2022).

This illustration of the event agrees well with the magnetograms and FIELDS data, but does not explain the plasma region from t_{2a} to t_{4a} . We would expect PSP to be within the ICME region as this interval exhibits multiple ICME in situ signatures. However, this spatial domain could be between the sheath and MC regions as PSP only travels through the flank and wake of the structure. In addition, this model does not explain the stable plasma conditions observed between t_{5a} and t_6 . The ICME could have a more complex three dimensional structure, experiencing some deformation, deflection, or rotation as it propagates from the Sun. Furthermore, this interpretation does not include a second portion of the CME, modeled to agree with white-light images. A slightly more complex interpretation with a two-part CME, as informed by the GCS modeling discussed in Section 3.4, is shown in Figure 3.12b. The illustration portrays a deformed ICME to emphasize asymmetry within the structure. After time t_{2a} , PSP passes through the flank of CME1 (blue region), observing numerous ICME signatures described above. At time t_{4a} , the spacecraft exits the ICME and enters a region of space still perturbed by the eruption (grey region). This region could exist if the ICME is concave, signifying a more complex structure in three dimensions. Once the spacecraft exits the ICME, the magnetic field lines are inverted, indicating a magnetic inversion line. The perturbed plasma begins to fade away at time t_{5a} , when ambient solar wind fills the ICME wake. Plasma parameters such as n , v , and T return to expected conditions based on solar distance during this interval. PSP re-enters CME1 at t_6 , with field lines still oriented away from the Sun. The spacecraft enters CME2 (orange region) at time t_{7a} , crossing another magnetic inversion line. The region after t_{7b} is highly sub-Alfvénic and could be caused due to the expansion of CME2 within the wake of CME1 (Chané et al. 2021). This interval persists until time t_9 when PSP exits the entire ICME perturbed plasma. This model agrees with the majority of in situ signatures observed by PSP and the inclusion of CME2 from Section 3.4. However, Figure 3.12b does not set the second magnetic inversion line as the current sheet between the CME legs, but instead the current sheet between the two parts. If CME1 and CME2 originate from the same source region, we would expect these structures to be connected.

Both illustrations in Figure 3.12 offer reasonable interpretations for different aspects of this event. A more complete description would need to be more complex in three dimensions to explain the evolution of the structure as it propagates from the Sun. While not every small-scale feature of the event can be explained with the proposed models, it is clear that the HCS experienced a major reconfiguration during this period. The coronal magnetic field had a four sector structure after the event, persisting for about two months, as evidenced independently by PSP and SOLO magnetic field data. Based on the magnetograms and PFSS model, this corresponds to time t_7 when PSP is directly above the active region aligned with one of the new sectors. Thus, this CME event led to a structural change in the HCS that continued well after the eruption. CMEs release accumulated magnetic helicity and trapped energy from the corona (Low 1994). This phenomenon continues the magnetic flux transport

that began from the solar interior out to interplanetary space, and plays a strong role in controlling the magnetic flux budget of the Sun (Low 2001). As the Sun approaches solar maximum, the frequency of CMEs will increase and the cumulative effect of these eruptions will reconfigure the HCS more frequently until the coronal field switches polarity. Therefore, the release of magnetic helicity through CMEs could predominately control the shape of the HCS and its evolution over the solar cycle.

3.7 Conclusion

On 05 September 2022 at 17:27:20 UTC, PSP observed an ICME event until 07 September 2022 at 11:30:16 UTC. Numerous in situ signatures that typically correspond with ICMEs were observed, such as bidirectional strahl, low T_p , low plasma β , $n_a/n_p > 8\%$, directional discontinuities, and a fast-forward shock. The magnetic field varied from region to region with periods of enhancements, small rotations, and turbulence. In addition, PSP measured intervals of density depletion corresponding to sub-Alfvénic regions. Besides in situ measurements, several remote sensing observations of the event were obtained from multiple viewpoints including PSP, STEREO-A, and SOHO. Based on white-light coronagraphs, two portions of the ICME were modeled separately to fit the extent of the event.

Due to the complexity of the ICME, a simple description of the event has proven difficult when comparing in situ and remote sensing measurements. Both proposed scenarios (see Figure 3.12) possess aspects that agree with PSP observations, but neither offers a complete explanation of the event. It is important to note that these interpretations only consider the large-scale structure, and that the full description of the event is more complex in reality. Regardless of the necessity for a second portion of the CME, PSP still crossed two magnetic inversion lines near the Sun in the wake of an ICME. The in situ polarity data taken crossing the ICME-perturbed longitudes, as well as PFSS extrapolations, show a global reconfiguration of the HCS after the ICME event. This would suggest these eruptions on the Sun strongly affect the HCS structure. Therefore, CMEs may have a key involvement in the evolution of the HCS throughout the solar cycle. In addition, the events described in this study may provide insight on how CMEs influence the development of young solar wind, e.g., during the extreme low density region after t_7 , addressing a key science objective of the PSP mission.

To further investigate this event, future work could implement in situ plasma parameters observed by PSP into the CME forward modeling described in Section 3.4. The models presented in this study only considered observations from coronagraph and heliospheric imagers. The inclusion of PSP plasma data during different periods of the event may produce improved results. In addition, future study of the ICME could incorporate measurements from the Integrated Science Investigation of the Sun (IS $\odot$ IS; McComas et al. 2016) instrument suite. IS $\odot$ IS is composed of two instruments that measure lower (EPI-Lo) and higher (EPI-Hi) energetic ions and electrons above the energy range of SPAN. Although EPI-Lo was not operational during this event, EPI-Hi was powered on and collected measurements

of energetic particles. A comparison of SPAN, FIELDS and EPI-Hi data would provide more information on various features of the event, such as high energy particle acceleration. In addition to PSP, SOLO observed the event on 06 September 2022 around 12:33 UT with in situ detectors. SOLO was located $\sim 150 R_S$ from the Sun at the start of the event, and was closer in proximity to the ICME apex than PSP. Further research with SOLO and PSP data could study CME morphology and evolution as the structure propagates farther from the Sun into interplanetary space.

3.8 Appendix: QTN Electron Density Algorithm

To determine the total electron density n_e , we implement a new simplified heuristic algorithm based on Quasi-Thermal Noise (QTN) spectroscopy theory (Meyer-Vernet et al. 2017). The electron plasma frequency f_p is derived from the peak with the slope of greatest ascent in a given QTN spectrum measured by a dipole antenna (see Figure 2 in Meyer-Vernet et al. 2017). By isolating the peak from other features in the spectrum, we can calculate the total electron density. To measure QTN spectra, we use the Radio Frequency Spectrometer (RFS) from the FIELDS suite that includes four electric antennas and a single axis of the PSP search coil magnetometer (Pulupa et al. 2017). We average the RFS Low Frequency Receiver (LFR) data from the V1-V2 and V3-V4 antenna dipole pairs, covering a frequency range of 10 kHz–1.7 MHz with 64 logarithmically spaced frequencies, providing $\sim 4.5\%$ resolution. The LFR spectra are then flattened by removing a fitted quadratic function to emphasize the plasma frequency peak. We convolve the scaled spectra with a step function and apply high and low pass filters centered at 80 and 700 kHz, respectively, to highlight and smooth the plasma frequency line. Finally, we normalize the convolution for each sample by the maximum value in the spectrum.

The heuristic search algorithm is composed of three parts: an outward search, an inward search, and a comparison between the two methods. Beginning with the outward search, we identify all the local maxima for N convolved normalized spectra in a given day. A reference plasma frequency f_r is selected from an ideal spectra with only one peak (the plasma line) at time t_r . From time t_r , an outward heuristic search is iterated forward in time to select frequencies along the plasma line until time t_N . The heuristic function implements both the local maxima values of the signal's convolution and the relative change in frequency between samples. This method assumes a smooth and continuous plasma line, which accounts for solar Type III radio bursts that contaminate the plasma line.

Given the plasma frequency at t_r , we determine f_p at time t_{r+1} by identifying each frequency bin number F (among the 64 frequency bins of the LFR spectra) with a local maxima m at time t_{r+1} . We then compute the normalized difference in frequency bins $\Delta F(t_{r+1})$ between the previous plasma line frequency bin $F_p(t_r)$ and the frequency bin for a given local maxima $F_m(t_{r+1})$ as

$$\Delta F(t_{r+1}) = |F_p(t_r) - F_m(t_{r+1})|/M \quad (3.1)$$

where M is the total number of frequency bins. For each local maxima, we apply $\Delta F(t_{r+1})$ to a frequency heuristic function $h_F(\Delta F)$, shown in Figure 3.13. It is important to note that we are using the frequency bin number F (1 through 64) and not the numerical frequency value f for the heuristic function to avoid a bias towards smaller frequencies. In addition, for each local maxima, we apply its normalized convolution value $c(t_{r+1})$ to a convolution heuristic function $h_c(c)$, shown in Figure 3.13. An outward search heuristic function is then computed weighing both parameters as $h = 0.4 h_c + 0.6 h_F$. Based on Figure 3.13, the local maximum with a combination of the largest convolution value and smallest frequency bin difference is assigned the plasma frequency f_p at time t_{r+1} . This process is iterated forward in time until t_N , and iterated backward from time t_r to t_1 in a given day. Given that this algorithm identifies the plasma line from time t_r towards the start and end of the day, it is referred to as an “outward search.”

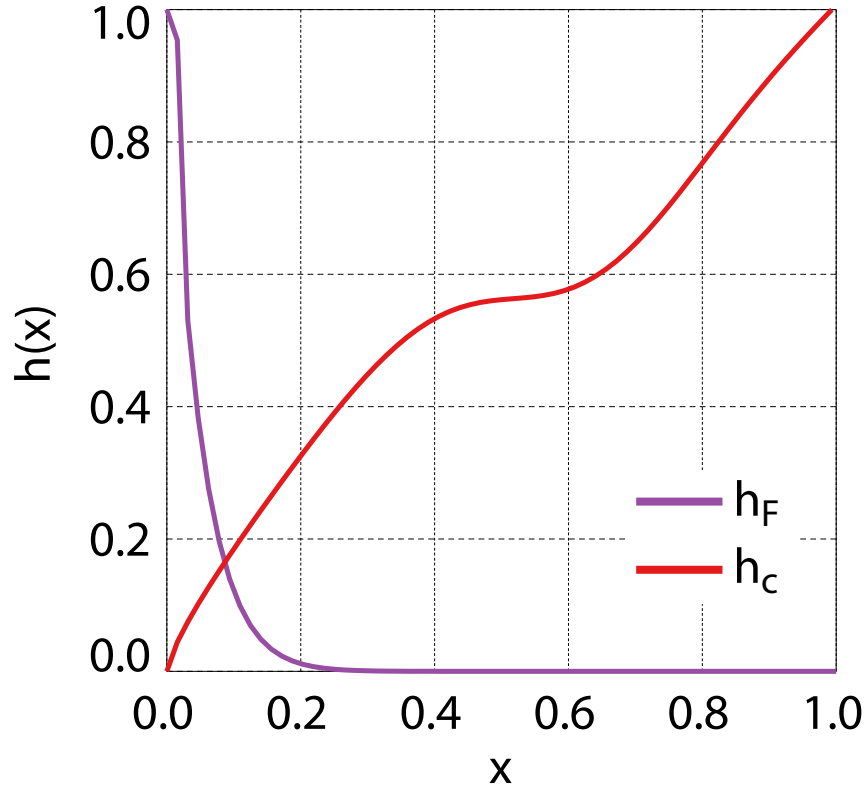


Figure 3.13: Heuristic functions $h(x)$ for a given normalized parameter x . The frequency heuristic function h_F (purple) is a function of the normalized frequency bin difference ΔF . The convolution heuristic function h_c (red) is a function of the normalized convolution value of the resulting spectra.

The “inward search” algorithm is performed in a similar manner, except that the search begins at time t_1 . We select a new reference plasma frequency at the beginning of the day and iterate forward in time to t_r using the same heuristic functions in Figure 3.13. In parallel, another search is performed from time t_N to t_r , iterating the process backwards in time to determine the plasma line. We manually select the plasma frequency at times t_1 and t_N given the spectra for a complete day. This process is known as an “inward search” as both searches are iterated separately towards time t_r . The purpose of bidirectional searches is to more easily determine the plasma line in cases of apparent bifurcation. In theory, the resulting frequency from either the outward search f_p^o or inward search f_p^i is the best estimate for the true plasma frequency. The final step of the algorithm compares the plasma frequencies from the two searches in an effort to determine a single value of f_p at each time.

Figure 3.14 emphasizes the importance of comparing the outward and inward searches for the plasma line. Outside the interval between time t_a and t_b , both search algorithms agree upon the estimated plasma frequency. However, within this interval, the outward and inward searches follow different paths due to a Type III radio burst. To account for this discrepancy ($f_p^o \neq f_p^i$), we compute a new modified heuristic function for each search, solely based on the relative change in frequency bins between t_a and t_b . The inward h^i and outward h^o heuristic functions are defined as

$$h^i = \sum_{t=a}^b \left\{ |F_p^i(t) - F_p(t_a)|^4 + |F_p^i(t) - F_p(t_b)|^4 \right\} \quad (3.2)$$

$$h^o = \sum_{t=a}^b \left\{ |F_p^o(t) - F_p(t_a)|^4 + |F_p^o(t) - F_p(t_b)|^4 \right\} \quad (3.3)$$

where F_p^o and F_p^i are the plasma frequency bins determined from the outward and inward searches, respectively. The search method that results in the smallest heuristic value will be used to set f_p . Since $h^o < h^i$ within the time interval in Figure 3.14, the outward search plasma frequency is selected. This is done for all time intervals that the inward and outward search algorithms do not agree.

Once the plasma frequency line is determined for a given day, we compute a geometric weighted mean frequency $\langle f_p \rangle$ between the neighboring frequency bins of the plasma line. This process is then repeated for each day to determine the total electron density with the equation $n_e = (\langle f_p \rangle / (8.98 \times 10^3))^2$. To ensure our algorithm is performing correctly, we use the QTN electron density estimate from Moncuquet et al. (2020) for PSP encounters 1 through 10. We have determined the frequency and convolution heuristic functions shown in Figure 3.13 by comparing the results of Moncuquet et al. (2020). The two data sets agree within 5 – 10%, which is reasonable given the 4.5% resolution of the instrument. This method offers several advantages, such as its simplicity, efficient run-time of 0.5–2 minutes per day, and handling of radio bursts, large fluctuations in the plasma line, and other events with low signal-to-noise ratio.

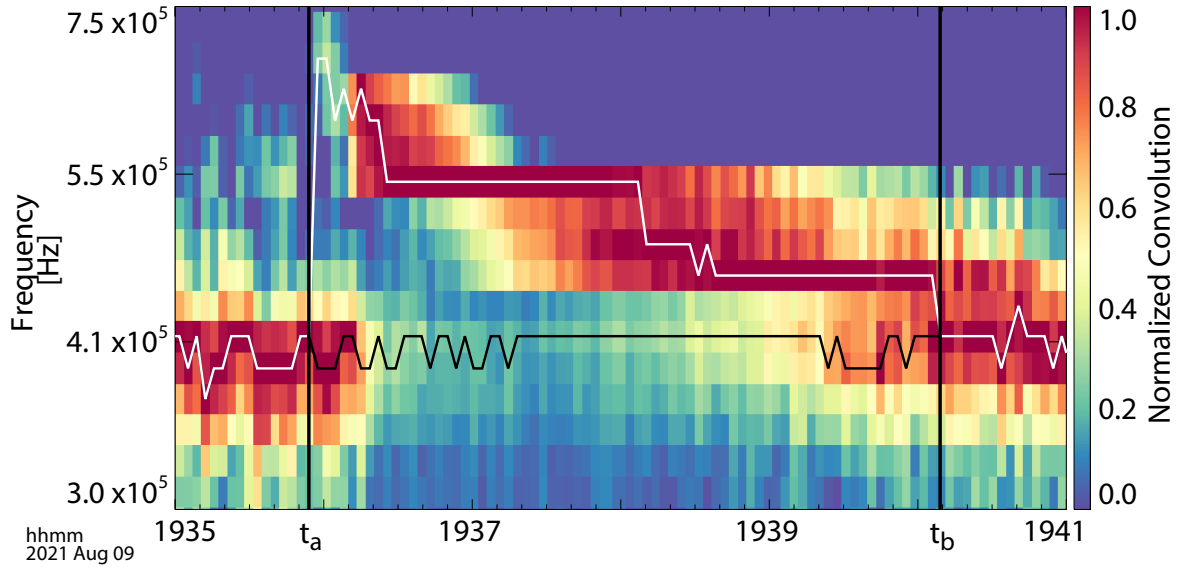


Figure 3.14: Normalized convolution spectra from RFS LFR data on 09 August 2021, targeted on the plasma line. the plasma frequencies selected from the inward and outward searches are designated by the white and black lines, respectively. Within the interval between times t_a and t_b , the two search algorithms do not follow the same path.

3.9 Appendix: Empirical Solar Wind Velocity & Temperature Relationship

Based on the bulk proton speed v_p , an empirical relationship can be found with the plasma proton temperature T_p for the ambient solar wind (Lopez 1987). However, there are instances when the measured T_p is lower than the expected proton temperature T_{exp} from the $v_p - T_p$ relation. These intervals could be associated with magnetically isolated, ejected plasma expanding faster than the ambient solar wind. They typically are preceded by interplanetary shocks, suggesting a connection to ICME shock drivers. We can expect ICMEs to follow the condition $T_p < 0.5 T_{exp}$ for magnetic clouds (Richardson et al. 1995).

Previous work has correlated v_p and T_p using Wind data (Huang et al. 2017; Yu et al. 2014). However, we must determine a new empirical relationship for v_p and T_p as PSP orbits much closer to the Sun than Wind. Thus, we follow a similar method from these studies to fit the expected proton temperature to the bulk proton speed using SPC and SPAN-I data for encounters 1 through 10. The data set is divided into two velocity regimes: slow solar wind ($v_p < 400 \text{ km s}^{-1}$) and fast solar wind ($v_p > 400 \text{ km s}^{-1}$). We perform a separate fit for each velocity regime based on the statistical formula from Lopez (1987). In addition, as the temperature scales with radial distance R from the center of the Sun, we include a power

law R^a in each fit. Figure 3.15 displays the empirical relationship between v_p and scaled T_p by R^a for both the slow and fast wind, described by

$$T_{exp} R^{0.79} = (0.122 v_p + 21.41)^2 \cdot 10^3 \quad (v_p < 400 \text{ km s}^{-1}) \quad (3.4a)$$

$$T_{exp} R^{1.10} = (0.537 v_p + 86.80)^2 \cdot 10^3 \quad (v_p > 400 \text{ km s}^{-1}) \quad (3.4b)$$

where R is in units of solar radius R_S . Based on these relationships, we use v_p to compute an expected proton temperature and compare T_{exp} with the measured proton temperature from SPAN-I.

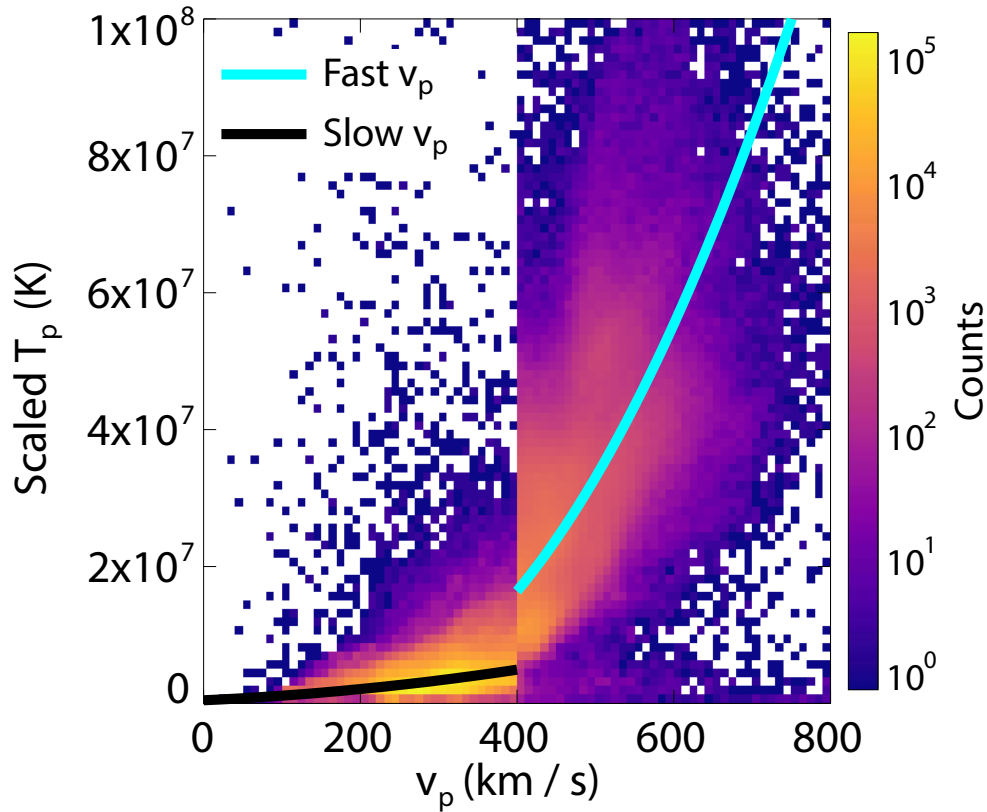


Figure 3.15: 2D Histogram of scaled proton temperature T_p by R^a as a function of SW velocity v_p for PSP encounters 1 through 10. The proton temperature is scaled by $R^{0.79}$ in the slow solar wind and $R^{1.10}$ in the fast wind. The fitted curves for the slow and fast solar wind are shown in black and cyan, respectively.

Chapter 4

Scales of Martian crustal magnetization constrained by MAVEN, InSight, and Zhurong

4.1 Abstract

While Mars does not possess a current active geodynamo, remanent crustal magnetization has been found across the planet and contains records of the origin, scale, and timing of Martian magnetization. The first in situ measurements of the Martian crustal magnetic field at the InSight and Zhurong landing sites allow for better constraints on magnetization coherence and depth scales near the surface, as crustal fields are closely related to a variety of geological and topographic features. We develop Monte Carlo models of the Martian crustal magnetization near the two landing sites on small-scales (< 50 km) to meso-scales (100–1000 km) to compute altitude profiles of the magnetic field intensity. By comparing our simulations with MAVEN orbital magnetic field observations, we determine a 20 km thick magnetized source close to the surface with coherence scales roughly 275 km near InSight, with larger coherence scales and possibly thicker crustal magnetization near Zhurong. Motivated by these length scales, we create additional magnetization models based on the geological unit regions near each lander to relate them to different time periods of Martian history. Our results suggest at least one polar field reversal in Martian history based on the simulated magnetization near the North-South dichotomy boundary. Furthermore, we propose that the Martian geodynamo weakened or suspended during the late Noachian, followed by the reactivation of a new core dynamo during the Hesperian period.

4.2 Introduction

In contrast to other terrestrial planets such as Earth and Mercury, Mars currently lacks an intrinsic global dynamo-driven magnetic field (Acuña et al. 2001; Smith et al. 1965).

However, the discovery of localized crustal magnetic fields near the Martian surface has provided evidence for the existence of an internal geodynamo during a previous stage of Mars' history (Acuña et al. 1998; Mittelholz et al. 2022). Investigating the nature of these crustal fields can offer insight into the timing of the Martian dynamo, which has been linked to global climate evolution on Mars (Chassefière et al. 2004; Jakosky et al. 2001). Furthermore, detailed analysis of crustal magnetic fields can be used to study the composition, structure, and formation of the Martian crust (e.g., Connerney et al. 2001; Lillis et al. 2013; Moore et al. 2017; Vervelidou et al. 2017).

Crustal magnetization can occur on Mars and other planetary bodies through various processes, collectively known as natural remanent magnetization (NRM), to acquire and preserve strong magnetic fields near the surface (Dunlop et al. 1997). Material that is cooled below the Curie temperature can become magnetized through thermal remanent magnetization (TRM) from the global field present during cooling. Shock energy from impacts can also alter surface magnetization through shock remanent magnetization (SRM), either enhancing or diminishing the local field depending on the large-scale field surrounding the impact site. In addition, chemical remanent magnetization (CRM) can arise from rock-water interactions, which are influenced by factors such as oxygen fugacity and the rate of cooling (Grant 1985; Quesnel et al. 2009). These NRM processes can be connected to large-scale geological features on Mars, such as the North-South dichotomy, impact basins, and volcanic regions (Smith et al. 1999).

To estimate the global distribution of crustal fields on the Martian surface, numerous magnetic field models have been created based on satellite observations (e.g., Arkani-Hamed 2004; Gao et al. 2021; Langlais et al. 2019; Morschhauser et al. 2014). These studies have utilized orbital magnetic field measurements from the Mars Global Surveyor (MGS; Albee et al. 1998) and Mars Atmosphere and Volatile Evolution (MAVEN; Jakosky et al. 2015) missions. MGS and MAVEN have sampled the magnetic field at altitudes above the surface as low as ~ 90 km and 110 km, respectively. The surface magnetization maps derived non-uniquely from orbital magnetic field data can be used to examine relationships between geomorphic features and areas of strong or weak magnetization on a global scale. Varying spatial coherence scales of crustal fields can then be associated with different NRM processes that preserved these fields since the ancient Martian geodynamo. Unfortunately, the spatial resolution of these models is approximately limited to orbital observation altitudes. Therefore, new models using surface measurements, such as magnetometer data from the Interior Exploration using Seismic Investigation, Geodesy and Heat Transport (InSight; Banerdt et al. 2020) and Zhurong (Wan et al. 2020) missions, are valuable to better constrain the magnetic coherence scales near each landing site because these landers are much closer to the source of crustal magnetic fields.

In this chapter, we follow a similar method presented in Lillis et al. (2010) and Halekas et al. (2010) to determine the small-scale to meso-scale magnetization structure near the Martian surface at each lander site. Section 4.3 describes the surface and orbital magnetic field measurements implemented in this study, providing a comparison with the (Langlais et al. 2019) model of Martian crustal fields. Section 4.4 details the two different methods used

to generate altitude profiles of the crustal field in this study. We first present a Monte Carlo approach using statistical modeling and continue with a geological representation model unique to each landing site. The resulting altitude profiles of each method are compared with the (Langlais et al. 2019) model in Section 4.5. The discussion and conclusions are presented in Sections 4.6 and 4.7, respectively.

4.3 Martian Magnetic Field Observations

Surface Measurements

InSight landed on Mars in western Elysium Planitia at 4.50° N and 135.62° E (see Figure 4.1) on 26 November 2018 to study the interior structure of Mars (Banerdt et al. 2020). The lander is located within a 27 m diameter degraded impact crater, known as Homestead hollow, at an elevation of -2.61 km (Golombek et al. 2020). InSight is equipped with the Auxiliary Payload Sensor Suite (APSS) to measure a variety of environmental conditions in support of the mission’s seismology efforts (Banfield et al. 2018). The APSS suite includes the InSight FluxGate (IFG) magnetometer to provide the first in situ surface magnetic field measurements on Mars (Johnson et al. 2020).

The IFG sensor is a single triaxial fluxgate magnetometer, with a range of $\pm 20,000$ nT and a sample rate of up to 20 Hz (Banfield et al. 2018). To obtain an accurate measurement of the Martian crustal field, the IFG must be calibrated to remove external and instrumental magnetic fields. External fields can originate from the interplanetary magnetic field (IMF), ionospheric currents, and weather events, while the lander can produce fields from the robotic arm, solar arrays, radio communications, and temperature deformations (Mittelholz et al. 2020a, 2023b). After accounting for these sources, Johnson et al. (2020) determined the average remnant crustal field vector measured by InSight as $(B_X, B_Y, B_Z) = (-1353 \pm 19, 1168 \pm 83, -925 \pm 40)$ nT, with a magnitude of 2013 ± 53 nT. In this local lander coordinate frame, X points north, Y is directed east, and Z completes the right-handed system in the vertical direction. The measured field strength is almost an order of magnitude larger than satellite based estimates (Langlais et al. 2019), highlighting the need for more low altitude measurements of the ambient static crustal field at different locations on Mars (Mittelholz et al. 2023a).

Following InSight, the Tianwen-1 mission launched in July 2020 to study the Martian atmosphere, geological structure, and mineral composition with an orbiter, lander, and rover (Wan et al. 2020). The orbiter arrived at Mars on 10 February 2021, with the Zhurong rover landing in the southern Utopia Basin at 25.07° N and 109.93° E on 15 May 2021 (see Figure 4.1), with an elevation of -4.10 km (Liu et al. 2022a). Once the rover was operational, a magnetic survey was performed using the Zhurong Mars Rover Magnetometers (RoMAG) payload, which provided the first magnetic field measurements at 16 distinct locations on the Martian surface (Du et al. 2020).

RoMAG consists of two identical tri-axial fluxgate magnetometers placed at the top and

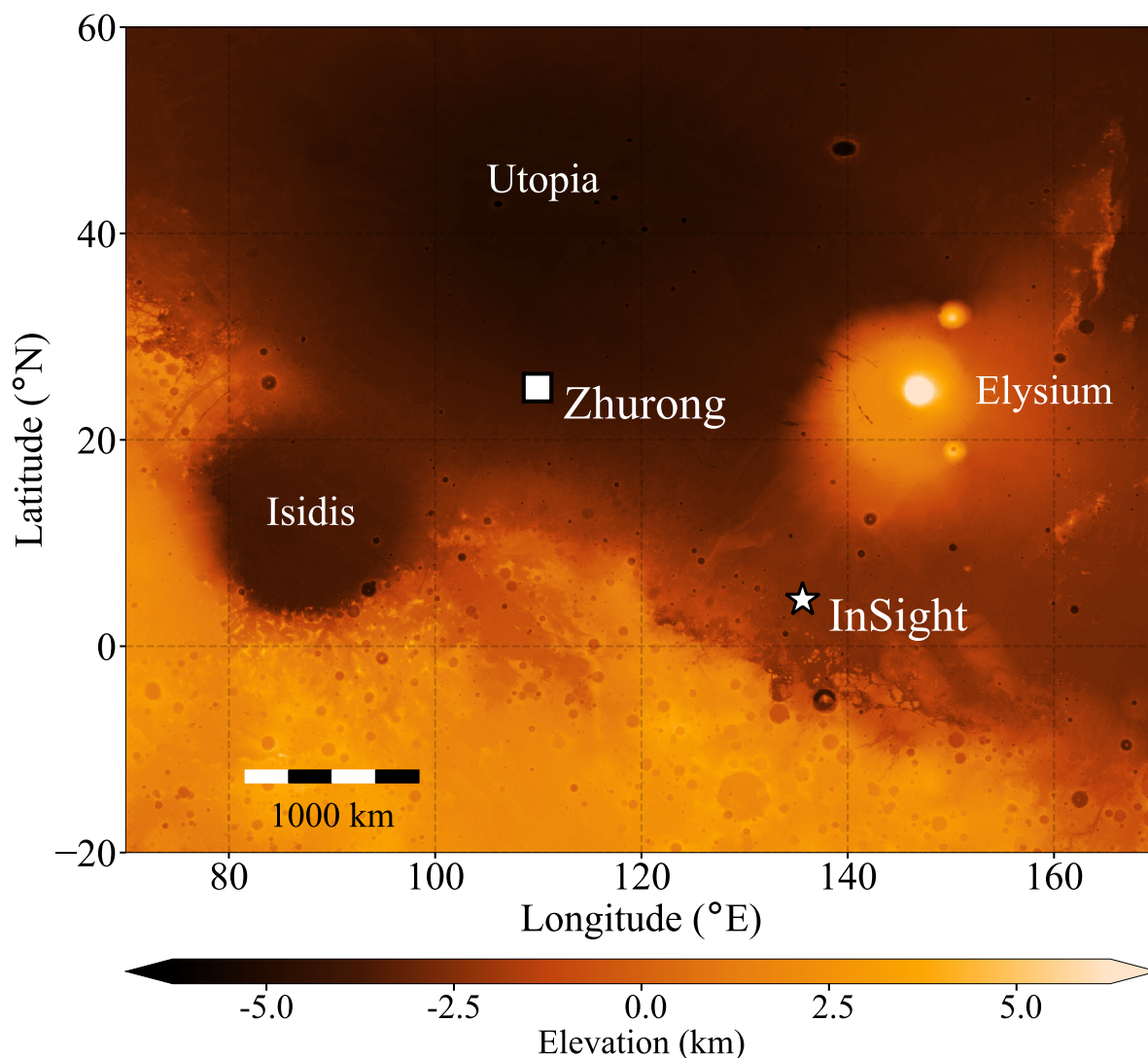


Figure 4.1: Topographic map of Mars near the InSight (white star) and Zhurong (white square) landing sites from the Mars Orbiter Laser Altimeter (MOLA) instrument onboard the Mars Global Surveyor (MGS; Smith et al. 2001). The North-South dichotomy is shown extending from northwest to southeast, with two impact basins Utopia and Isidis and the volcanic region Elysium labeled on the map. The selected elevation scale is set to emphasize the relative difference between the two landing sites.

bottom of the rover mast. The sensors have a range of $\pm 65,000$ nT and a sample rate of up to 32 Hz, with a noise level of 1 Hz (Du et al. 2020). For 90 sols, magnetic measurements were taken at 16 different locations along a 1089 m path within the Utopia basin (Du

et al. 2023). For context, the entire ~ 1 km path that Zhurong traveled falls within one orbital pixel. After accounting for external and instrumental magnetic field contributions in combination with assumptions on external fields, Du et al. (2023) determined an average horizontal crustal magnetic field strength of 11.2 ± 10.9 nT. The vertical magnetic component could not be calculated due to the orientation of the two magnetometers. The horizontal magnetic field magnitude B_H was calculated using the individual arithmetic mean of the north (-5.1 ± 13.6 nT) and east (9.9 ± 11.0 nT) magnetic field components separately. However, to properly compute B_H , we first calculate the horizontal magnetic field magnitude at each location to combine both the north and east component distributions, shown in Figure 4.2a as a function of distance along the path of Zhurong. As these values can only be positive ($B_H > 0$ nT) and their uncertainties are comparable to their magnitude, we cannot apply Gaussian statistics to determine the most likely value of B_H . Instead, we compute a Cumulative Distribution Function (CDF) of the B_H measurements and fit the data to a Weibull CDF, given by equation 7.1 in Hallinan (1993) as

$$f(E) = 1 - \exp \left[- \left(\frac{B_H}{a} \right)^b \right] \quad (4.1)$$

where a and b are referred to as the scale and shape parameters, respectively. From equation 4.1, we can compute the mode $\tilde{B}_H$ of the distribution using the following expression

$$\tilde{B}_H = a \left(\frac{b-1}{b} \right)^{1/b} \quad (4.2)$$

in cases where $a > 1$. Using equation 4.1 to perform a least-squares fit on the B_H measurements, shown in Figure 4.2b, we obtain the values $a = 17.8$ and $b = 1.5$, corresponding to $\tilde{B}_H = 9.04 \pm 10.64$ nT based on the standard deviation σ_{B_H} of the distribution, given by

$$\sigma_{B_H} = a \sqrt{\Gamma \left(1 + \frac{2}{b} \right) - \Gamma \left(1 + \frac{1}{b} \right)^2} \quad (4.3)$$

where $\Gamma(z) = \int_0^\infty s^{z-1} e^{-s} ds$ is the Gamma function for an integration variable s .

The Zhurong magnetometers could not measure the vertical component B_V of the crustal fields. However, we make the assumption that the vertical component is roughly on the order of a single horizontal component of the magnetic field vector. The horizontal magnetic field $\tilde{B}_H$ is composed of two components along the planetary surface (north and east directions). Therefore, we can relate the vertical and horizontal components as $\tilde{B}_H = \sqrt{2} \tilde{B}_V$, where the $\sqrt{2}$ term accounts for an averaged single horizontal component. This leads to $\tilde{B}_V = 6.39 \pm 5.32$ nT, and a final estimated crustal magnetic field magnitude of 11.07 ± 9.22 nT near Zhurong. We note that since the error is comparable to the field value, there is a very significant probability of no detectable crustal magnetic field near Zhurong.

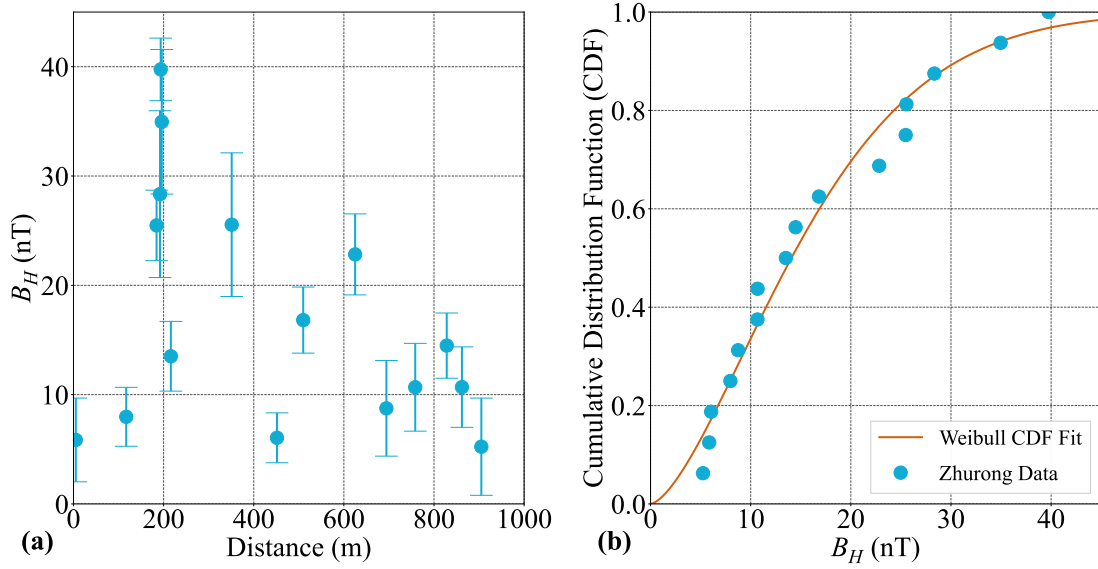


Figure 4.2: (a) Zhurong horizontal magnetic field B_H measurements with uncertainties at 16 different locations along the ~ 1 km path. (b) Cumulative distribution function (CDF) of B_H measurements (blue) fitted to a Weibull CDF (orange), with $a = 17.8$ and $b = 1.5$.

Orbital Measurements

The MAVEN spacecraft began its primary science operation at Mars in November 2014 to study the atmospheric composition, structure, and mechanisms that led to atmospheric loss. Over the course of a decade, MAVEN has executed multiple orbital maneuvers that led to crustal magnetic field measurements over a large range of altitudes. The initial orbit had a nominal periapsis altitude of 150 km, an apoapsis altitude of 6220 km, and a period of about 4.5 hours (Jakosky et al. 2015). MAVEN also performed nine Deep Dip (DD) week-long campaigns, reaching ~ 125 km above the Martian surface. Between February and April 2019, MAVEN executed an aerobraking maneuver to better support relay communications with other Martian missions, lowering the nominal periapsis and apoapsis to ~ 130 km and ~ 4500 km, respectively. After August 2020, the orbital periapsis was raised above ~ 180 km to create a more fuel efficient and stable orbit.

To measure the magnetic field, the MAVEN magnetometer (MAG) instrumentation consists of two fluxgate magnetometers that provide vector magnetic field measurements over a broad range (to 65,536 nT per axis) with a sampling frequency of 32 Hz and accuracy of 0.25 nT (Connerney et al. 2015). In this statistical study, we utilize a subset of night-side MAVEN MAG data from November 2014 through February 2023. Night-side magnetic measurements are identified when MAVEN is within the cylindrical shadow of Mars with respect to the Sun to observe crustal fields with minimal influences from the IMF and ionospheric currents. Due to the spacecraft's orbital precession and various maneuvers over the course

of the mission lifetime, MAVEN infrequently orbits directly above the InSight and Zhurong sites at varying altitudes. Therefore, we compile all night-side orbital measurements within a 200 km radius of each mission landing site to estimate the altitude profile of crustal field strength above each lander on the surface. Selection radii between 50 and 500 km were considered around each landing site, but a radius of 200 km resulted in the smallest area with continuous sampling at all altitudes between 150 and 400 km.

Once the orbital dataset was finalized, we compared several methods to calculate the altitude profile above each lander. For each method, we separate data points in 10 km altitude bins from 150 to 400 km. The first method computed a simple arithmetic average of all data points within 200 km at each altitude. We extended this procedure to geometric and weighted averages based on the relative distance between a given measurement and the lander. We also tried a nearest neighbor approach from each individual point. The second set of approaches we examined averaged the MAG data in 3D voxels with a bin size of 10, 30 and 30 km for altitude, longitude and latitude, respectively. From this sparse 3D grid, we implemented Radial Basis Function (RBF) interpolation to find the altitude profile directly above each landing site (Skala 2016). We compared the generated profiles either interpolated on a 2D grid plane for each altitude separately or a 3D grid simultaneously for all altitudes. The 3D RBF interpolation from spatially binned MAVEN data resulted in consistent and smooth results when varying the selection area near each lander.

Figure 4.3a displays the distribution of the number of MAVEN night-side measurements above InSight and Zhurong at each altitude from 150 to 400 km. While MAVEN did measure crustal fields below 150 km, the lowest orbital altitude within 200 km of each lander with more than 50 measurements was around 150 km. Figure 4.3b shows the resulting altitude profile of magnetic field intensity based on MAVEN orbital data using 3D RBF interpolation. The error bars for each MAVEN data point are derived from the standard deviation of all measurements within each altitude bin. The surface measurements from InSight and Zhurong are also depicted in Figure 4.3b.

To determine whether the interpolated altitude profiles derived from MAVEN data are still affected by external magnetic fields, we compare the data with the Langlais et al. (2019) model (hereafter referred to as L19). The L19 model isolates the crustal fields from the external magnetic fields at satellite altitudes near Mars. For each MAVEN orbital sample, we find the corresponding L19 value at that location. We then perform a similar 3D RBF interpolation from spatially binned L19 values to compute a new altitude profile, shown in Figure 4.3b. For both InSight and Zhurong, the MAVEN data agree with L19 within 3σ , but overestimate the magnetic field profile compared to L19. Therefore, we will use the L19 interpolated altitude profile to characterize the magnetic field above each lander site in later sections.

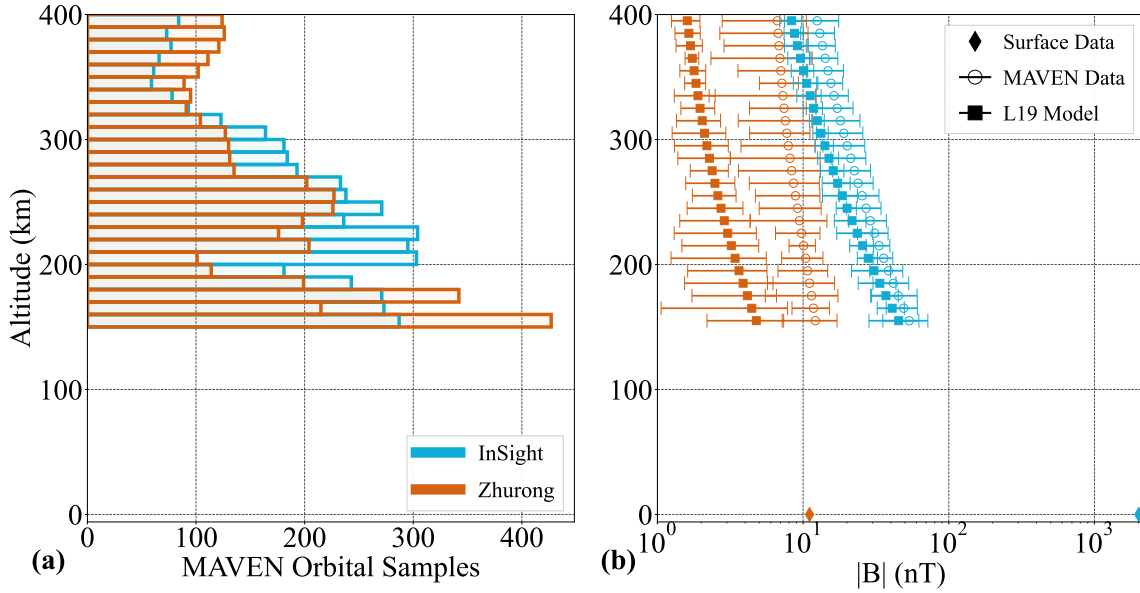


Figure 4.3: **(a)** Histogram of MAVEN magnetic field measurements at each orbital altitude, from 150 to 400 km, above the InSight (cyan) and Zhurong (orange) landing sites. For each 10 km altitude bin, only night-side magnetic field measurements within a radius of 200 km from each mission landing site are considered. **(b)** Altitude profiles of magnetic intensity above the InSight (cyan) and Zhurong (orange) landing sites based on MAVEN data (circles with error bars) and the Langlais et al. (2019) spherical harmonics model (L19; squares with error bars), with surface measurements (diamonds). The error bars for each point at a given altitude are computed as the standard deviation of all night-side measurements within a radius of 200 km around each lander.

4.4 Methodology of Crustal Magnetization Models

Monte Carlo Simulations

We first implement a Monte Carlo approach on simulated crustal magnetization distributions to statistically correlate magnetic and spectral properties of the Martian crust near each lander. The distributions are generated from a Fourier domain expansion of randomized magnetic dipole vectors in a 3D grid. This technique avoids inherent non-uniqueness from subsurface magnetic inversion methods and issues regarding dominating magnetization wavelengths at comparable orbital altitudes (Blakely 1995). The method discussed in this section follows a similar approach by Lillis et al. (2010) to study impact demagnetization on Mars and Halekas et al. (2010) to analyze lunar crustal magnetization.

We first create a 3D rectangular grid of a given size and spatial resolution with equal x and y dimensions. The grid parameters in the z direction (vertical) are controlled separately

since we study the vertical and horizontal domains independently. Certain NRM processes can result in different magnetization scales depending on the influence of gravity, leading to anisotropic behavior. For example, the cooling of magmatic sills may dominate the horizontal magnetization coherence scales, whereas magmatic dikes may have a stronger effect on the vertical scale. After testing several grid sizes in the lateral direction, we set the grid size to 1024×1024 km in the xy plane with a 2 km spatial resolution to compromise between computational runtime and the range of possible coherence scales. The number of voxels in the horizontal plane is fixed to a power of two to improve the computational efficiency when computing Fourier transforms.

Once the 3D grid is defined, we assign each voxel a magnetization vector parallel (or anti-parallel) to the local magnetic field near each lander. The unit vector is defined in the lander coordinate frame discussed above, and set to $\hat{B}_{InSight} = (-0.67, 0.58, -0.46)$ nT based on Johnson et al. (2020) and $\hat{B}_{Zhurong} = (0.0, 0.71, 0.71)$ nT based on Du et al. (2023) for InSight and Zhurong, respectively. The vector magnitude for each voxel is randomly sampled from a Gaussian distribution centered on zero with a width (standard deviation) of 10 A m^{-1} . Other models, such as Lorentzian and uniform distributions, were explored but did not result in significantly different altitude profiles of crustal fields. It is important to note that a negative magnetization strength refers to an anti-parallel vector relative to the local magnetic field direction.

A Fourier transform is then applied on the randomly magnetized 3D grid to the wavenumber domain, calculating the wavenumbers k_x , k_y , and k_z in each direction, where $k = 2\pi f$. As we are only interested in distinguishing the horizontal and vertical directions, we compute a lateral wavenumber $k_L = \sqrt{k_x^2 + k_y^2}$. The vertical wavenumber k_V and lateral wavenumber k_L are used to apply a spatial filter in Fourier space to set different patterns of magnetization in each direction. While numerous coherence filters exist, we focus our work on two well known distributions considered in previous studies. The first is a Gaussian distribution of the form $F(k) \sim e^{-(k-c)^2/2w^2}$, used by Lillis et al. (2010) to study Martian impact demagnetization. The peak of the distribution c is referred to as the coherence wavenumber, centered on a specific spatial wavelength to set the magnetization scale. Following Lillis et al. (2010), we assign the width of the Gaussian filter to be half of the coherence wavenumber. The second distribution follows a power law with the form $F(k) \sim k^{-\alpha}$, used in Pilkington et al. (1993) to examine aeromagnetic and borehole measurements of Earth's lithosphere. The power law filter only depends on α , which represents the slope of the function in the wavenumber domain. As α approaches larger values, the dominance of larger coherence wavelengths on surface magnetization increases. Once the coherence filter is applied, we perform an inverse Fourier transform on the 3D grid and normalize the magnetic field strength to the original maximum value.

Figure 4.4 presents an example of the Gaussian and power law coherence filters applied on a magnetization distribution in the lateral and vertical dimensions. Figure 4.4a shows normalized Gaussian distributions as a function of wavenumber, with peak coherence wavelengths $c_L = 100$ km and $c_V = 10$ km in the lateral and vertical dimensions, respectively. The

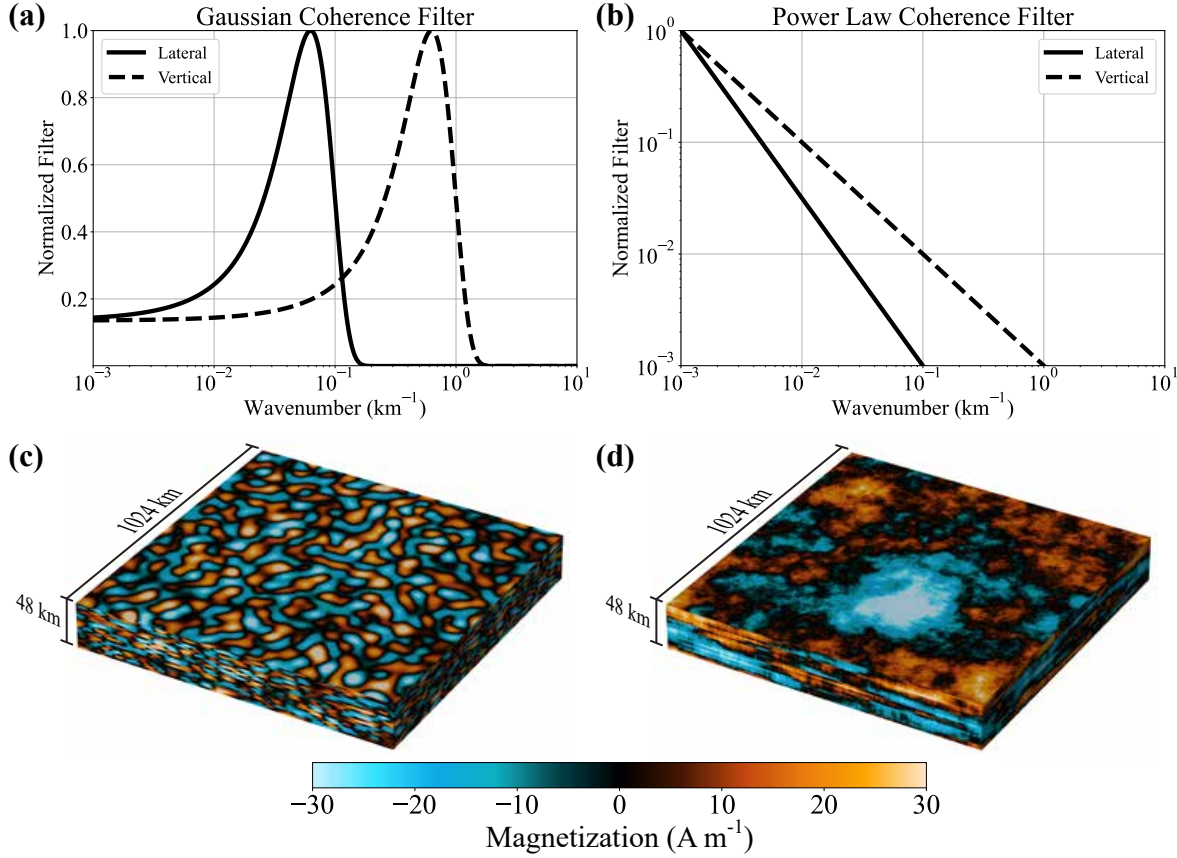


Figure 4.4: Example of normalized coherence filters in the Fourier domain based on a (a) Gaussian and (b) power law distribution. The lateral and vertical Gaussian filters have a peak coherence wavelength of 100 and 10 km, respectively, with a width equal to half of the peak wavelength. The lateral and vertical power law filters are generated with powers of 1.5 and 1.0, respectively. A magnetized 3D grid is then generated with dimensions $1024 \times 1024 \times 48$ km and spatial resolutions of 2 and 1 km in the lateral and vertical directions, respectively. Each voxel in the grid is randomly assigned a magnetization strength from a Gaussian distribution centered on zero with a width of 10 A m^{-1} , where negative values correspond to anti-parallel vectors to the set magnetic field direction. The coherence filters are later applied to each magnetization distribution in wavenumber space for each dimension individually. To obtain the final coherent magnetization distribution from a (c) Gaussian and (d) power law filter, the inverse Fourier transform is performed to the space domain.

widths of the lateral and vertical coherence distributions are set to half of the peak coherence wavelength. We then apply the Gaussian filters along each axis on a randomly magnetized grid, with dimensions $1024 \times 1024 \times 48$ km, in the Fourier domain. Once the magnetization strength is normalized across the entire grid, we perform the inverse Fourier transform to

obtain the 3D simulation box illustrated in Figure 4.4c. Figures 4.4b and 4.4d demonstrate the same procedure using a power law coherence filter with $\alpha_L = 1.5$ and $\alpha_V = 1.0$ for the lateral and vertical dimensions, respectively.

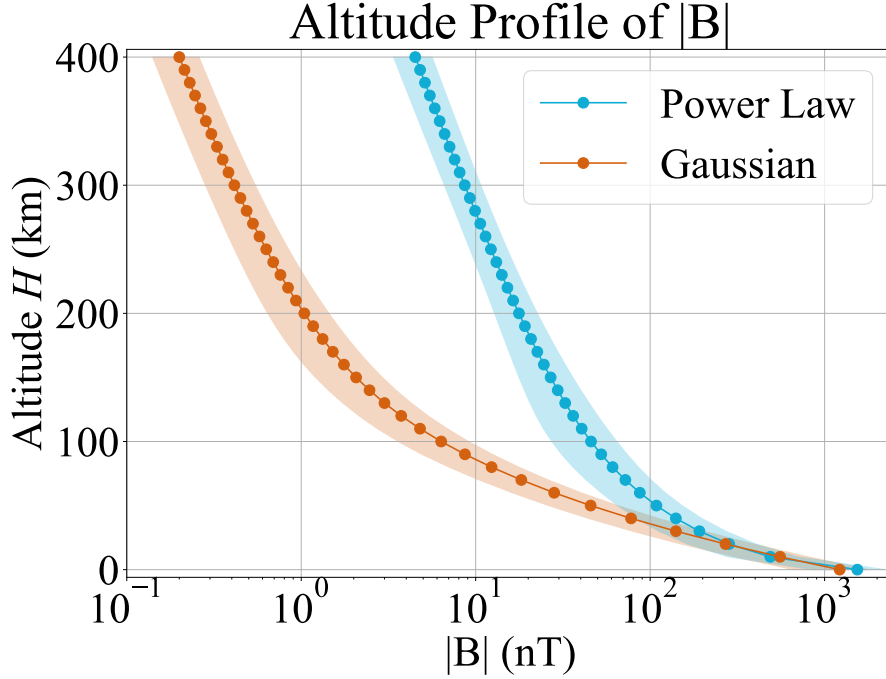


Figure 4.5: Example of altitude profiles of the magnetic field strength $|\mathbf{B}|$ calculated from the example power law (blue) and Gaussian (orange) coherence filters portrayed in Figure 4.4. The lightly shaded blue and orange regions represent the spread (standard deviation) from averaging all 20 simulations.

For both the Gaussian and power law coherent magnetization grid, we implement a Fourier domain modeling technique from Blakely (1995) to calculate the resulting magnetic field at a given location above the distribution. The process is detailed in Appendix 4.8, and summarized below. For a grid with thickness t and depth d from the surface, we divide the slab into horizontal layers. We then forward model the magnetic field at a given point above the surface at altitude H by summing the contributions of each layer. We only apply periodic boundary conditions in the lateral direction and set the vertical thickness t of the grid as a free parameter.

To statistically simulate the magnetization, we find the altitude profiles of magnetic strength above 100 random locations on the grid surface. Each individual altitude profile is scaled relative to the corresponding surface measurement at each lander. This process is repeated for 20 grid simulations, averaging 2000 altitude profiles to obtain a final profile of magnetic intensity above each lander. An example altitude profile is shown in Figure 4.5

using the Gaussian and power law coherence distributions presented in Figure 4.4. The modeled magnetic field at $H = 0$ is normalized to agree with surface measurements near InSight or Zhurong. The values of $|\mathbf{B}|$ at all other altitudes are then linearly shifted by the same amount. Thus, we can vary the grid depth d , thickness t , and vertical/lateral coherence parameters of the grid for each filter to generate characteristic altitude profiles above each lander. The reduced χ^2_ν is calculated between the stochastic models and L19 crustal field estimates to determine the set of parameters that minimizes χ^2_ν .

Geological Representation Models

The method presented in the previous section utilizes a statistical technique to non-uniquely simulate the crustal magnetization on Mars. We now introduce an alternative approach to create magnetization grids near each lander based on Martian stratigraphy and topography rather than random initial conditions. In addition to incorporating the spatial extent of the Martian geology, this technique considers the temporal evolution of the surface. Over time, the Martian surface has been transformed by various geological processes, such as volcanic activity and impact cratering. These mechanisms have created surface features that can be divided into three different time periods (Tanaka 1986). In order from oldest to youngest, the Martian time periods are defined as the Noachian (ending ~ 3.7 Ga ago), Hesperian (ending $\sim 2.9 - 3.3$ Ga ago), and Amazonian (extending to the present day) (Hartmann et al. 2001). During each of these epochs, different NRM processes could have dominated depending on the presence of a global geodynamo.

To correlate surface evolution over time and geological features near each lander, we utilize the 44 geological units determined by Tanaka et al. (2014) through digital photogeologic mapping techniques of multiple spacecraft datasets from the late 1990s. The relative age of each unit is determined by detailed stratigraphic relations from the reconstructed geological map and crater density statistics. The unit grouping is based on geographic settings, lithologic divisions, and epoch ages of Mars. The geographic settings of interest include the highland, lowland, and transitional regions, while the lithologic divisions consist of impact and volcanic units. The regions of interest near each lander encompass the middle Noachian highland (mNh), Hesperian Noachian transition (HNt), Hesperian transition undivided (Htu), early Hesperian transition (eHt), late Hesperian transition (lHt), late Hesperian lowland (lHl), Amazonian Hesperian transition undivided (AHtu), Amazonian Hesperian impact (AHi), Amazonian Hesperian volcanic (AHv), and late Amazonian volcanic (lAv) units. Figure 4.6 illustrates each corresponding unit and organizes them by the geological time scale of Mars.

For each lander, we map the geological units on a 1024×1024 km horizontal grid and assign a given magnetization strength and direction to each region. This magnetization extends downward from an initial elevation d_e to the grid thickness t , where we assume that geological surface characteristics may either be representative of the crust at lower depths, or affected the bulk material below the surface. Magnetization vectors are removed for grid voxels above the measured elevation at each point on the topographic map generated from

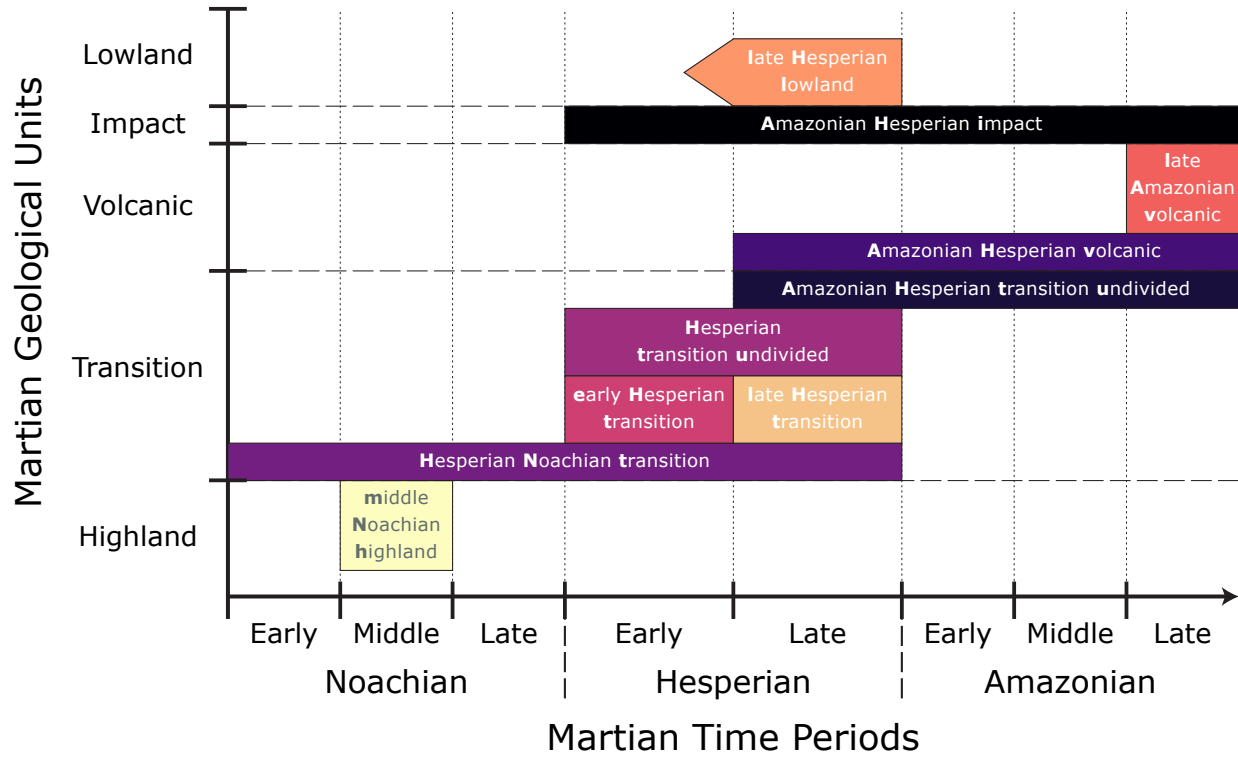


Figure 4.6: Martian geological map units near the InSight and Zhurong landing sites, organized by the relative Noachian, Hesperian, and Amazonian time periods and epochs (not to scale). The bold letters denote each geological unit acronym, and the color scheme matches the unit legend of Figures 4.7a and 4.7c. The triangular tip of the late Hesperian lowland suggests the unit may have originated earlier in Mars’ history. Adapted from Tanaka et al. (2014).

the Mars Orbiter Laser Altimeter (MOLA) instrument onboard the Mars Global Surveyor (Smith et al. 2001). Therefore, it is important to select an appropriate initial grid elevation d_e to accurately represent the full elevation range near each lander. Finally, a lateral convolution filter with a spatial size of 5 km is applied to smooth the magnetization near the region boundaries with opposite field directions. Figure 4.7 demonstrates this process for simulating possible magnetization near InSight and Zhurong based on the local stratigraphy and topography.

The altitude profile of magnetic intensity is determined using the same procedure for the Monte Carlo models described in the previous section. However, only one simulation per unique combination of regional magnetization is modeled due to the specific configuration of the geological units within each grid. While this method does not provide as rigorous a statistical approach compared to the stochastic simulations, we can examine the tendency of

certain geological units to be magnetized based on the combinations of regional magnetization that agree with spacecraft observations of the crustal field. The grid thickness t and the magnetization strength of each region (nine geological units near InSight and four geological units near Zhurong) can be tuned to minimize the reduced χ^2_ν between the model and L19 crustal field values. For simplification, we assume the magnetization depth is relative to the elevation of each grid voxel. This method relies on several underlying assumptions and is based on the premise that surface observations provide insights into subsurface conditions beneath each lander. However, it is important to note that the source of magnetization could be located deep underground, where surface variations would not necessarily provide any indication on magnetized Martian crust.

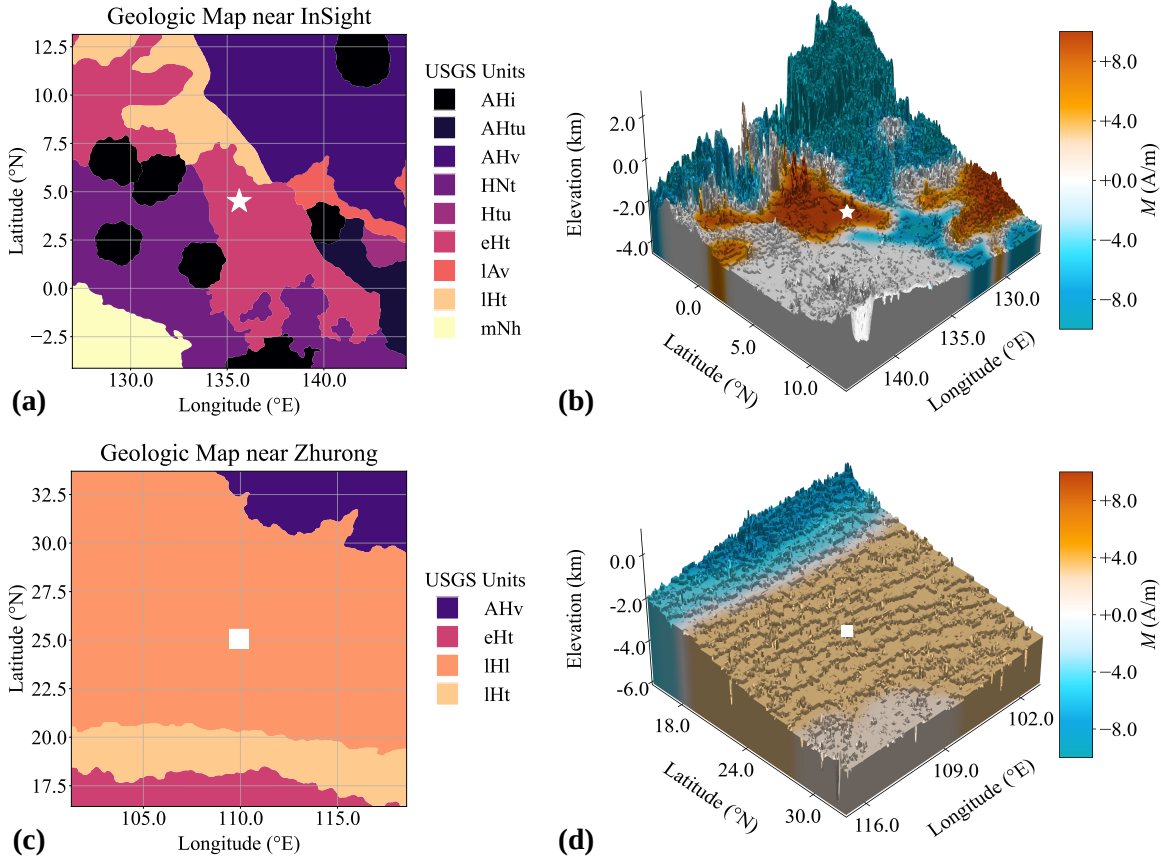


Figure 4.7: The geological map, derived from Tanaka et al. (2014), near (a) InSight (white star) and (c) Zhurong (white square). Each map is 1024×1024 km in horizontal extent, where the geological units are defined in Figure 4.6. An example of the corresponding magnetization grids, which also account for local elevation, are shown for (b) InSight (white star) and (d) Zhurong (white square). The magnetization for each region was selected on a scale of $\pm 10 \text{ A m}^{-1}$ and spatially filtered to smooth the magnetic field near unit boundaries.

4.5 Results

Stochastic Modeling Results: Gaussian Distribution

We conduct a broad search in χ_ν^2 space to estimate the Gaussian model parameters that agree with the lander measurements and L19 crustal field estimates. We clarify that the stochastic Gaussian models are referred to as our simulations and the binned Langlais et al. (2019) crustal field values as L19. For each set of unique parameters, we average the altitude profiles above 100 locations on 20 simulated magnetized grids, characterized by a Gaussian coherence distribution described in Section 4.4. These altitude profiles are generated for a range of simulations with Gaussian lateral and vertical coherence wavelengths from 1 to 1000 km and 1 to 50 km, respectively. In addition, we simulate grid depths between 0 and 30 km and grid thicknesses between 1 and 50 km. Grid simulations with a thickness smaller than the applied vertical coherence scale are excluded from the search in parameter space. It is also important to note that computing the reduced χ_ν^2 in logarithmic space results in similar minima of the grid parameters.

Prior to calculating χ_ν^2 for each set of simulations, we examine the dependence of different coherence scales on the surface magnetization strength M . While M at a single location is not a fitted parameter, it is constrained by the lander measurements. Figures 4.8 and 4.9 show grid magnetization as a function of lateral and vertical Gaussian peak wavelengths near InSight and Zhurong, respectively. We first consider simulations with zero depth from the surface, where the lateral coherence curves are derived from grids with $c_V = 1$ km and the vertical coherence curves are calculated with $c_L = 250$ km. For all simulated grid thicknesses, the magnetization strength is linearly proportional to the lateral coherence wavelength, while inversely proportional to the vertical coherence wavelength. This trend remains consistent when extended to other coherence scales and grid depths. The magnetization near InSight ranges from 1 to ~ 300 A m⁻¹ for $c_V = 1$ km across all lateral coherence wavelengths, while the magnetization near Zhurong is about two orders of magnitude smaller. However, grids with $c_L = 250$ km for all vertical scales are generated with magnetization below 110 A m⁻¹ and 0.6 A m⁻¹ for InSight and Zhurong, respectively.

We then compute 3D histograms of χ_ν^2 and magnetization strength as a function of grid depth and thickness for InSight and Zhurong, shown in Figures 4.10 and 4.11, respectively. For InSight, we examine slices in parameter space of $c_L = [200, 250, 275]$ km and $c_V = 1$ km. As the lateral coherence wavelength increases, the χ_ν^2 minima shift towards shallower grid depths. The surface magnetization is observed to decrease with smaller grid depths and larger grid thicknesses across all coherence scales. For Zhurong, we consider slices in parameter space of $c_L = [800, 900, 1000]$ km and $c_V = 1$ km. As the lateral coherence scale increases, the average value of χ_ν^2 decreases for all grid sizes, with a minimum near $t = 20$ km for large coherence scales. The lack of a smoothly varying χ_ν^2 as a function of grid thickness may be a result of small magnetic field amplitudes with low signal measured both near and above the Zhurong landing site. Similar to InSight, the surface magnetization decreases with larger grids closer to the surface.

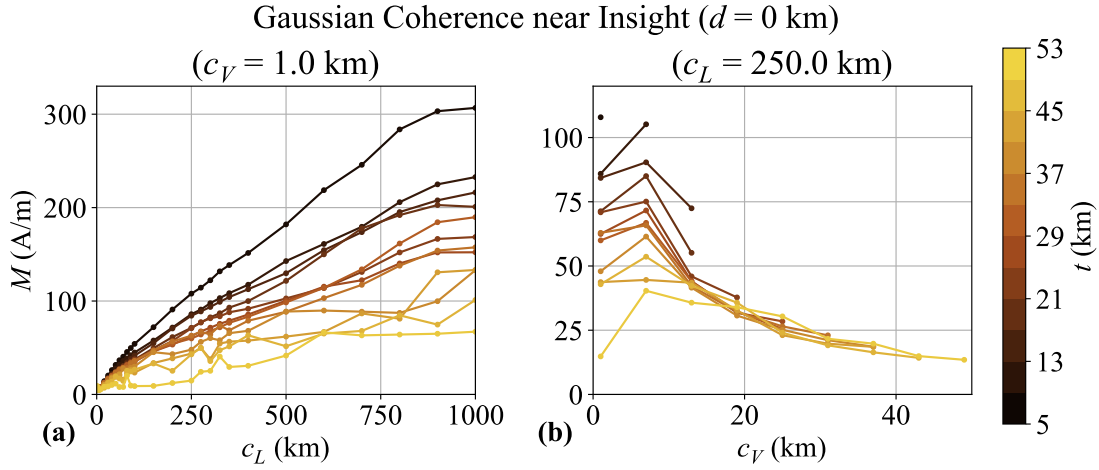


Figure 4.8: Surface magnetization M (A m^{-1}) as a function of **(a)** lateral Gaussian coherence c_L (km) and **(b)** vertical Gaussian coherence c_V (km) for grid simulations with zero depth d below the InSight landing site. The lateral coherence curves are computed with $c_V = 1$ km and the vertical coherence curves are calculated with $c_L = 250$ km. Each curve represents a different grid thickness, excluding simulations where the vertical coherence scale c_V is greater than the grid thickness t .

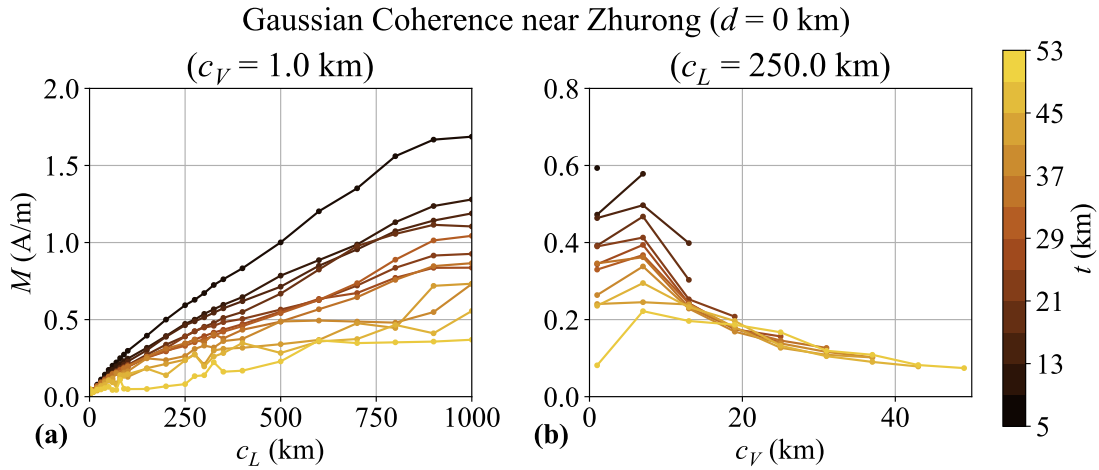


Figure 4.9: Similar to Figure 4.8, but for the Zhurong landing site.

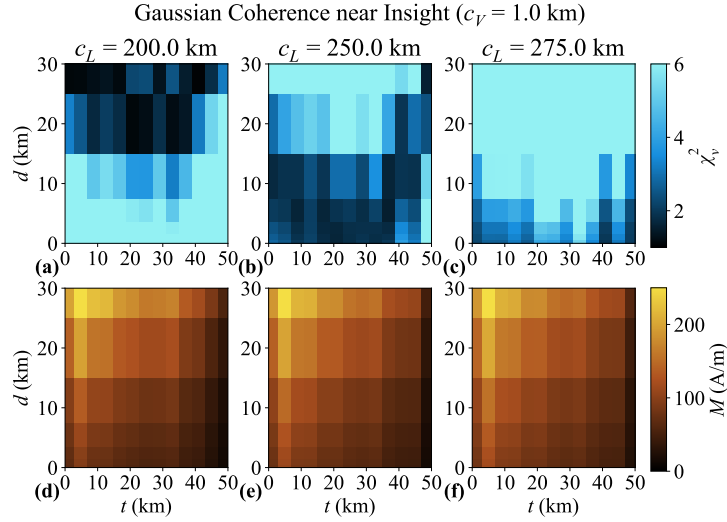


Figure 4.10: 3D χ_v^2 histograms of Gaussian coherence simulations near InSight, arranged by grid depth d (km) as a function of grid thickness t (km), with the lateral coherence c_L set to (a) 200 km, (b) 250 km, and (c) 275 km. 3D surface magnetization M (A m^{-1}) histogram of Gaussian coherence simulations near InSight, arranged by grid depth d (km) as a function of grid thickness t (km), with the lateral coherence c_L set to (d) 200 km, (e) 250 km, and (f) 275 km. In each case, the vertical coherence scale is $c_V = 1$ km.

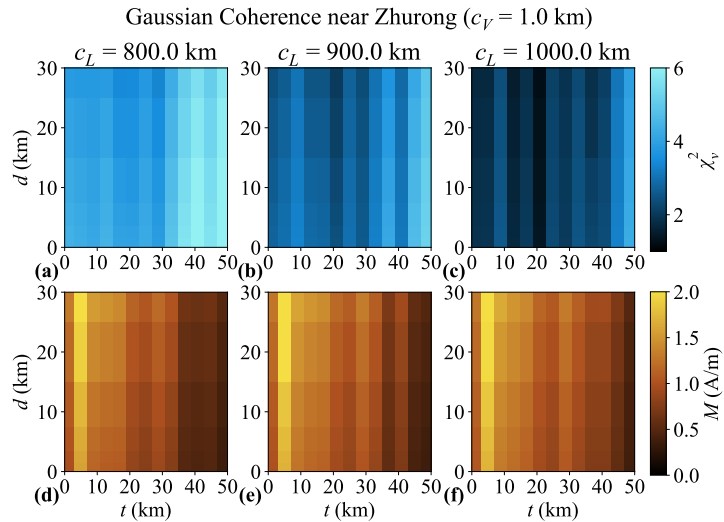


Figure 4.11: Similar to Figure 4.10, but for the Zhurong landing site.

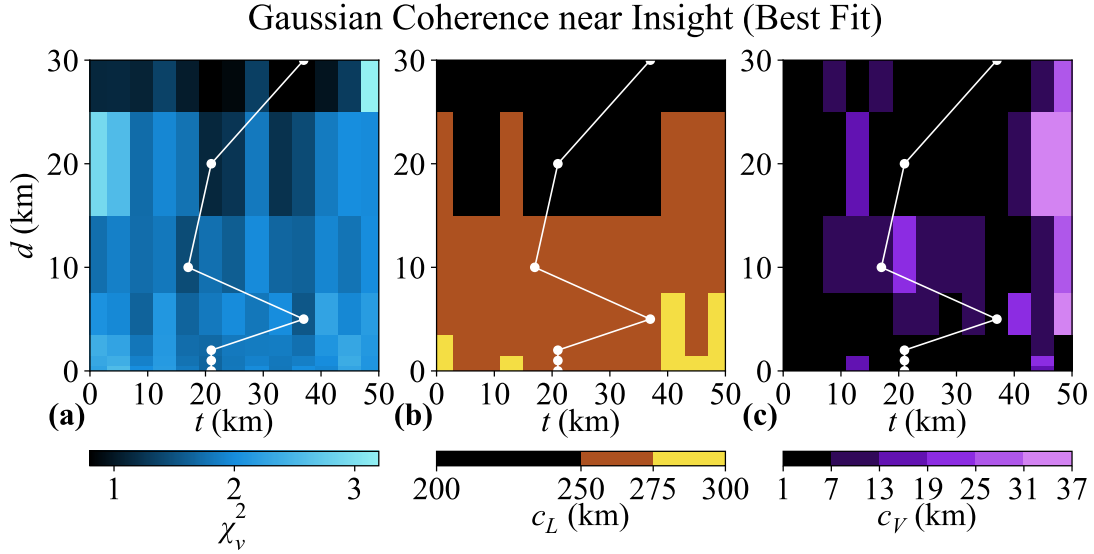


Figure 4.12: 3D histograms of (a) χ_v^2 , (b) lateral coherence c_L (km), and (c) vertical coherence c_V (km) arranged by grid depth d (km) as a function of grid thickness t (km) for the best-fit Gaussian simulations. The white line portrays the location of the absolute minimum value of χ_v^2 for each grid depth.

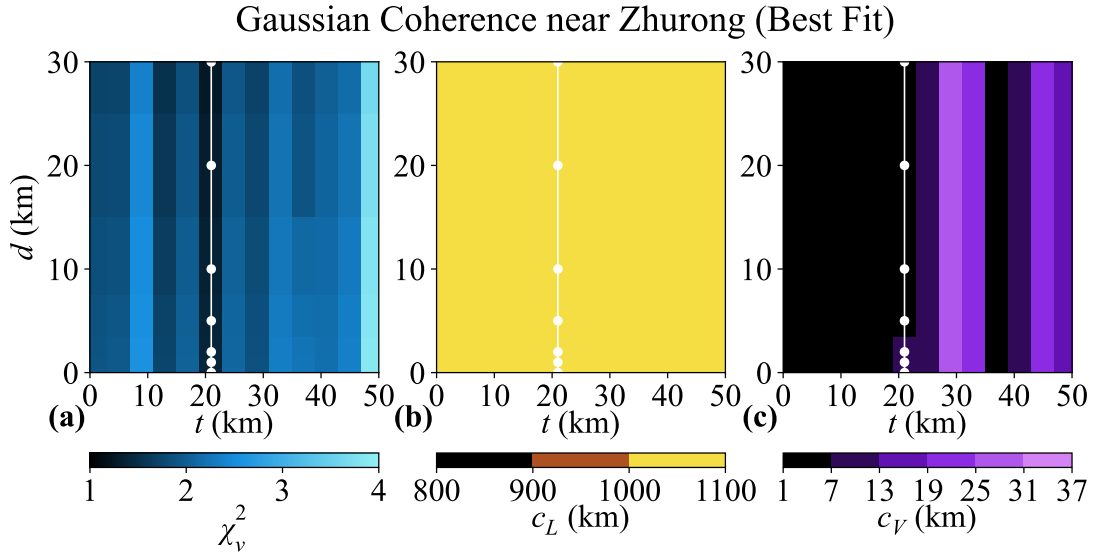


Figure 4.13: Similar to Figure 4.12, but for the Zhurong landing site.

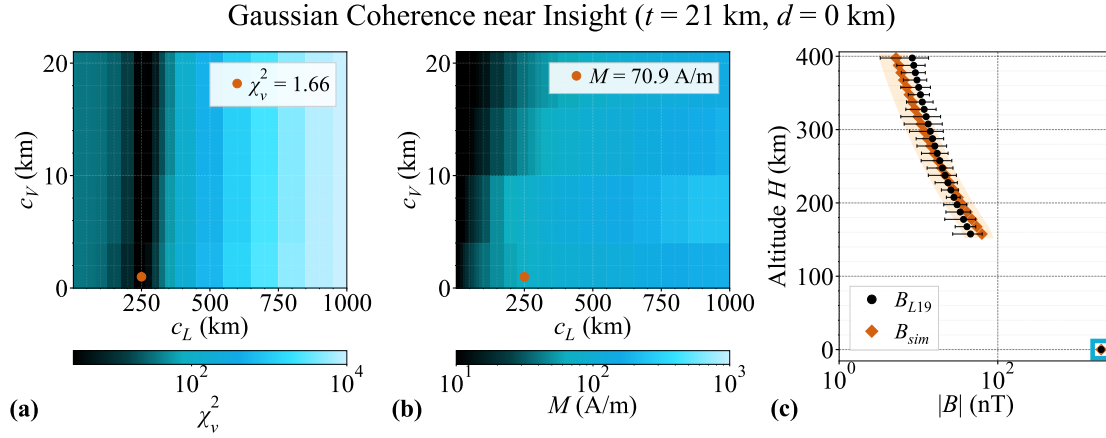


Figure 4.14: Gaussian coherence simulation results near InSight for a grid thickness of 21 km at the surface. **(a)** 3D χ_v^2 histogram as a function of lateral coherence c_L (km) and vertical coherence c_V (km) scales. The minimum χ_v^2 value is shown (orange circle) for $c_L = 250$ km and $c_V = 1$ km. **(b)** 3D surface magnetization M (A m⁻¹) histogram as a function of lateral and vertical coherence scales. The minimum χ_v^2 value (orange circle) corresponds to $M = 70.9$ A m⁻¹. **(c)** Altitude profile of the L19 (black) and simulated (orange) magnetic field strength for $c_L = 250$ km and $c_V = 1$ km. The black error bars represent the standard deviation of binning the L19 model values. The light orange shaded area indicates the standard deviation calculated from averaging all 20 simulations. The blue box at $H = 0$ km highlights the constraint on the magnetic field at the surface.

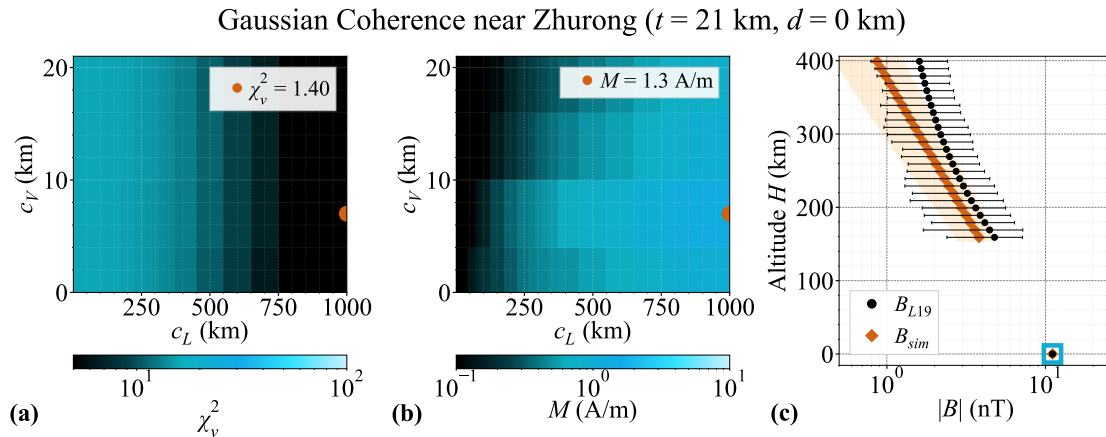


Figure 4.15: Similar to Figure 4.10, but for the Zhurong landing site. The minimum χ_v^2 value (orange circle) is observed for $c_L = 1000$ km, $c_V = 7$ km, and $M = 1.3$ A m⁻¹.

For each simulated grid depth and thickness, we identify the coherence scales that minimize the reduced χ_ν^2 . Figures 4.12 and 4.13 show the minimum χ_ν^2 and corresponding coherence scales for InSight and Zhurong, respectively. For InSight, the values of χ_ν^2 tend to decrease with increasing grid depth and decreasing lateral coherence wavelength. No clear trend is apparent between χ_ν^2 and either grid thickness or vertical coherence wavelength. The selected vertical coherence scale generally increases with larger grid size in the vertical direction. The white vertical lines in Figures 4.12 and 4.13 denote the absolute χ_ν^2 minimum for each grid depth. Across all depths, the absolute minima are observed for $250 < c_L < 275$ km and $c_V < 10$ km. The fitted grid thicknesses tend to cluster either near 20 km or 40 km, but do not follow a clear trend. For Zhurong, there is no evident trend between χ_ν^2 and grid depth, grid thickness, lateral coherence, or vertical coherence. The absolute minimum χ_ν^2 for each depth of the grid occurs for grids with $t \sim 20$ km and $c_L \sim 1000$ km and $c_V < 10$ km.

The results of the Gaussian model for a simulation with a grid thickness of 21 km at the surface are presented in Figures 4.14 and 4.15 for InSight and Zhurong, respectively. The χ_ν^2 and surface magnetization maps are shown as functions of the vertical and lateral coherence scales. The χ_ν^2 minimum (orange circle) for InSight is found for a magnetized grid with $c_L = 250$ km, $c_V = 1$ km and $M = 70.9$ A m⁻¹. For Zhurong, the χ_ν^2 minimum is generated for a magnetized grid with $c_L = 1000$ km, $c_V = 7$ km and $M = 1.3$ A m⁻¹. The selected lateral coherence scale is at the edge of the explored domain, and could be larger than shown in Figure 4.15. The corresponding altitude profile is shown in panel (c) of Figures 4.14 and 4.15. The simulated surface magnetic field magnitude is constrained by the surface measurements from each lander, highlighted by the blue square outline. In both cases, the simulated magnetic field is within 1σ of the L19 profile, with slightly larger deviations at higher altitudes.

Stochastic Modeling Results: Power Law Distribution

A similar search in χ_ν^2 space is performed for the power law coherence simulations based on stochastic modeling. The simulated grid depth and thickness range from 0 to 30 km and from 1 to 50 km, respectively. For these simulations, the lateral coherence power α_L varies from 0.1 to 10 and the vertical coherence power α_V varies from 0.1 to 5.0. Figures 4.16 and 4.17 portray the grid magnetization as a function of lateral and vertical coherence powers near InSight and Zhurong, respectively. The top layer of these simulations is positioned on the surface, where the lateral scale curves are calculated from grids with $\alpha_V = 2.1$ and the vertical scale curves are computed with $\alpha_L = 2.0$. Across all simulated grid sizes, the surface magnetization is non-linearly proportional to the lateral coherence scale, while inversely proportional to the vertical coherence. The magnetization strength ranges from 0.3 to ~ 1000 A m⁻¹ below InSight and from 0.003 to ~ 10 A m⁻¹ near Zhurong for $\alpha_V = 2.1$ across all lateral coherence wavelengths.

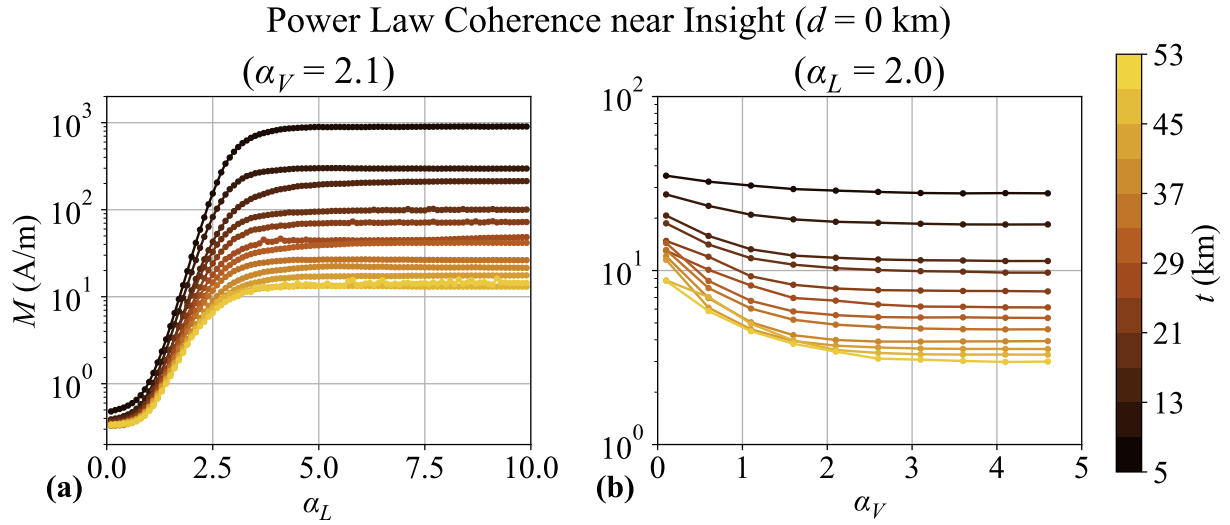


Figure 4.16: Surface magnetization M (A m^{-1}) as a function of (a) lateral power law coherence α_L and (b) vertical power law coherence α_V for grid simulations with zero depth d below the InSight landing site. The lateral coherence curves are computed with $\alpha_V = 2.1$ and the vertical coherence curves are calculated with $\alpha_L = 2.0$. Each curve represents a different grid thickness t (km).

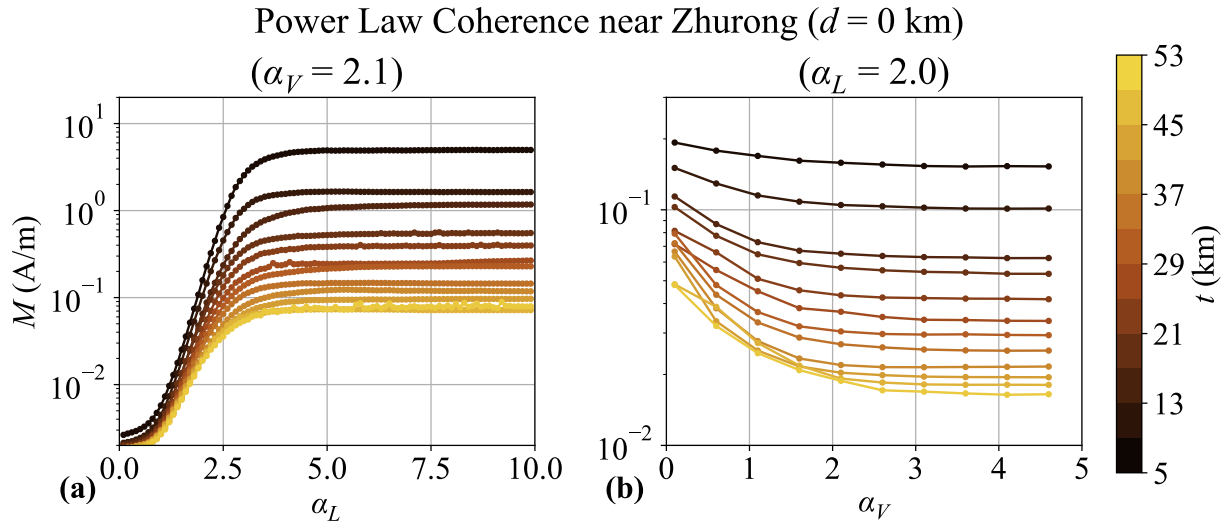


Figure 4.17: Similar to Figure 4.16, but for the Zhurong landing site.

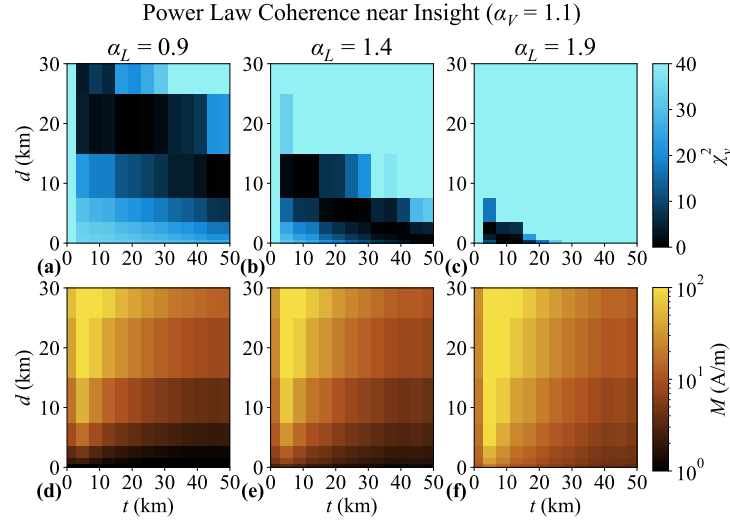


Figure 4.18: 3D χ_v^2 histograms of power law coherence simulations near InSight, arranged by grid depth d (km) as a function of grid thickness t (km), with the lateral coherence α_L set to (a) 0.9, (b) 1.4, and (c) 1.9. 3D surface magnetization M (A m^{-1}) histogram of power law coherence simulations near InSight, arranged by grid depth d as a function of grid thickness t , with the lateral coherence α_L set to (d) 0.9, (e) 1.4, and (f) 1.9. The vertical coherence scale α_V is 1.1.

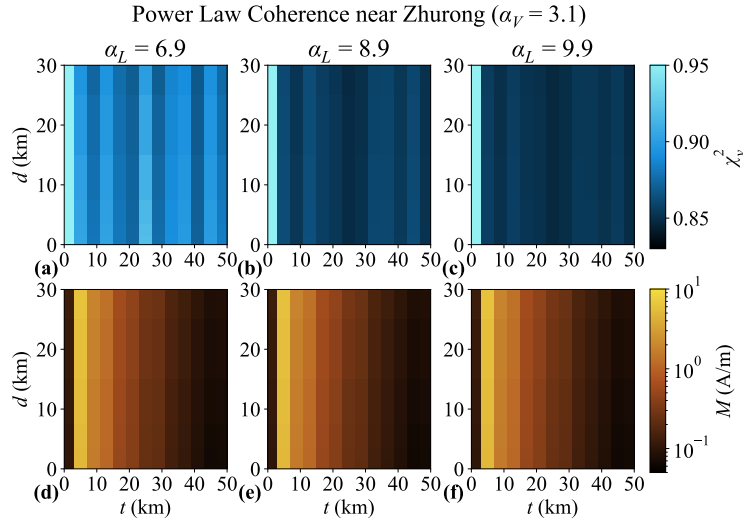


Figure 4.19: Similar to Figure 4.18, but for the Zhurong landing site.

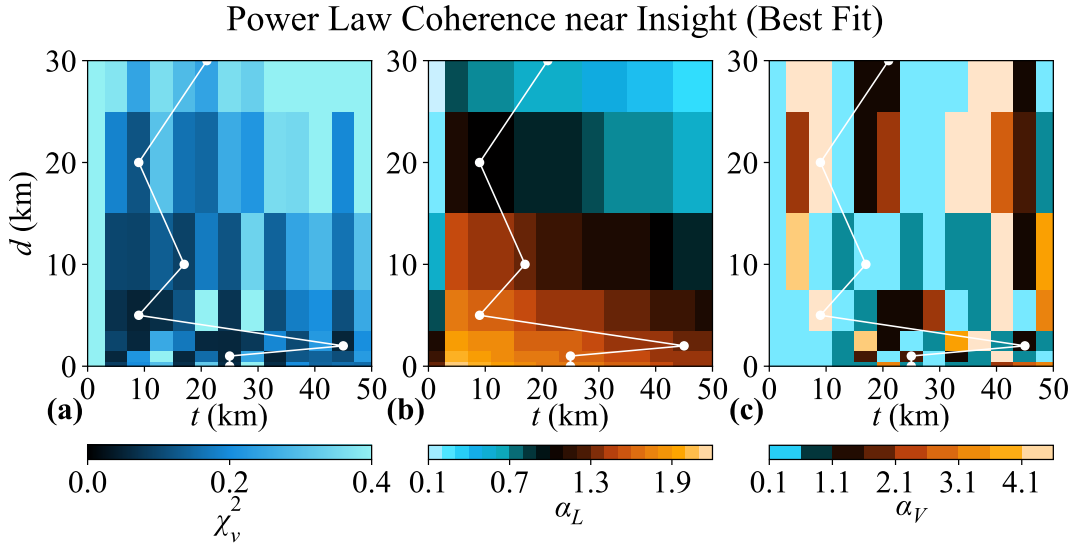


Figure 4.20: 3D histograms of (a) χ_v^2 , (b) lateral coherence α_L , and (c) vertical coherence α_v arranged by grid depth d (km) as a function of grid thickness t (km) for the best-fit power law simulations. The white line portrays the location of the absolute minimum value of χ_v^2 for each grid depth.

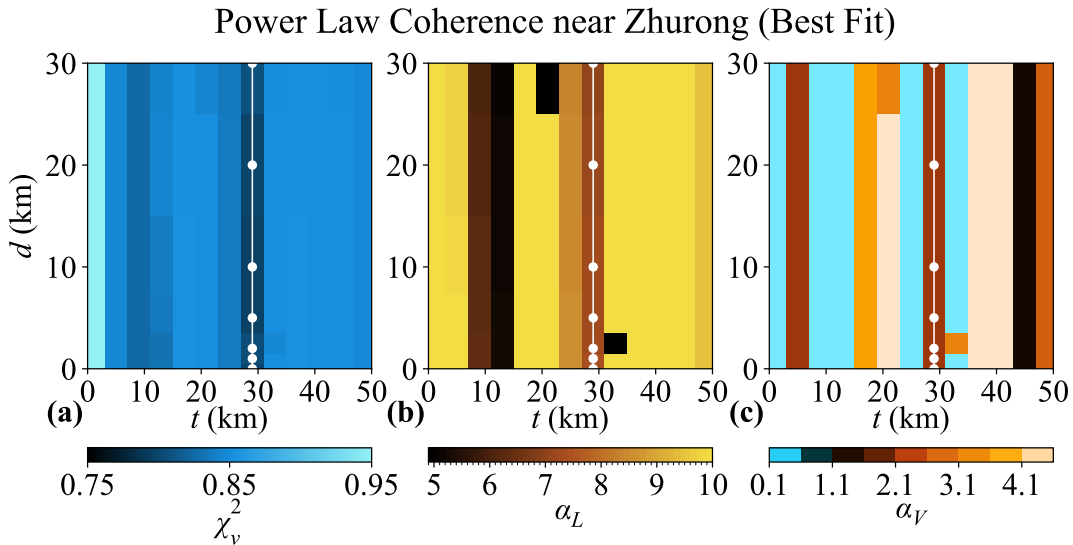


Figure 4.21: Similar to Figure 4.20, but for the Zhurong landing site.

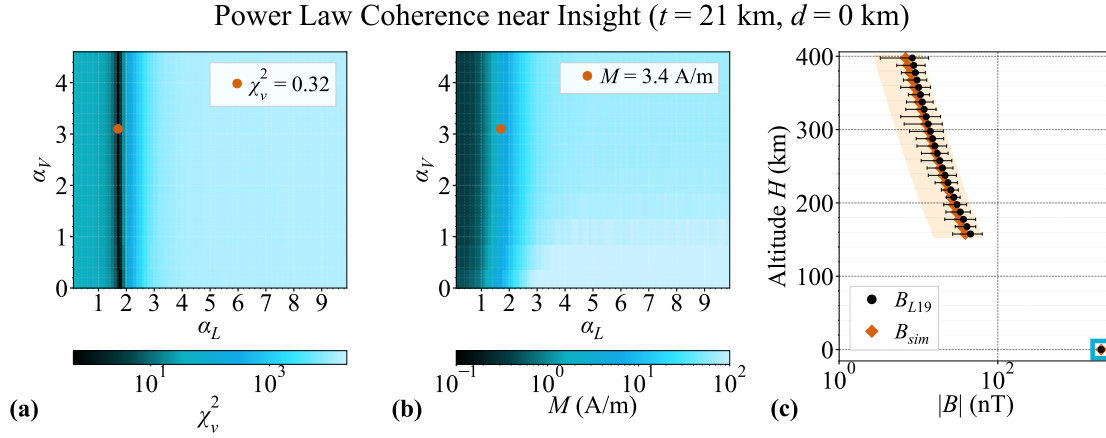


Figure 4.22: Power law coherence simulation results near InSight for a grid thickness of 21 km at the surface. **(a)** 3D χ_v^2 histogram as a function of lateral coherence α_L and vertical coherence α_V scales. The minimum χ_v^2 value is shown (orange circle) for $\alpha_L = 1.7$ and $\alpha_V = 3.1$. **(b)** 3D surface magnetization M (A m⁻¹) histogram as a function of lateral and vertical coherence scales. The minimum χ_v^2 value (orange circle) corresponds to $M = 3.4$ A m⁻¹. **(c)** Altitude profile of the L19 (black) and simulated (orange) magnetic field strength for $\alpha_L = 1.7$ and $\alpha_V = 3.1$. The black error bars represent the standard deviation of binning the L19 model values. The light orange shaded area indicates the standard deviation calculated from averaging all 20 simulations. The blue box at $H = 0$ km highlights the constraint on the magnetic field at the surface.

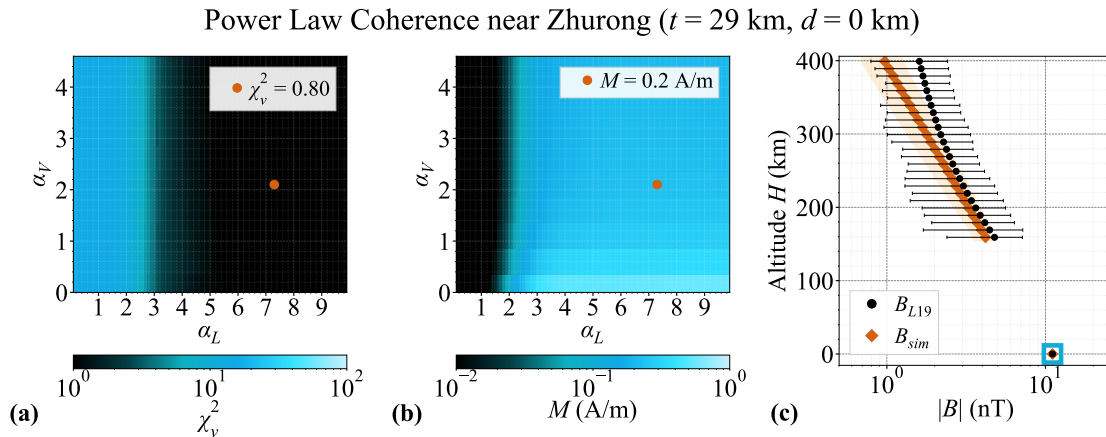


Figure 4.23: Similar to Figure 4.22, but for the Zhurong landing site with a grid thickness of 29 km. The minimum χ_v^2 value is observed for $\alpha_L = 7.3$, $\alpha_V = 2.1$, and $M = 0.2$ A m⁻¹.

Figures 4.18 and 4.19 show the χ_ν^2 and magnetization strength maps as a function of grid depth and thickness for InSight and Zhurong, respectively. For InSight, we consider a subset of coherence powers, focusing on $\alpha_L = [0.9, 1.4, 1.9]$ and $\alpha_V = 1.1$. Similar to the Gaussian model, the χ_ν^2 minima shifts toward smaller grid depths for larger lateral coherence scales. However, as the lateral power increases, the χ_ν^2 minima shifts toward thinner grid sizes. The surface magnetization is found to decrease with smaller grid depths and thicker grids for all coherence scales. For Zhurong, we examine simulations with $\alpha_L = [6.9, 8.9, 9.9]$ and $\alpha_V = 3.1$. As the lateral coherence scale increases, the average value of χ_ν^2 decreases for all grid sizes, with a minimum between 20 and 30 km in grid thickness for large scales. Once again, the sharp variations in the χ_ν^2 map as a function of grid thickness is attributed to small magnetic field measurements near the lander. No significant changes are observed in the values of χ_ν^2 and M when varying the depths of the grid from the surface.

For each grid depth and thickness, we investigate the coherence scales that minimize the reduced χ_ν^2 using a power law distribution. The minimum χ_ν^2 and the corresponding coherence scales are presented in Figures 4.20 and 4.21 for InSight and Zhurong, respectively. The simulations for InSight show smaller χ_ν^2 values for shallower grids and larger lateral coherence powers in general. It is also important to note that the lateral coherence increases for shallower and thinner grids. No relation is observed between χ_ν^2 and the vertical length scale. The values of χ_ν^2 are minimized for grids with $d < 10$ km, $1.3 < \alpha_L < 1.8$ and $\alpha_V < 1.0$. In contrast, the simulation results for Zhurong show no apparent relationship among any of the parameters. However, a definitive χ_ν^2 minimum is found for $t = 29$ km, $\alpha_L = 7.2$, and $\alpha_V = 2.1$.

An example of the simulation results using a power law distribution is shown in Figures 4.22 and 4.23 for InSight ($t = 21$ km) and Zhurong ($t = 29$ km), respectively. The χ_ν^2 and surface magnetization histograms are depicted as a function of vertical and lateral coherence power. The χ_ν^2 minimum (orange circle) for InSight is found for a magnetized grid with $\alpha_L = 1.7$, $\alpha_V = 3.1$ and $M = 3.4$ A m⁻¹. For Zhurong, the minimum χ_ν^2 is generated for a grid with $\alpha_L = 7.3$, $\alpha_V = 2.1$ and $M = 0.2$ A m⁻¹. The altitude profile corresponding to the χ_ν^2 minimum is shown in panel (c) of Figures 4.14 and 4.15. The simulated surface magnetic field magnitude is constrained by the lander surface measurements, highlighted by the blue square outline. The simulated altitude profile above InSight and Zhurong are well within 1σ of the L19 profile above each lander.

Geological Representation Modeling Results

We examine different combinations of magnetization for each neighboring geological region near InSight and Zhurong. Grid thicknesses range from 10 to 40 km in 5 km increments, where a small depth of $d = 200$ m from the surface is assumed for each horizontal position on the grid. For each region, the magnetization is varied from $-M_{max}$ to M_{max} in increments of $M_{max}/2$, where M_{max} is the maximum magnetization that can be assigned to a region, restricted by surface measurements. We vary the magnetization of the mNh, HNt, eHt, lHt, and Htu regions for InSight and the eHt, lHt, and lHl units for Zhurong. All Amazonian

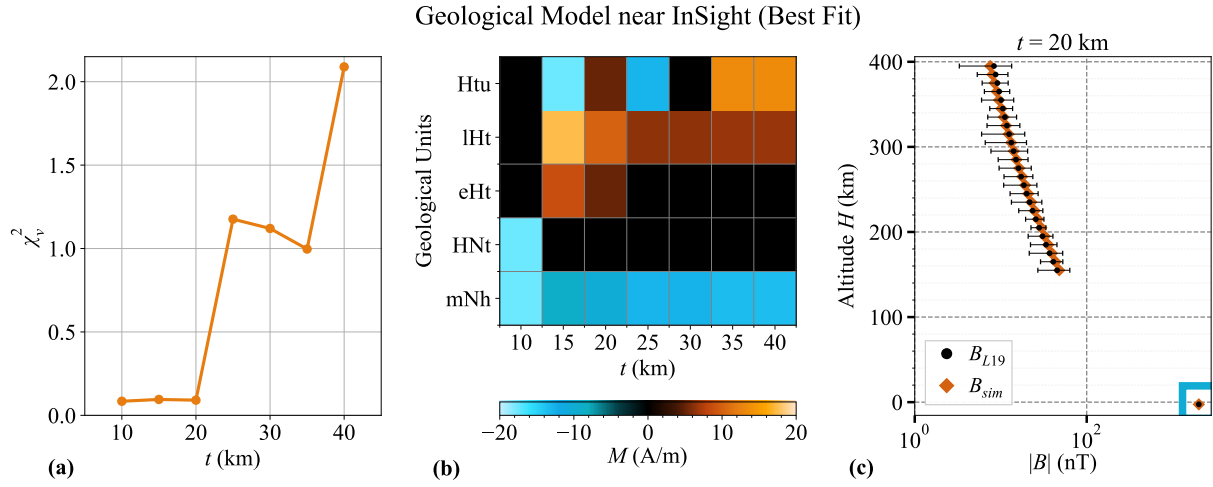


Figure 4.24: Results of the best-fit geological magnetization models near InSight. (a) χ^2_ν minimum for each simulation grid thickness t (km). (b) The corresponding regional magnetization for each geological unit and simulation grid thickness, where positive and negative magnetization refers to parallel and anti-parallel magnetization vectors relative to the local magnetic field direction at each lander. (c) Altitude profile of the L19 (black) magnetic field strength and simulated (orange) magnetic intensity derived from the optimal 20 km thick grid. The black error bars represent the standard deviation of binning L19 model values. The blue box at $H = 0$ km highlights the magnetic field constraint at the surface.

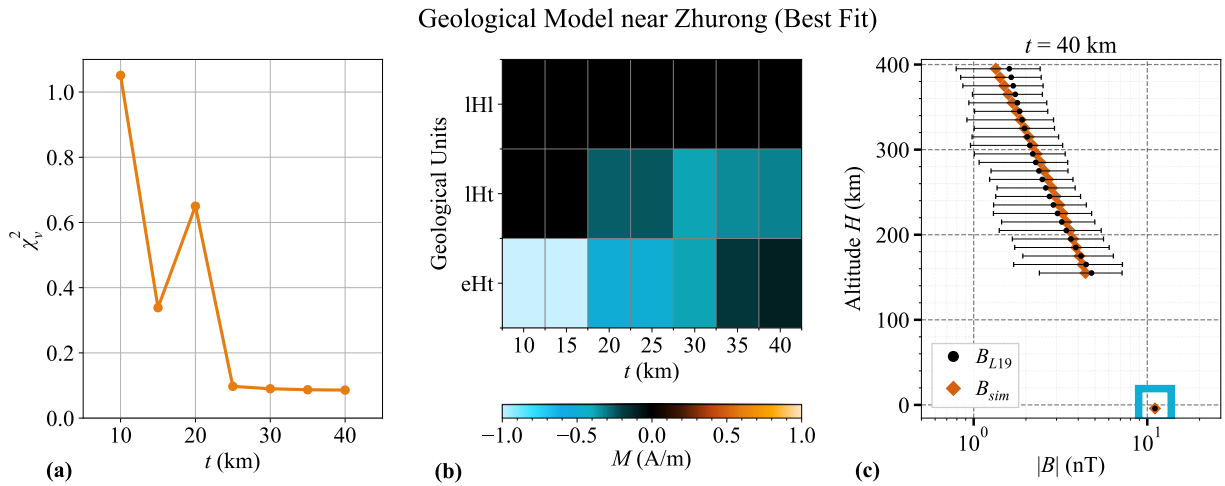


Figure 4.25: Similar to Figure 4.24, but for the Zhurong landing site.

units are assigned $M = 0 \text{ A m}^{-1}$, except for the small AHi regions, which are set to match the magnetization of the underlying region. It is important to note that although all regional combinations of magnetization are simulated, inverting the polarity of the entire grid would yield similar model results.

Figure 4.24 shows the grid parameters that best fit the L19 altitude profile above InSight. The minimum values of χ_ν^2 are found for thin grids with $t < 25 \text{ km}$. Thin grids ($t \lesssim 10 \text{ km}$) exhibit magnetization for most regions with the same polarity. Thicker grids ($t \geq 15 \text{ km}$) are shown to have regions with opposite magnetization between the Noachian and Hesperian periods. For $t = 20 \text{ km}$, the magnetic polarization of mNh is opposite to the direction of the magnetization simulated in the eHt, lHt, and Htu regions. The corresponding altitude profile for the optimal 20 km thick grid simulation is depicted in Figure 4.24c. For all grid thicknesses, the range of maximum magnetization is between 10 and 40 A m^{-1} . The Zhurong simulations are shown in Figure 4.25, where the minimum χ_ν^2 values are found for grids thicker than 25 km. Each simulation demonstrates magnetization with roughly the same polarity in general, with stronger magnetization in lHt region for thicker grids. The maximum magnetization values range from 0.25 to 1.5 A m^{-1} . Figure 4.24c portrays the altitude profile for the 40 km thick grid simulation, where the eHt and lHt units are both magnetized, with little to no magnetization in the lHl region.

As multiple regional combinations resulted in good agreement with the L19 estimates, we analyze all simulations with $\chi_\nu^2 < 2.0$. Figure 4.26 presents the relative frequency of these simulations near InSight for each region with either negative, zero, or positive magnetization. Among the 3125 possible combinations of regional magnetization, 1160, 264, and 56 total simulations with $\chi_\nu^2 < 2.0$ are found for 10, 15, and 20 km thick grids, respectively. Most simulations show negative magnetization for the mNh unit, positive magnetization for the lHt region, and non-zero magnetization for the Htu unit. Non-zero magnetization for the HNt and eHt units is prevalent for thinner grids, but the majority of simulations with thicker grids ($t > 15 \text{ km}$) show no magnetization for these regions. The relative frequency of simulations near Zhurong with $\chi_\nu^2 < 2.0$ is depicted in Figure 4.27. For grid thicknesses of 20, 35, and 40 km, a total of 6, 20, and 20 simulations are found, respectively, from 125 possible regional magnetization combinations. Most simulations across all grid sizes show non-zero magnetization within the eHt region, with little to no magnetization in the lHl unit. Furthermore, the lHt region is magnetized with the same polarity as the eHt geological unit, especially for grid thicknesses near 20 km.

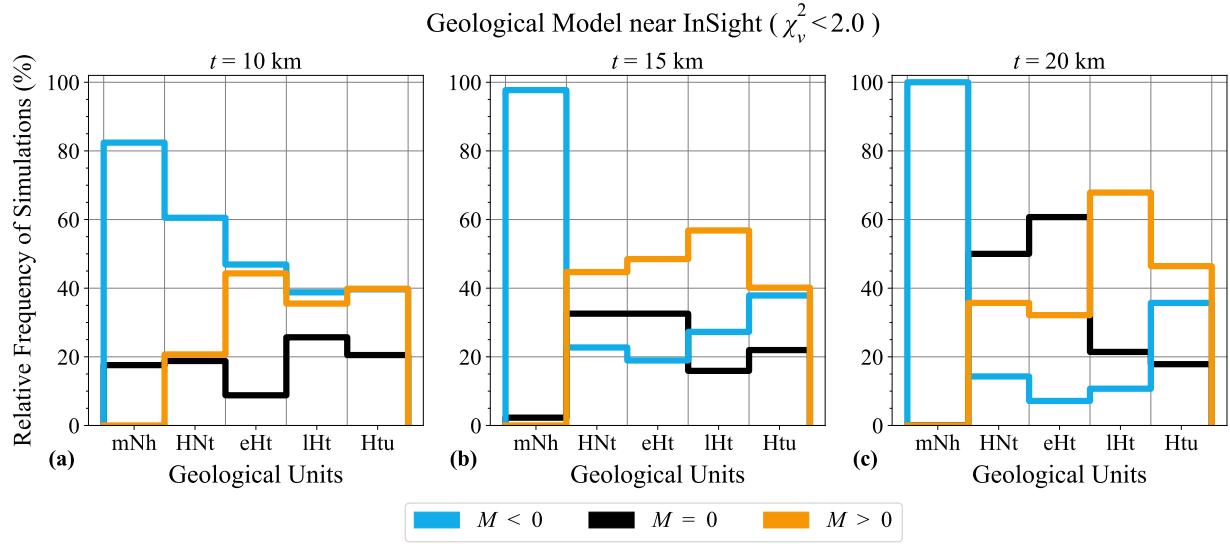


Figure 4.26: Relative frequency of geological models with $\chi_v^2 < 2.0$ for each lithologic region with negative (blue), zero (black), or positive (orange) magnetization for grids with a thickness of (a) 10 km, (b) 15 km, (c) 20 km.

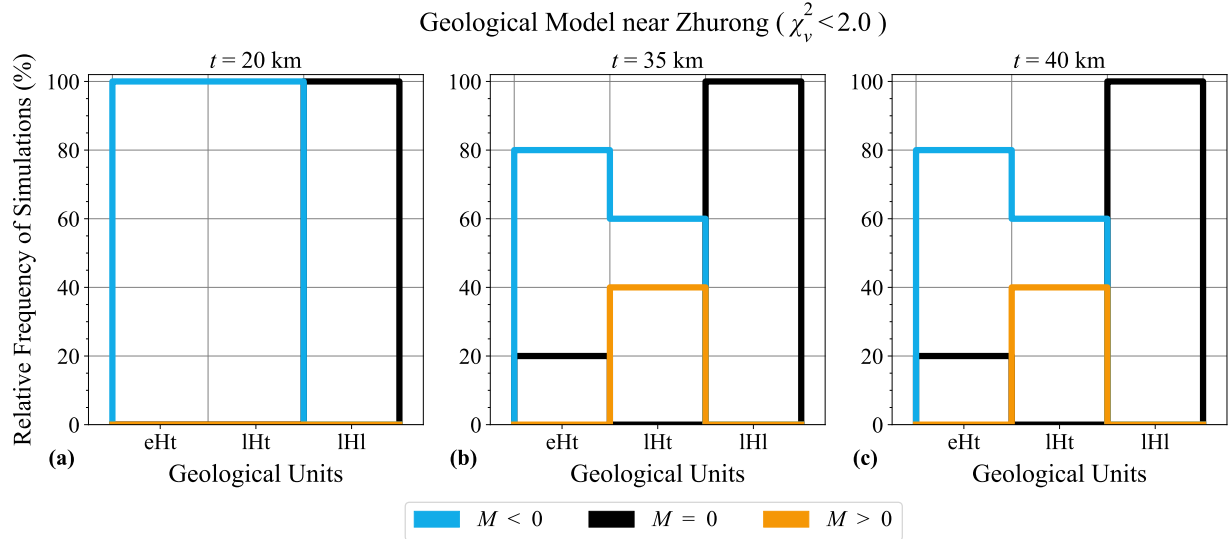


Figure 4.27: Similar to Figure 4.26, but near Zhurong for grids with a thickness of (a) 20 km, (b) 35 km, (c) 40 km.

4.6 Discussion

Our first approach to model Martian crustal fields using stochastic simulations depends primarily on the assumed magnetic coherence distribution. For a Gaussian coherence filter, a lateral coherence scale of $c_L = 250 \pm 50$ km provides the best fit to the altitude profile above InSight, with greater discrepancy at the lower and upper limits of the MAVEN altitude range. The vertical coherence is found to predominantly influence the magnetization strength M , resulting in a range from ~ 15 to 115 A m^{-1} depending on the grid thickness for $d = 0$ km and $c_L = 250$ km. For thinner grids 10 km below the surface, the magnetization strength can reach above 200 A m^{-1} . Ravat (2011) has shown that magnetization values up to 50 A m^{-1} may be possible on Mars due to stronger magnetizing fields on the surface and higher quantities of magnetic iron oxide in the crust compared to Earth. However, it is difficult to account for magnetization above 40 A m^{-1} based on the possible types of iron-bearing crustal minerals on Mars (Dunlop et al. 2005; Nimmo et al. 2001). Thus, simulations requiring intense magnetization strengths are not realistic for the Martian crust. Regardless of magnetization strength, lateral scales between 200 and 300 km consistently match the InSight altitude profile for any grid thickness, depth, or vertical coherence scale. Therefore, a thicker grid ($t \gtrsim 20$ km) closer to the surface with $c_L \approx 250$ km and larger vertical coherence may be more realistic to achieve magnetization strengths less than 40 A m^{-1} . Lillis et al. (2009) reported similar magnetic lateral coherence scales near the southwestern Tharsis region, linked to volcanic activity, using a grid thickness of 30 km. These coherence scales are more than a factor of two smaller than the magnetic length scales identified by Voorhies (2008) to model the global magnetic field distribution of Mars (~ 650 km). Lillis et al. (2010) attributed smaller coherence scales to younger crust which may have developed over a prolonged period of time later in Martian history compared to the rest of the planet.

In contrast to InSight, simulated altitude profiles near Zhurong using a Gaussian filter with $c_L = 1000 \pm 100$ km achieves the best fit compared to L19. The corresponding magnetization values are also more reasonable than the InSight simulations, ranging from 0.75 to 1.75 A m^{-1} for $c_L = 1000$ km and $c_V = 1$ km. We find a clear χ^2_ν minimum in Figure 4.13 for $t = 20$ km, associated with magnetization strengths between 0.9 and 1.3 A m^{-1} depending on the grid depth. Larger coherence scales tend to result in better agreement with the L19 altitude profile, but still consistently underestimate the L19 crustal field estimates above Zhurong. It is possible that the surface measurement calculation near Zhurong is closer to zero, indicating no magnetization near the lander and within that region of the Utopia Basin. To investigate broader length scales, additional simulations are modeled separately (not shown) with lateral coherence scales ranging from 1000 to 2000 km. However, Gaussian coherence filters with $c_L \sim 2000$ km only provide a slight improvement compared to the simulated altitude profile in Figure 4.15. These large coherence wavelengths may be reflected by the demagnetization diameter (~ 3100 km) for the Utopia impact basin determined by Lillis et al. (2010) using a similar Monte Carlo method. This may indicate that the magnetization (or lack thereof) is dominated by Utopia basin length scales above Zhurong. Due to the absence of a demagnetization basin in our models, our simulation results cannot be

directly compared to the lateral coherence $c_L = 788$ km observed by Lillis et al. (2010) for Utopia. Furthermore, the simulated grid in this study is centered near the Zhurong landing site instead of the basin and does not include the full extent of Utopia. In comparison to the magnetic coherence scale of 1 km determined by Du et al. (2023) near Zhurong based on surface measurements, our defined grid resolution may be too large to capture length scales this small. Future work could implement finer grid simulations with a demagnetized basin using a Gaussian distribution to directly compare with previous studies.

We find better agreement between our model and L19 using the power law coherence filter compared to the Gaussian distribution for both landers. For $d < 10$ km near InSight, the optimal lateral scale powers α_L are between 1.5 and 1.8. Similar to the Gaussian coherence filter, thicker grids ($t \gtrsim 20$ km) closer to the surface may be more realistic to the magnetization strengths possible on Mars. Although power law distributions do not determine a single peak wavelength, powers closer to 1 indicate significant contributions from finer structures. On Earth, Bouligand et al. (2009) compared the fractal parameter β , defined as the slope of the 3D magnetization power spectrum, from numerous terrestrial studies. In order to directly compare our estimates of power α_L , we define $\beta = \alpha_L + 1$, where α_L is computed in the lateral direction for a 2D power spectrum. Therefore, we determine the fractal parameter β to range from 2.5 to 2.8 near InSight on Mars. Terrestrial fractal parameters were found between 2.5 and 5.0 across different lithologies based on maps of magnetic susceptibility measurements. Fractal powers near 4.0 have been correlated with deep crustal magnetization in continental domains, while smaller fractal powers less than 3.0 have been linked to volcanic rocks in western North America (Bouligand et al. 2009) and volcanic regions of the Japanese Islands (Okubo et al. 2023). Furthermore, Salem et al. (2014) reported lower fractal parameters between 2.4 and 2.7 offshore the Central Red Sea, an area of divergence between two continental sub-plates. Thus, the fractal power near InSight on Mars is comparable with magnetic distributions on Earth near volcanic regions and divergent plate boundaries. However, it has been argued that constant fractal parameters cannot accurately model extended regions with different geological terranes and rock types (Bansal et al. 2010; Gettings 2005; Maus et al. 1995; Pecknold et al. 2001). Modeling entire magnetic regions with a single value of β could combine distributions of various lithologies, neglecting the possibility of β to vary with wavenumber.

Simulated Zhurong altitude profiles using a power law filter show a distinct χ_v^2 minimum for a grid thickness of 30 km across all depths from the surface. Most simulations in Figure 4.21 result in large fractal powers, but 30 km thick grids exhibit lateral coherence of $\alpha_L = 7.3$ and vertical coherence of $\alpha_V = 2.1$, with a magnetization strength near 0.2 A m^{-1} . The vertical coherence power is roughly constant as a function of depth, with no clear relationship with the thickness of the grid. The power law filter demonstrates better agreement with the L19 altitude profile compared to the Gaussian distribution, but continues to deviate from the L19 estimates at higher altitudes. Converting the 2D fractal power α_L to the fractal parameter β , we find a new lateral coherence near 8.3, suggesting a suppression of finer magnetic scales. Fractal powers greater than 8.0 are larger than the highest values of β observed at Earth. For example, high values of β have been modeled near the British Isles

comprising two different geological terranes, resulting in a range of fractal powers between 6.0 and 7.0 (Mather et al. 2019). The large fractal parameter observed below Zhurong may be caused by the length scales of the Utopia basin, which encompasses roughly 70% of our 1024×1024 km magnetization grids.

As previously discussed, the magnetic coherence distribution can depend on both rock type and geological features in the lithosphere. Notably, the ~ 250 km coherence observed near InSight using a Gaussian filter roughly corresponds to the size of some geological units in a 10^6 km area around the lander. Similarly, the magnetic coherence scale exceeding 1000 km near Zhurong agrees to first order with the scale size of the IHI unit within Utopia. This motivates our study to examine different magnetization distributions for each geological unit on the Martian surface. We consider a unique distribution of uniform magnetization in each region solely based on topography, smoothing magnetization boundaries between units. In addition, we investigate cases of opposite magnetization, originating from possible paleomagnetic reversals throughout Martian history (Arkani-Hamed 2001; Hood et al. 2005). We assume (1) a constant coherence scale within each region and (2) magnetization depths solely based on the local surface elevation. Our study provides a simple attempt to directly relate magnetization length scales to geological time scales on Mars, where future work could implement more complex magnetic distributions for each geological unit. However, it is important to note that the magnetized source below the surface could be deep underground, in which case the surface features we observe on Mars may have no connection with the distribution of crustal magnetization.

The regional magnetization simulations modeling the geology and topography near InSight result in a better fit compared to either stochastic model. Despite this improvement, both our stochastic and geological simulations suggest a magnetized crustal thickness of roughly 20 km below InSight (40 km grids also are shown to match the L19 profile using stochastic modeling). After accounting for the local elevation at the lander, this would indicate a magnetization thickness of roughly 17 km below InSight. Our estimate of magnetic crustal thickness is similar to the two layer crustal model with a thickness of 20 ± 5 km reported by Knapmeyer-Endrun et al. (2021), assuming most of the crust is magnetized. However, it is important to note that this two layer model is no longer probable based on analysis from InSight seismic data (Kim et al. 2021), and it more likely that a three layer model exists with an average crustal thickness greater than 30 km (Wieczorek et al. 2022). Our estimate only provides a lower bound on the crustal thickness, as it is possible that not all of the crust is magnetized. Thicker grids of 40 km could better explain the crustal field altitude profile, allowing smaller values of grid magnetization. Based solely on our geological model, Figure 4.28a shows one example of an optimal regional magnetization configuration near InSight for a grid thickness of 20 km. To compare with the coherence distributions implemented in the stochastic simulations, we calculate the lateral power spectrum of grid magnetization for the geological model, portrayed in panel (c). The spectrum predominantly scales as a power law with $\beta \sim 2.8$, where the coherence of smaller wavelengths is governed by the window averaging size to smooth the magnetization between region boundaries. In addition, small peaks in the power spectrum are observed at coherence wavelengths of 125,

300, and 600 km, roughly corresponding to the scale size of nearby regions. This suggests that accurately modeling the magnetization below InSight may require a combination of multi-peak Gaussian and power law distributions.

Most of the best-fit simulation grids thicker than 15 km consistently show the Noachian region to have opposite magnetization compared to the Hesperian regions, especially the IHT unit. This change in magnetization direction may indicate at least one magnetic field polar reversal between the Noachian and Hesperian time periods (Milbury et al. 2010, 2012). Furthermore, geological units aged near the transition between the Noachian and early Hesperian tend to have little to no magnetization, possibly signifying a weakened or suspended geodynamo between these two periods. However, a magnetized late Hesperian unit may suggest a longer-lived dynamo (Mittelholz et al. 2020b; Steele et al. 2024, 2023), or that the dynamo reactivated and magnetized the crustal material through NRM processes possibly derived from volcanic activity (Lillis et al. 2006) or serpentinization of ultramafic rocks (Quesnel et al. 2009) to create magnetically susceptible materials. Previous studies (e.g., Lillis et al. 2006; Zhang et al. 2023) have proposed a similar two-period dynamo model, where large impacts may have played a key role in its evolution. Roberts et al. (2009) has suggested that the Martian dynamo could have ceased around the mid-Noachian due to a significant reduction in core heat flow caused by large impact heating. Conversely, Arkani-Hamed et al. (2011) has argued that large impacts could heat the interior of Mars via shock waves, potentially stratifying the core, suppressing thermal convection, and suspending a core dynamo. For Mars, it could take tens to hundreds of millions of years to resume global internal convection and restart a new geodynamo (Arkani-Hamed 2011; Arkani-Hamed et al. 2011).

The geological representation models near Zhurong also resulted in better agreement with the L19 crustal field profile above the lander compared to the Monte Carlo simulations. While magnetized grids with $t = 20$ km are consistent with the magnetization source thickness near InSight and the regional depth of a second crustal layer near the dichotomy boundary (Li et al. 2022), 40 km thick magnetized grids also result in a comparable fit with the L19 profile. This thicker grid is reasonable based on the estimate of the third crustal interface by Knapmeyer-Endrun et al. (2021) of 39 ± 8 km, assuming similar crustal depths for regions near the dichotomy boundary. For all grid sizes, the majority of optimal simulations show little to zero magnetization in the lowland unit, with greater magnetization south of Zhurong at higher elevations, illustrated in Figure 4.28b. The magnetic power spectrum near Zhurong in panel (c) primarily follows a power law distribution, dominated by the large scales of the Utopia basin. However, a small spectral peak near 200 km may be related to the combined latitudinal size of the early and late Hesperian transition units. This magnetic configuration suggests that the Utopia basin was demagnetized from a large impact due to the lack of a core dynamo. Other surface NRM processes could have created remanent magnetization on small local length scales, observed by Du et al. (2023), later in Martian history if a geodynamo was reactivated.

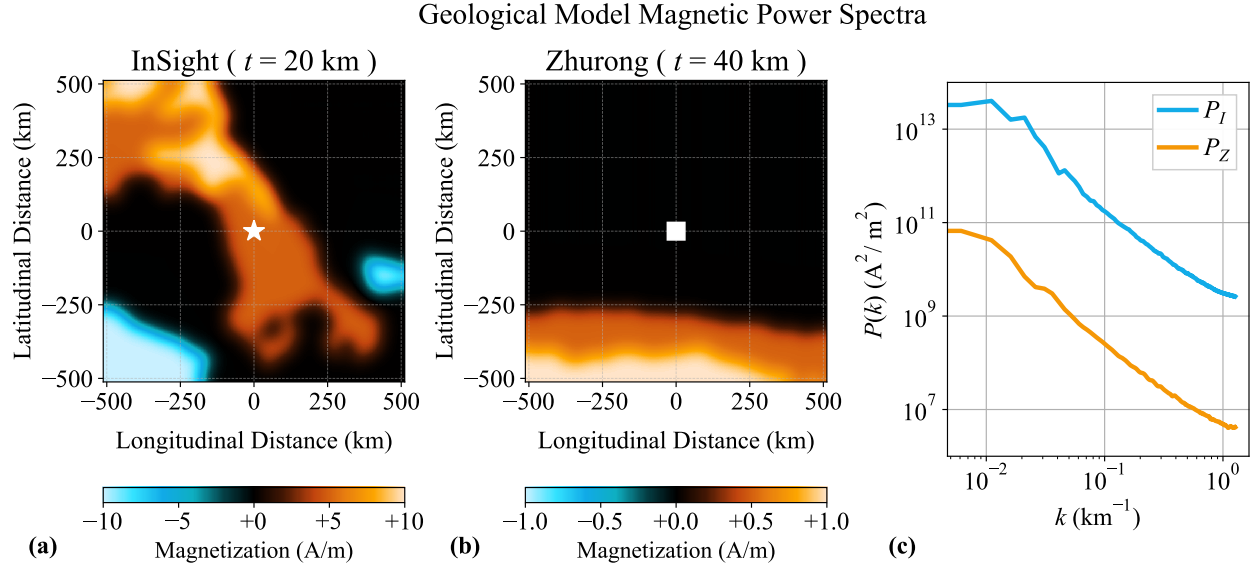


Figure 4.28: Regional magnetic configuration near (a) InSight (white star) and (b) Zhurong (white square) for simulated grids with 20 and 40 km thickness, respectively. (c) Lateral power spectra of the simulated magnetic grids near InSight (blue) and Zhurong (orange) as a function of wavenumber k (km⁻¹).

4.7 Conclusion

Our study provides new insights into the Martian crustal field by comparing stochastic and geological simulations of magnetization with the first in situ magnetic field measurements by InSight and Zhurong. For InSight, Gaussian coherence filters with lateral coherence scales around 250 km are consistent with L19 altitude profiles of magnetic intensity above the lander. While power law distributions demonstrate stronger agreement with the L19 model, both stochastic models suggest a magnetization source thickness of roughly 20 km or 40 km close to the surface. Contrary to the magnetization near InSight, larger coherence scales of 1000 km are observed near Zhurong, primarily dominated by the lack of strong magnetization within the Utopia impact basin. The estimated range of crustal magnetization thickness surrounding Zhurong varies from roughly 20 to 30 km, depending on the type of stochastic model.

Motivated by the length scales of coherent magnetization, we compute geologically and topographically representative simulations near each lander. We determine a similar estimate of crustal thickness below InSight, with the possibility of thicker magnetization sources near Zhurong. By relating geological regions to different time periods on Mars, we are able to determine the evolution of the Martian geodynamo. We find opposite polarity of magnetization in different regions, suggesting at least one polar reversal of the global magnetic

field between the Noachian and Hesperian periods. Furthermore, the absence of simulated magnetization in geological units near the Noachian-Hesperian transition could indicate a temporary suspension of the Martian core dynamo. This suggests that the geodynamo was active at some point during the Noachian period, weakened or ceased between the Noachian and Hesperian eras, and restarted during the Hesperian period.

These findings emphasize the importance of considering different geological and topographic features when modeling local Martian magnetization and suggest that a combination of multi-peak Gaussian and power law distributions may be necessary to simulate local Martian crustal fields. Future work could further refine these models by incorporating more complex distributions within each geological region near InSight, and re-exploring the effects of demagnetization basins on the Martian crustal field near Zhurong. In addition, finer grid resolution could be implemented in future simulations to study smaller coherence scales observed by these novel in situ measurements of the surface magnetic field.

4.8 Appendix: Derivation of Fourier Domain Modeling

To compute the resulting magnetic field vector at a given altitude above a 3D magnetized grid, we derive the magnetic field $\mathbf{B}$ based on the Fourier domain modeling from Blakely (1995). We start with the Fredholm equation of the first kind to describe a potential field $f(P)$ created from a source distribution s at an observation point P

$$f(P) = \int_R s(Q)\psi(P, Q)dV \quad (4.4)$$

where $s(Q)$ represents the source value at a given point Q in the distribution and $\psi(P, Q)$ is the Green's function that depends on the geometric location of P and Q . The integral is evaluated over the entire region R occupied by the source distribution for a volume element dV . Equation 4.4 can be applied to forward model the magnetic field $B(P)$, where $s(Q)$ is the magnetization distribution M and $\psi(P, Q)$ is the magnetic field from a single dipole, described below.

Consider a 3D slab of magnetization, shown in Figure 4.29, where x points north, y is directed east, and the z axis is pointed downward. To calculate the magnetic field $B(x, y, z_0)$ at an altitude z_0 above the surface, we divide the magnetization source grid into N horizontal layers. The first layer of the entire 3D slab, extending from z_1 to z_2 , is defined to have a magnetization distribution M_1 . Each layer has an associated magnetization distribution M_j , which is only a function of x' and y' . We require $z_0 > z_j$ and $z_{j+1} > z_j$ for all x' and y' , where $j = 1, 2, 3 \dots N$ is the layer index. From equation 4.4, the observed magnetic field $B_j(x, y, z_0)$ computed solely from the j th layer is represented as

$$\begin{aligned}
B_j(x, y, z_0) &= \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} \int_{z_j}^{z_{j+1}} M_j(x', y') \psi_j(x - x', y - y', z_0 - z') dx' dy' dz' \\
&= \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} M_j(x', y') \int_{z_j}^{z_{j+1}} \psi_j(x - x', y - y', z_0 - z') dz' dx' dy'
\end{aligned} \tag{4.5}$$

where ψ_j is the field at (x, y, z_0) due to a point source below the surface at (x', y', z') . The next step is to apply a Fourier transform on both sides of equation 4.5. In this context, the Fourier transform $\mathcal{F}[f]$ maps a function $f(x, y)$ from the space domain into the wavenumber domain, defined as

$$\mathcal{F}[f] = F(k_x, k_y) = \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} f(x, y) e^{-i(k_x x + k_y y)} dx dy \tag{4.6}$$

Applying the Fourier transform on each side, equation 4.5 becomes

$$\begin{aligned}
\mathcal{F}[B_j(x, y, z)] &= \mathcal{F} \left[\int_{-\infty}^{\infty} \int_{-\infty}^{\infty} M_j(x', y') \int_{z_j}^{z_{j+1}} \psi_j(x - x', y - y', z_0 - z') dz' dx' dy' \right] \\
&= \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} M_j(x', y') \int_{z_j}^{z_{j+1}} \mathcal{F}[\psi_j(x - x', y - y', z_0 - z') dz' dx' dy']
\end{aligned} \tag{4.7}$$

where we separated the Fourier operation on the second line using the Fourier-convolution theorem for linear systems, given by

$$\mathcal{F}[B_j] = \mathcal{F}[M_j] \mathcal{F}[\psi_j] \tag{4.8}$$

It is also important to note that the Fourier transform is only applied with respect to x and y , and not with respect to x' and y' . To remove the Green's function dependence on the field point such that $\psi = \psi(x, y)$, we employ the shifting property of Fourier transforms, defined as

$$\mathcal{F}[\psi(x - x', y - y')] = \mathcal{F}[\psi(x, y)] e^{-i(k_x x' + k_y y')} \tag{4.9}$$

We can rewrite equation 4.7 to separate the shift of ψ in the lateral dimensions as

$$\mathcal{F}[B_j(x, y, z_0)] = \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} M_j(x', y') \int_{z_j}^{z_{j+1}} \mathcal{F}[\psi_j(x, y, z_0 - z')] e^{-i(k_x x' + k_y y')} dz' dx' dy' \tag{4.10}$$

To find the Fourier transform of $\psi(x, y, z)$, we equate the function to a magnetic field $\mathbf{B}_D$ from a single dipole. We start with $|\mathbf{B}_D| = |-\nabla_p V|$, where the magnetic potential V of a dipole with a unit moment vector $\hat{\mathbf{m}}$ is given by

$$\begin{aligned}
V &= -C_m \hat{\mathbf{m}} \cdot \nabla_p \frac{1}{r} \\
&= -C_m \left(\hat{m}_x \frac{\partial}{\partial x} \frac{1}{r} + \hat{m}_y \frac{\partial}{\partial y} \frac{1}{r} + \hat{m}_z \frac{\partial}{\partial z} \frac{1}{r} \right)
\end{aligned} \tag{4.11}$$

where C_m is the proportionality constant to convert between units and is defined as $C_m \equiv 10^{-7} \text{ H m}^{-1}$. Applying the Fourier transform on the potential in the horizontal plane yields

$$\begin{aligned}\mathcal{F}[V] &= -C_m \left(\hat{m}_x \mathcal{F} \left[\frac{\partial}{\partial x} \frac{1}{r} \right] + \hat{m}_y \mathcal{F} \left[\frac{\partial}{\partial y} \frac{1}{r} \right] + \hat{m}_z \mathcal{F} \left[\frac{\partial}{\partial z} \frac{1}{r} \right] \right) \\ &= -C_m \left(\hat{m}_x i k_x \mathcal{F} \left[\frac{1}{r} \right] + \hat{m}_y i k_y \mathcal{F} \left[\frac{1}{r} \right] + \hat{m}_z \mathcal{F} \left[\frac{1}{r} \right] \right)\end{aligned}\quad (4.12)$$

The second line in equation 4.12 is obtained using the differentiation property of Fourier transforms,

$$\frac{\partial^n}{\partial x^n} \frac{\partial^m}{\partial y^m} f(x, y) \leftrightarrow (i k_x)^n (i k_y)^m F(k_x, k_y) \quad (4.13)$$

in which the differentiation of a function $f(x, y)$ in the space domain is equivalent to a multiplication by a power of wavenumbers in each dimension of the Fourier domain. To solve the Fourier transform of $\frac{1}{r}$, we utilize the solution of a Hankel transform presented by Blakely (1995), resulting in the expression

$$\mathcal{F}[V] = -2\pi C_m \Theta_m e^{|k|(z_0 - z')} \quad (4.14)$$

for $z' > z_0$, where

$$\Theta_m = \hat{m}_z + i \frac{\hat{m}_x k_x + \hat{m}_y k_y}{|k|} \quad (4.15)$$

is a complex function of the wavenumbers k_x and k_y in the horizontal plane. We then expand $|\mathbf{B}_D| = |-\nabla_p V|$ by following a similar approach using the Fourier transform differentiation property to find the Fourier transform of $|\mathbf{B}_D|$, given by

$$\mathcal{F}[B_D] = -2\pi C_m \Theta_m |k| e^{|k|(z_0 - z')} \quad (4.16)$$

for $z' > z_0$. $\mathcal{F}[\psi_j]$ is then substituted in equation 4.10 for the above expression describing the Fourier transform of the magnetic field at a single point to obtain

$$\begin{aligned}\mathcal{F}[B_j(x, y, z_0)] &= \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} M_j(x', y') \int_{z_j}^{z_{j+1}} 2\pi C_m \Theta_m |k| e^{|k|(z_0 - z')} e^{-i(k_x x' + k_y y')} dz' dx' dy' \\ &= 2\pi C_m \Theta_m e^{|k|z_0} (e^{-|k|z_j} - e^{-|k|z_{j+1}}) \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} M_j(x', y') e^{-i(k_x x' + k_y y')} dx' dy' \\ &= 2\pi C_m \Theta_m e^{|k|z_0} (e^{-|k|z_j} - e^{-|k|z_{j+1}}) \mathcal{F}[M_j(x', y')]\end{aligned}\quad (4.17)$$

Finally, we replace the z parameters with the height $H_j = z_j - z_0$ between the source point Q and the observation point P , and the grid layer thickness $t_j = z_{j+1} - z_j$ to acquire

$$\mathcal{F}[B_j(x, y, z)] = 2\pi C_m \{\mathcal{F}[M_j]\} \{\Theta_m e^{-|k|H_j} (1 - e^{-|k|t_j})\} \quad (4.18)$$

It is important to note that the depth d below the surface described in Section 4.4 is measured from the surface to the grid layer, defined as $d = H - h$, where h is the altitude above the average Martian surface to the orbiter and H is the total altitude height from the top layer of the magnetization grid to the orbiter altitude. Equation 4.18 is then summed over each horizontal layer to calculate the total magnetic field at the observation point P above the surface, given by

$$\mathbf{B}(x, y, z) = \sum_{j=1}^N \mathcal{F}^{-1} \left[\mathcal{F}[B_j] \left(\frac{ik_x}{|k|} \hat{\mathbf{x}} + \frac{ik_y}{|k|} \hat{\mathbf{y}} + \hat{\mathbf{z}} \right) \right] \quad (4.19)$$

where $\mathcal{F}^{-1}[f]$ is the inverse Fourier transform on a function $f(x, y)$. Once the magnetic field vector is computed at each altitude H , the magnitude of the field can be calculated to generate an altitude profile of magnetic intensity above a fixed position $L(x, y)$ on the surface of the 3D grid. In the context of studying Martian magnetization, the point $L(x, y)$ could represent the center of an impact basin or the location of a lander on the surface.

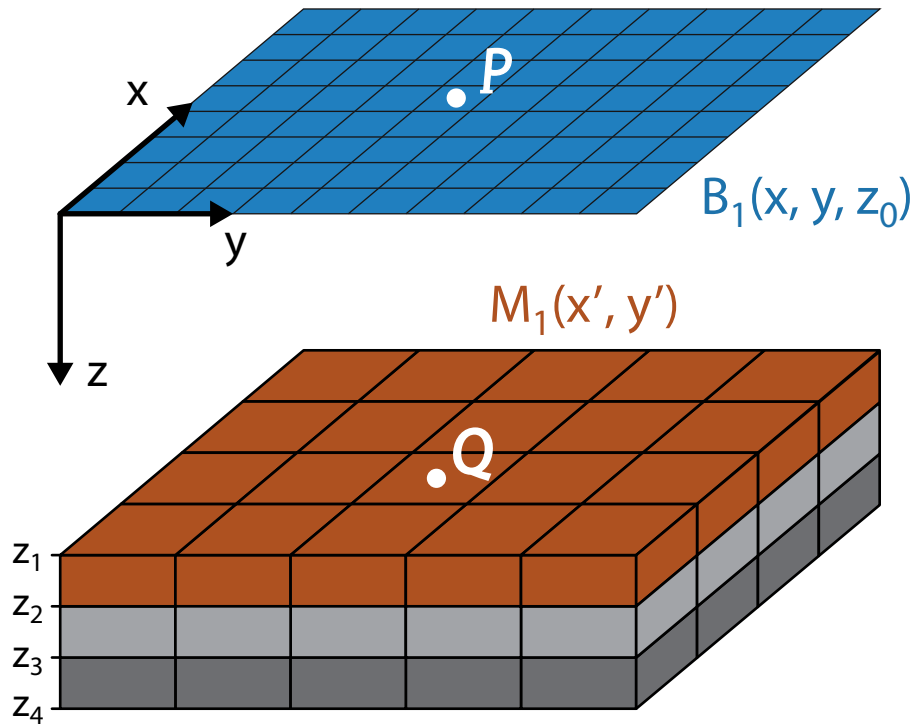


Figure 4.29: Simplified representation of calculating the magnetic field at some observation point P on a horizontal plane (blue) at an altitude z_0 above the surface caused by a 3D magnetization grid. The coordinate frame is defined such that x points north, y points east, and z is directed downward. The grid is divided vertically into three horizontal layers, each with a unique magnetization distribution M that only varies in the lateral dimension. The first layer (orange) extends from z_1 to z_2 with a source magnetization $M_1(x', y')$. The source point Q represents a given point in the grid layer with a source value following the form of a single dipole. The magnetic field $B_1(x, y, z_0)$ is computed solely from the influence of the first layer. The total magnetic field B at point P is then calculated by summing each individual contribution from the three layers.

Chapter 5

Conclusions & Future Work

In this dissertation, we have explored multiple domains of Heliophysics through the use of novel in situ measurements to showcase the range of scales and complexity within the inner heliosphere. To address the questions presented in each chapter, we conducted instrument calibration and processing, applied various data analysis techniques, and performed numerical simulations. The common goal across all chapters is to improve our understanding of the dynamics of electrons and magnetic fields in different environments between the Sun and Mars. Given the vast scale of the heliosphere, these regions are separated by considerable distances, spanning from 10 to $325 R_S$ (solar radii). However, these domains are intrinsically linked since the heliosphere is an interconnected system, where advancements in one area can contribute to our knowledge in an adjacent sub-field.

A basic introduction related to each study in this work was presented in Chapter 1. We began by examining the typical solar structure and solar wind conditions observed today and explored how solar behavior varies over different timescales, from tens to billions of years. In contrast to the ambient solar wind, we also discussed solar eruptive events, such as Coronal Mass Ejections (CMEs) and solar flares. The chapter further explored the influence of the Sun, both through steady solar wind conditions and these eruptive structures, on various planets. We specifically focused on the impacts of space weather at Mars, examining planetary atmospheric loss, the ancient geodynamo, and the remnant crustal magnetization observed on the surface of the red planet.

Our first study was introduced in Chapter 2, investigating the evolution of the electron strahl with heliocentric distance and other solar wind parameters. We calibrated the two electron electrostatic analyzers (ESAs) on board Parker Solar Probe (PSP) to produce accurate measurements of the electron velocity distribution function (eVDF). Prior to the PSP mission, the closest electron measurements were obtained by Helios (Schwenn et al. 1975), which reached approximately $62 R_S$ from the Sun. Recent observations made by PSP mark the closest in situ eVDF measurements to the Sun, spanning roughly 11 to $100 R_S$ from the solar center. Once bulk plasma parameters, such as electron density and temperature, were determined across this entire range of distances, we then calculated the angular width and direction of the strahl beam. We found the strahl beam to narrow with increasing electron

energy and vary with other parameters such as heliocentric distance, electron density, and electron collisional age. Our results suggest that Coulomb collisions are the main mechanism that scatters the strahl below roughly $30 R_S$. Outside of this heliocentric distance, other plasma wave instabilities such as whistler waves could scatter the strahl electrons more effectively as they propagate from the Sun.

Future work could study the electron measurements from upcoming PSP orbital encounters, which will reach within $10 R_S$ of the Sun. As our work is intrinsically biased by fewer measurements during the closest approach due to the orbit of PSP, obtaining additional measurements closer to the Sun can help validate our results in Chapter 2, especially as we enter a new phase of the solar cycle. In addition, PSP observations in conjunction with Solar Orbiter may provide more insight into electron strahl scattering along similar solar wind streams. Beyond the scope of this study, the calibration of ESAs is valuable to other areas of research in order to obtain bulk plasma properties in different contexts. Similar instrument and calibration techniques could be implemented for future missions to explore the outer heliosphere to provide new observations of the plasma environments near Uranus and Neptune (Mousis et al. 2018). While these methods of instrumentation are similar, the questions they help to address are crucial for assembling a comprehensive understanding of the heliosphere.

Chapter 3 continued to analyze electron observations made by PSP near the Sun, but for an Interplanetary Coronal Mass Ejection (ICME). The 05 September 2022 “Labor Day” ICME event became the first highly energetic and fast ICME encountered by PSP within $14 R_S$ from the Sun. Various ICME signatures were observed as the spacecraft traveled through the flank and wake of the structure, where some features were typical of ICMEs near 1 AU. Both in situ and remote-sensing measurements from PSP and other spacecraft led to a more complete picture of the event, with our model suggesting two separate CME structures. Moreover, this ICME event appeared to be related to a large reconfiguration of the Heliospheric Current Sheet (HCS). Thus, ICMEs may play a role in balancing the energy budget of the Sun, having a greater impact on the heliosphere as a whole.

In addition to the “Labor Day” ICME, numerous studies have been published (e.g., Braga et al. 2024; Palmerio et al. 2024) on other solar eruptive events observed by PSP. Many of these studies incorporate in situ measurements from other spacecraft, such as Solar Orbiter or BepiColombo, to refine our understanding of ICME evolution. Improving our models of ICMEs as they propagate from the Sun can benefit our knowledge of solar eruptive mechanisms, ICME arrival times, and potential impacts on planetary magnetospheres. As we approach solar maximum of solar cycle 25, PSP is expected to continue observing ICMEs, either in situ or remotely. These future events observed near the Sun can contribute to our understanding of ICME eruption and propagation throughout the heliosphere, as well as the effects on solar and planetary environments.

In Chapter 4, we analyzed surface magnetic field measurements at Mars to model crustal magnetization near the InSight and Zhurong landing sites. The crustal fields observed near the surface are theorized to be remanent of an ancient global magnetic field once present on Mars. These localized regions of magnetization could have undergone various natural

remanent magnetization processes linked to the geology and chemistry near the surface. We were able to constrain the magnetization characteristics near each lander using the first surface measurements of the Martian magnetic field. Furthermore, we developed new models of crustal magnetization to link length scales of magnetization to geological regions corresponding to different time periods in Martian history. We provided evidence that Mars experienced at least one polar field reversal during the presence of a global magnetic field. In addition, our work suggests that a new geodynamo was generated sometime during the Hesperian period after an initial dynamo on Mars weakened or disappeared.

The surface magnetic field measurements at Mars contributed significantly to constraining the magnetization below each lander. Before InSight and Zhurong, all crustal field models (e.g., Langlais et al. 2019) relied on satellite-based measurements above roughly 150 km from the surface. At these orbital altitudes, length scales of magnetization smaller than 150 km can not be resolved due to the limitation of downward continuation. This issue motivates the need for more measurements obtained below 150 km at different locations on Mars. The only landers equipped with magnetometers are located in the northern hemisphere of the planet, but most of the strong magnetization lies in the heavily cratered southern highlands. Future landers with magnetometers could be positioned in other areas of the planet to refine our current global crustal models. Another advancement would be to mount magnetometers on helicopters (Bapst et al. 2021), balloons (Blamont 1987), or eventually sounding rockets (Noga et al. 2023) at Mars. For example, a small helicopter could collect measurements both at the surface and above the ground during flight at very low altitudes. A sounding rocket mission on Mars could then provide the remaining observations between helicopter and orbital magnetic field measurements for a complete altitude profile of the crustal fields. Only through the combination of these methods can we better constrain the timing of the Martian geodynamo, which will enhance our understanding of other unresolved questions, such as climate evolution on Mars.

As our technology advances, our approach to conducting science must also follow suit. With the growth of our diverse database driven by upcoming space missions and improved instrumentation, we must maintain the tools and methodology from previous studies to allow for seamless integration with future research. This will further facilitate collaborations and interdisciplinary frameworks to provide significant benefits to the scientific community, with the intention of connecting different aspects of the world around us.

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