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Radially sheared azimuthal flows and turbulent transport in a cylindrical plasma

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A radially sheared azimuthal flow is observed in a cylindrical helicon plasma device without any apparent external sources of angular momentum input. Broadband fluctuations combined with a chain of coherent structures are observed, turbulent particle transport across the shear layer is inhibited, and energy appears to be transferred from linearly unstable intermediate wave numbers into both larger and smaller spatial scales that are linearly stable. The shape of the radial plasma potential profile associated with the azimuthal flow is in agreement with published theory, and the flow magnitude is consistent with estimates of the turbulent Reynolds stress. © 2004 American Institute of Physics. [DOI: 10.1063/1.1794752]

I. INTRODUCTION

Theory suggests that sheared flows in toroidal plasma devices,¹ associated with the low to high-confinement mode (L-H) transition² and the formation of internal transport barriers in the hot central plasma of such fusion confinement devices³ are both triggered by the turbulent Reynolds stress.^{4,5} Once formed, these shear flows can act to reduce the rate of cross-field turbulent mixing as has been demonstrated in the edge of tokamaks.^{6–8} This reduction of the cross-field particle (heat) flux can cause the mean density (temperature) profiles to evolve steeper gradients across the shear layer region. The resulting pressure gradient can then sustain a radially sheared radial electric field which maintains the reduced transport rates by the process of turbulent shear decorrelation.^{9,10}

Even in the absence of large mean shear flows, linearly damped time-varying sheared zonal flows can be nonlinearly coupled to the smaller turbulent scales by the turbulent Reynolds stress, resulting in a transfer of kinetic energy from the turbulence into the fluctuating shear flows.¹¹ This energy transfer rate depends upon the shear flow amplitude which, in turn, depends upon the nonviscous shear flow damping processes. Examples of such damping processes include ion-neutral elastic collisions (as in our experiment), collisions between trapped and untrapped populations (as in a torus), or topographic friction (in the case of geophysical flows). As a result turbulence/shear flow interactions are also thought to play an important role in the saturation of turbulent transport¹² even in the absence of transport bifurcations and transport barrier formation. Thus turbulence-driven shear flows are thought to play a critical role in the overall confinement characteristics and qualities of confined magnetic fusion plasmas. Similar processes have also been implicated

in the formation of shear flows in planetary and stellar atmospheres¹³ and thus are also of broader scientific interest.

Motivated by the significance of turbulence/shear flow interactions, a number of recent experimental papers have addressed aspects of this problem. The Reynolds stress was reported to evolve in a manner consistent with shear layer formation during a transition to improved Ohmic confinement (a form of L-H transition) at the edge of the HT-6M tokamak.¹⁴ Recent probe studies in DIII-D show that the phase coherence between low-frequency and high-frequency fluctuations increased during the L-H transition^{15,16} in a manner that is qualitatively consistent with the existence of Reynolds stress driven L-H transitions. Turbulence imaging studies of time-stationary L-mode plasmas using beam emission spectroscopy (BES) in DIII-D and probe studies on the H-1 heliac both indicate the presence of time-varying shear flows that are nonlinearly coupled to higher frequency fluctuations^{17,18} and, in the case of the H-1 heliac, modulate the higher frequency fluctuation amplitudes. Thus several pieces of evidence exist that support the hypothesis that plasma turbulence can generate organized shear flows.

These experimental studies leave a number of relevant turbulence/shear flow interactions physics issues unanswered. In the work on H-1 the ions' gyro-orbits themselves are large compared to the other spatial scales in the system. In addition, in the fluctuating shear, flow was nonlinearly coupled to distinct coherent finite amplitude perturbations in the plasma and not with broadband turbulent fluctuations. Thus the results are in a qualitatively different regime from that encountered in larger confinement devices. Furthermore the drift velocity ordering that was assumed in the theoretical models that underpin these studies is violated, rendering conclusions about the existence of an inverse energy transfer suspect. The increase in phase coherency during the L-H transition in DIII-D does not indicate the direction of energy transfer and thus provides no conclusive indication that the transition is actually triggered by the Reynolds stress. We

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also note that there is disagreement in the literature over the usefulness of the autobicoherence as a means of detecting Reynolds-stress driven shear flows.^{19,20} The BES observations show interactions between turbulence and zonal flows in an L-mode discharge but provide no evidence for the direction of energy exchange between the two types of fluctuations, and results from L-H transitions are not presently available due to diagnostic limitations. Finally, in all of these experiments the free energy sources and the dissipation mechanisms as well as the magnetic geometry are complex, and the plasma neoclassical flow damping has a nonlinear velocity dependence that can lead to bifurcation phenomena that are unrelated to the turbulent Reynolds stress.²¹ Thus observation of shear flow generation from turbulence in simpler physical systems which can then be studied in detail would be a valuable addition to the experimental tests of the basic theory.

II. DESCRIPTION OF EXPERIMENT

In this paper we report observations of a radially sheared azimuthal flow that appears to be consistent with a turbulence-driven shear flow in a cylindrical plasma device with no apparent external angular momentum input. The experiments were performed in the controlled shear decorrelation experimental (CSDX) device in which collisional plasmas (density $\sim 10^{19} \text{ m}^{-3}$ electron temperature $T_e \sim 3 \text{ eV}$, ion temperatures $T_i \sim 0.2\text{--}0.8 \text{ eV}$) are generated.²² The plasma source is composed of a 10 cm diameter quartz tube surrounded by a two-loop azimuthally symmetric ($m=0$) helicon wave antenna, and is attached to a 3 m long vacuum chamber with an inside diameter of 20 cm. Both source and vacuum chamber are immersed within a solenoidal magnetic field that can be varied from 0 up to 1000 G. The working gas is injected at the exit plane of the plasma source, and variable speed pumping is provided by a manual throttle valve and 1000 l/s turbopump located 3 m downstream from the source. For all of these experiments argon plasmas were used with a fill pressure of 3 mTorr and a source power of 1500 W at a radio frequency (rf) of 13.56 MHz and a magnetic field of 1000 G. The rf matching network was adjusted such that less than 10–20 W of power were reflected, and measurements of the antenna voltage and phase in an identical plasma source used for plasma processing applications indicates that typically greater than 80% of the forward-going rf power is absorbed in the plasma itself (presumably the balance of the power is dissipated in the antenna, match network, and electrically conducting structure immediately surrounding the antenna).

In the plasma source region the magnetic field lines terminate on insulating surfaces, while at a distance of 1.9 m downstream from the exit plane of the plasma source the magnetic field lines terminate on a series of ten conducting annular rings. Each ring has a width of about 3 mm and is separated from the next ring by a gap of 4 mm (a distance that is slightly larger than the ion gyroradius $\rho_i \sim 2\text{--}3 \text{ mm}$ at these conditions). In all of the experiments reported here the end rings were electrically insulated from each other and from the vacuum chamber inhibiting the radial transport of

FIG. 1. Time-averaged radial profiles of (a) plasma density n_0 , (b) electron temperature T_e , (c) measured (open circles) and predicted (solid line) plasma potential, (d) broadband potential fluctuation amplitude (solid line) and low-frequency azimuthally symmetric potential fluctuation amplitude (dash-dot), (e) turbulent cross-field particle flux $\tilde{\Gamma}_r = \langle \tilde{n} \tilde{v}_r \rangle$, and (f) effective turbulent diffusivity, $D_r^{\text{eff}} = \tilde{\Gamma}_r / \nabla n_0$. Data in panel (d) are normalized to equilibrium electron temperature at $r=3 \text{ cm}$. The electron temperature errors are estimated from the quality of the Langmuir probe fitting; this error is then propagated through a model of the sheath potential drop to estimate the errors in the space potential profile shown in panel (c).

current between rings. In addition, the vacuum chamber wall in the region behind the end-ring assembly was covered with dielectric material to prevent currents from flowing through the residual plasma, field lines, and chamber wall in this region. Thus any radial currents must flow within the plasma located between the source and end-ring assembly via non-ambipolar transport mechanisms (which have been shown to be equivalent to a Reynolds-stress effect⁴).

Measurements of the equilibrium plasma conditions were made with a rf-compensated swept-bias Langmuir probe using established techniques²³ which could be scanned radially across the plasma column $\approx 1 \text{ m}$ downstream of the source exit plane. Measurements of the fluctuating ion saturation current (I_{sat}) and floating potential (V_f) were made with a four-tip probe assembly which had two I_{sat} probe tips separated azimuthally by 0.5 cm, aligned along the field line with two V_f tips that were also separated azimuthally by 0.5 cm.

III. RESULTS AND DISCUSSION

The measured equilibrium density and electron temperature are shown in Figs. 1(a) and 1(b), respectively. The plasma density peaks at a value of $2.4 \times 10^{19} \text{ m}^{-3}$ and decays with increasing radius. The temperature in the inner 2.5 cm region is about 3.2 eV and then decays at larger radii, consistent with a heat source being localized inside of the antenna radius. Thus the pressure gradient contains contributions from both temperature and density gradients. The measured equilibrium plasma potential profile, normalized to the T_e at $r=3 \text{ cm}$, which is the location of the maximum density gradient and taken as a reference point for all normalizations throughout this paper unless otherwise stated, is shown with open circles in Fig. 1(c). The plasma potential peaks at the center of the plasma, has a minimum at about

FIG. 2. (Color). (a) Radial profile of potential autopower spectrum and (b) two-point estimated spectrum $S(\omega, k_\theta)$ of potential fluctuations measured at $r=3$ cm. Log_{10} of the autopower amplitude is shown in panel (a). Data in (b) obtained at $r=3$ cm.

$r \sim 4$ cm, and then begins to increase in value for $r > 4$ cm. Thus a positive radial electric field exists in the inner portion of the plasma column and a negative electric field exists outside of a shear layer that exists at the potential minimum located near $r=4$ cm. The Doppler shifted azimuthal fluctuation phase velocity (not shown here) is consistent with the combination of the azimuthal $E \times B$ flow profile inferred from the measured space potential and the electron diamagnetic drift velocity.

Using arguments based upon the conservation of angular momentum, the shape of the time-averaged plasma potential profile associated with a radially sheared azimuthal $E \times B$ flow formed via an inverse energy transfer process was predicted using theory and numerical simulation.²⁴ In the theory of Ref. 24, the radial transport of turbulent angular momentum via the turbulent Reynolds stress, or equivalently a nonlinear transfer of turbulent kinetic energy to large azimuthal scales, leads to the build up of two regions of oppositely directed azimuthal flow that conserves angular momentum. A shear layer is then formed at the interface between these two regions of flow. The theory also assumes that the electron temperature is uniform and that the plasma density has a form $n(r) = n_0[0.9 \exp(-2r^2/a^2) + 0.1]$ where the radius a defines the size of the plasma in the simulations. The density profile shown in Fig. 1(a) is reasonably close to this form for $a \approx 3.8$ cm. Using this value to scale the system size we find that the predicted potential profile agrees well with the measured space potential profile [see solid line in Fig. 1(c)].

Since the published theory does not provide an absolute magnitude of the shear flow, only the relative shape of the shear profile can be compared with the experimental data. However, as we show later in this paper, the magnitude of the measured shear flow appears to be consistent with initial estimates of the turbulent Reynolds stress.

The radial profiles of the low-frequency (i.e., $\omega < \Omega_{ci}$) broadband electrostatic potential fluctuation amplitude, normalized to the electron temperature at 3 cm, is shown in Fig. 1(d). The density fluctuation amplitude profile, normalized to the mean density at $r=3$ cm, has a similar shape and magnitude. The broadband fluctuations range in frequency from about 3–4 kHz up to 25–30 kHz, and have a peak amplitude near the maximum equilibrium pressure gradient at $r=3$ cm. These fluctuations also satisfy the Boltzmann relation $|\tilde{n}/n_0| \sim |e\tilde{\phi}/kT_e|$, and density fluctuations lead the potential fluctuations by a $\sim(\pi/2)$ phase shift. Given the presence of large parallel electron dissipation, we note that these observations are consistent with expectations for highly resistive drift-wave turbulence.²⁵ In addition to these turbulent fluctuations, a large amplitude narrow-band (frequency width $\Delta f \approx 1$ kHz) coherent oscillation at 10–12 kHz is also observed, as shown in Fig. 2(a). Note that there are also weaker coherent oscillations at several other frequencies, but that the dominant coherent fluctuation activity is contained within the 10–12 kHz band. Integrating over this narrow-band spectral component, we find that the potential fluctuation has a root

mean square (rms) amplitude $\tilde{\phi} \sim 0.5$ V. The narrow-band mode amplitude is peaked near the minimum in the mean potential profile, suggesting that it is somehow associated with the presence of the shear flow. Using the azimuthally separated probes we have measured the frequency resolved local azimuthal wave number spectrum, $S(\omega, k_\theta)$, for the potential and density fluctuations using the two-point technique.²⁶ In this technique the measured azimuthal wave number is taken as the phase delay incurred as fluctuations propagate between two azimuthally separated probes, divided by the 0.5 cm separation distance and thus represents a measure of the local wave number. A typical result for potential fluctuations obtained at $r=3$ cm is shown in Fig. 2(b). The 10–12 kHz coherent oscillation has azimuthal wave numbers ranging from 1 – 1.6 cm^{-1} at $r=3$ cm, corresponding to an azimuthal mode numbers ranging from $m=3$ – 5 . The autocorrelation time of this narrow-band feature is ~ 1 ms, much longer than the estimated eddy turnover time $\tau_{\text{eddy}} = (L_\perp / \tilde{V}_E) \sim (2\pi B) / k_\perp^2 \tilde{\phi} \sim 10$ – 30 μs , consistent with a long-lived coherent structure. A detailed discussion of the transition to the turbulent state, the generation of these coherent structures, and the identification of the origins of the fluctuations were presented earlier²² and will be discussed in a separate paper.²⁷

Examination of Fig. 2(a) also shows that the potential fluctuations also contain a significant component over a frequency band of 0–2 kHz, most clearly seen in the $r > 4$ cm region. This class of fluctuation is much weaker in the density fluctuation spectrum. Examining Fig. 2(b), we see that the frequency is generally proportional to the azimuthal wave number, and thus we surmise that these very low-frequency potential fluctuations must correspond to very low k_θ fluctuations [note that the data in Fig. 2(b) are obtained at $r=3$ cm where these very low-frequency fluctuations are weak]. In order to quantify the azimuthal scale of these potential fluctuations we measured the local wave number spectrum, $S(k_\theta, \omega)$, for the potential and density fluctuations at a variety of radial positions, and then integrated over frequency to give the local wave number spectrum $S(k_\theta)$ for different radial locations. The result for potential fluctuations measured at $r=3$ cm, 4 cm, and 5 cm is shown in Fig. 3. This result indicates that these very low-frequency potential fluctuations located primarily between $r=4$ cm and $r=6$ cm have a low effective mean azimuthal wave number $\bar{k}_\theta \leq 0.1$ cm^{-1} and have a significantly broad variation about this mean $\Delta k_\theta \sim \pm 0.3$ cm^{-1} (again, these values represent measurements of a *local* wave number as discussed above). The small mean wave number of these potential fluctuations indicates that they are nearly azimuthally symmetric (i.e., $m=0$), and the finite width of Δk_θ corresponds to azimuthal perturbations with a wavelength between 10 and 20 cm (i.e., slightly smaller than the plasma circumference over the $r=4$ – 6 cm range where these fluctuations are largest). Integrating the potential fluctuation spectrum shown in Fig. 2(a) over the 0–2 kHz spectral range gives the radial profile of the azimuthally symmetric potential fluctuation amplitude $\tilde{\phi}_{m=0}$. This result, normalized by the electron temperature at $r=3$ cm, is shown as the dashed line in Fig. 1(d). The maxi-

FIG. 3. Local wave number estimate $S(k_\theta)$ for $r=3$ cm, 4 cm, and 5 cm found by integrating $S(\omega, k_\theta)$ over frequency.

mum fluctuating shear flow amplitude $\tilde{\phi}_{m=0}^{\text{max}} \approx 0.6$ V occurs around $r=5$ cm, slightly outside of the minimum in the mean plasma potential profile. There is a significant gradient in the amplitude of this potential fluctuation between $r=3$ cm and $r=5$ cm suggesting that there is some radial structure to this azimuthally symmetric potential fluctuation.

Simultaneous measurements of these potential fluctuations at multiple radial locations are not available and thus at this time we can only estimate the radial scale length of these azimuthally symmetric potential fluctuations. Noting that the gradient in $\tilde{\phi}_{m=0}$ is significant only between $r=3$ cm and $r=5$ cm, we estimate the radial scale length of the $m \sim 0$ potential fluctuations as $\tilde{L}_r^{m=0} \approx 2$ cm yielding an estimated range for the effective radial wave number k_r of these fluctuations as $\pi / \tilde{L}_r^{m=0} < k_r < (2\pi) / \tilde{L}_r^{m=0}$. Here the factor of π would denote a half-wave length potential perturbation existing between $r=3$ cm and $r=5$ cm, whereas the factor of 2π would denote a full wavelength perturbation existing between these points. Using our estimate for $\tilde{L}_r^{m=0}$ we then estimate $k_r \sim 2$ – 3 cm^{-1} . We conclude that a mean shear flow exists in the plasma, and also that nearly axisymmetric low-frequency potential fluctuations exist, and that this fluctuating plasma potential has $|k_r| \gg |\bar{k}_\theta|, |\Delta k_\theta|$ consistent with expectations for fluctuation-driven shear flows.

The radial profile of the time-averaged turbulent particle flux $\tilde{\Gamma}_r = \langle \tilde{n} \tilde{v}_r \rangle = \langle \tilde{n} \tilde{E}_\theta \rangle / B$ and the radial equilibrium density gradient ∇n_0 , computed from the data in Fig. 1(a), are shown in Fig. 1(e). These data are then used to calculate an effective radial turbulent diffusion coefficient $D_r^{\text{eff}} \equiv -\tilde{\Gamma}_r / (\nabla n_0)$ shown in Fig. 1(f). Note that here we calculate this diffusivity merely because it represents a readily understood measure of the rate of cross-field transport, and we note that transport may be, in part, caused by nondiffusive processes as has been recently shown in other work.^{28–31} The turbulent particle flux and effective diffusivity both vanish at the shear

FIG. 4. (Color). Autocorrelation $b(f_1, f_2)$ of potential fluctuations at $r = 4$ cm.

layer, indicating that particle transport is inhibited at the shear layer location.

A net transfer of energy between turbulent fluctuations and shear flows must be accompanied by a finite phase coherence between the two associated space-time scales. It has been argued that this phase coherence can be represented by the autocorrelation^{32,33} (Kim: 1979) computed from potential fluctuations.^{15,34} However, we also note that there is disagreement in the literature over the usefulness of the potential fluctuation autocorrelation in detecting Reynolds-stress driven shear flows. In the work of Holland *et al.*²⁰ this quantity was found to contain the expected signatures during the formation of a strong shear flow obtained in numerical simulations. However, in the work of Ramisch *et al.*¹⁹ it is argued that the autocorrelation does not provide a useful measure of the appropriate nonlinear dynamics and the crossbicoherence computed between fluctuating electric fields must instead be used to detect turbulent-driven shear flows. At this point in time this issue is unsettled and must be resolved by further work. For the moment we present measurements of the autocorrelation of the potential fluctuations in order to search for nonlinear coupling across different frequency scales. In particular, we are interested in determining if there is any evidence that the turbulence fluctuations might be coupled to the mean and fluctuating shear flow that exists in this experiment.

Measurements of the potential fluctuation autocorrelation $b(f_1, f_2)$, defined as

$$b^2(f_1, f_2) = \frac{|\langle \phi^*(f_1 + f_2) \phi(f_1) \phi(f_2) \rangle|^2}{\langle |\phi(f_1 + f_2)|^2 \rangle \langle |\phi(f_1) \phi(f_2)|^2 \rangle} \quad (1)$$

taken at $r = 4$ cm are shown in Fig. 4. It is important to use a large number of realizations with time-stationary data to ensure that the bicoherence estimator is converged³⁴ and has a small variance, given as $\text{var}(b) \approx 1/M[1 - b^2(f_1, f_2)]$ where $b^2(f_1, f_2)$ is the squared bicoherence for the frequency triad $(f_1 + f_2, f_1, f_2)$ and M represents the available number of statistically independent ensembles.³² In the results shown here $M = 10^3$ giving a converged estimate for the bicoherence²²

and typical values for the bicoherence lie in the range $b(f_1, f_2) \approx 0.1 - 0.2$, giving an estimated variance $\text{var}(b) \leq 10^{-3}$ in the bicoherence estimates shown below. Thus bicoherence values above $\sqrt{\text{var}(b)} \approx 0.03$ represent real phase coherence between frequencies satisfying the usual summation rule. The data in Fig. 4 show a non-negligible bicoherence $b(f_1, f_2) \approx 0.15 - 0.2$ occurring between the broadband fluctuations distributed from 3–30 kHz and the $m \sim 3 - 5$ coherent fluctuations at $\sim 10 - 12$ kHz. The bicoherence between the low-frequency shear-flow fluctuations (0–2 kHz) and the $m = 3 - 5$ coherent fluctuations ranges between 0.15–0.2 indicating that energy is also exchanged between these modes and the fluctuating shear flow. A weaker interaction (bicoherence of around 0.1) between the broadband turbulence (3–30 kHz) and the fluctuating shear flow (0–2 kHz) is also seen. Thus energy appears to be exchanged between all three types of fluctuations, with the turbulence/coherent mode and coherent mode/shear flow interactions being of equal strength, while the turbulence/shear flow interaction is slightly weaker. The coherent mode may act in some way to mediate energy exchange between the shear flow and the higher frequency turbulent fluctuations as was suggested elsewhere.²⁰ However, this point requires further investigation before firm conclusions can be drawn.

Obviously for a large-scale flow to be driven by smaller scaled turbulent velocity fluctuations, net energy must be transferred from the turbulence scales into the larger scaled shear flow. The bicoherence results above indicate the presence of phase coherence (which as noted above must exist if three-wave coupling is to nonlinearly transfer energy), but they say nothing about the crucial issue of the net energy transfer direction. We can make progress on this issue by noting that there are two free energy sources—a pressure gradient and an $E \times B$ velocity gradient—available to drive the fluctuations in our experiment. The development of the turbulent spectrum can then be described by a balance of the growth of the power at a given wave number, due to these instabilities, and nonlinear transfer of this power to other wave numbers where the fluctuations are then damped. For weakly dissipative drift turbulence (where the parameter $k_{\parallel}^2 v_{th}^2 / \nu_e \omega \gg 1$) this process is described by the wave kinetic equation of the form^{35,36}

$$\frac{1}{2} \frac{\partial \langle |\phi_k|^2 \rangle}{\partial t} = \gamma_k \langle |\phi_k|^2 \rangle + T_k, \quad (2)$$

where the nonlinear transfer of energy is given as $T_k = \sum_{k=k'+k''} \Lambda_{k',k''}^k \text{Re} \langle \phi_k^* \phi_{k'} \phi_{k''} \rangle$ with the coefficient $\Lambda_{k',k''}^k = 1/2 \rho_s^3 c_s (\vec{k}' \times \vec{k}'') \cdot \hat{z} (|\vec{k}''|^2 - |\vec{k}'|^2)$, and the potential fluctuations are written in normalized form with $\phi_k = (e \tilde{\phi}_k) / k T_e$; γ_k denotes the linear damping/growth rate of the k th mode. Note that the energy transfer is related to the bispectrum $\langle \phi_k^* \phi_{k'} \phi_{k''} \rangle$, which is closely related to the bicoherence discussed above (analogous to the relationship between a cross-spectrum and coherence). Thus the time-stationary net energy transfer in such a system can then be given as $T_k = -\gamma_k \langle |\phi_k|^2 \rangle$. We explicitly note that in the case of strong parallel electron dissipation $k_{\parallel}^2 v_{th}^2 / \nu_e \omega \ll 1$ this one-field description is supplanted by the two-field Hasegawa–Wakatani

FIG. 5. Azimuthal wave number spectrum of normalized density and potential fluctuations, and along with linear growth rate of collisional drift waves with mean shear flow effects included. Spectra are measured at $r=3.5$ cm.

model.³⁷ For our experiment we estimate that $0.1 < k_{\parallel}^2 v_{th_e}^2 / \nu_e \omega < 1$ and thus the one-field description summarized above does not strictly hold. However, even in this more complex two-field description of drift turbulence, the basic balance between free energy release within an unstable range of scales and the subsequent nonlinear transfer of this energy to stable scales *does* hold. This argument thus implies that an examination of the k spectrum of the fluctuations and an examination of the linear stability of the plasma equilibrium can provide insight into the direction of energy transfer in this experiment.

As discussed earlier, we have measured the local azimuthal density and potential wave number spectra $n_{k_{\theta}}^2$ and $\phi_{k_{\theta}}^2$ using the two-point technique,²⁶ which provides legitimate results for fluctuations that propagate in one direction (as is the case here). As shown in Fig. 5, these spectra peak at an intermediate wave number, $k_{\theta} \rho_s \sim 2$, and decay as $\sim k_{\theta}^{-5.4 \pm 0.2}$ for higher wave numbers. We caution the reader that the exponent a depends upon radial location (e.g., $|\phi_k|^2 \propto k_{\theta}^{-a(r)}$), and so the value of this exponent is not “universal.” This is probably an indication that the turbulence is far from being homogeneous, and that the boundaries and shear layer strongly influence the turbulence properties. Interestingly, the $k_{\theta} \rho_s < 1$ portion of the spectrum (corresponding to frequencies below 5 kHz) has appreciable amounts of power (i.e., significantly higher than the noise level of the measurement). If this spectral region is linearly stable, then by the arguments discussed above, the power that is contained in this spectral region likely appears there because it was transferred nonlinearly from other spectral regions.

Motivated by these considerations, we have examined the linear stability of the plasma equilibrium. In an experiment with large-scale flows it is crucial to include relevant descriptions of the linear damping of such flows in the stability analysis. Here such damping comes about from ion-neutral collisions which create a drag force on the ion fluid as it

drifts through the background stationary neutral gas. Ion viscous damping, effective on small spatial scales, also plays a role in dissipating the fluctuation energy. However, this viscous damping plays an insignificant role in the damping of the large-scaled shear flows. Thus in our linear analysis we include an effective ion-neutral drag on the ions in a form identical to that used in recent studies of the ionosphere,³⁸ but we have neglected viscous damping which will affect the high- k spectral region (and thus is not important for the consideration of the origins of the large-scale shear flows).

The linear radial eigenmode equation for poloidal mode-number m , including both the pressure profile and the mean shear flow profile along with ion-neutral flow damping, is then given as³⁹

$$\frac{\partial^2 \phi_m}{\partial x^2} + \left(\frac{1}{x} - \frac{a^2}{L_n^2 x} - \frac{1}{A_5} \frac{dA_5}{dx} \right) \frac{\partial \phi_m}{\partial x} + \left\{ \frac{m \Omega_{ci} A_1}{(\omega_- + i \nu_{in}) x} \left(\frac{1}{A_5} \frac{dA_5}{dx} - \frac{1}{A_1} \frac{dA_1}{dx} \right) + A_5 \frac{a^2}{\rho_s^2} \left[\frac{\omega_{\parallel}}{\omega_{\parallel} - \omega_-} \frac{\omega^* - \omega_-}{\omega_{\parallel} - \omega_- + i \nu_{in}} + \left(\frac{A_1}{A_5} - 1 \right) \frac{\omega^*}{\omega_- + i \nu_{in}} \right] - \frac{m^2}{x^2} \right\} \phi_m = 0, \quad (3)$$

where $x = r/a$, $\phi_m = e \delta \phi_m / T_e$, $\omega_{\parallel} = (k_{\parallel}^2 V_{Te}^2 / \nu_e) / (C_s / L_n)$, $\omega_- = \omega - \omega_{E \times B}$, $\omega_{E \times B} = -(m \rho_s C_s / r) d\phi_0 / dr$, $\omega^* = -(m \rho_s C_s / r) d \ln n_0 / dr$, $A_1 = 1 + 2 \omega_{E \times B} / m \Omega_{ci}$, $A_2 = 1 + 2 \omega_{E \times B} / m \Omega_{ci} + (r d \omega_{E \times B} / dr) / m \Omega_{ci}$, and $A_5 = A_1 A_2 - (\omega_- + i \nu_{in})^2 / \Omega_{ci}^2$. In our numerical solutions, we approximate the equilibrium density profile by a Gaussian, $n_0(r) = n_{peak} \exp[-0.5(r/L_n)^2]$, which implies that the diamagnetic frequency $\omega^* = m \rho_s C_s / L_n^2$ is independent of radius, and assume that the equilibrium potential ϕ_0 can be approximated by the predicted profile of Hasegawa and Wakatani [shown in Fig. 1(c) to agree reasonably well with the measured mean potential profile]. Solutions of the eigenmode equation using the measured CSDX parameters were found using a shooting technique. The results of this shooting code were validated against integral method calculations analogous to those discussed in Perkins and Jassby⁴⁰ with excellent agreement. It is found that two different branches (corresponding to the two lowest radial modenumbers) are unstable over a wide range of modenumbers (shown in Fig. 5); the first branch is unstable for $m \geq 1$ while the second is unstable at $m \geq 2$. However, it can also be shown from the eigenmode equation that the $m=0$ mode is always damped, with a damping rate γ_0 determined from the dispersion relation

$$\gamma_0^2 - [\omega_{\parallel}(1 + k_0^{-2} \rho_s^{-2}) + \nu_{in}] \gamma_0 + \omega_{\parallel} \nu_{in} = 0, \quad (4)$$

where k_0 forms an effective radial wave number for the perturbations and is given as

$$k_0^2 = \frac{\int_0^a dr r n_0(r) \left| \frac{\partial \phi_0}{\partial r} \right|^2}{\int_0^a dr r n_0(r) |\phi_0|^2} = \frac{\int_0^a dr r \exp\left(-\frac{1}{2}\left(\frac{r}{L_n}\right)^2\right) \left| \frac{\partial \phi_0}{\partial r} \right|^2}{\int_0^a dr r \exp\left(-\frac{1}{2}\left(\frac{r}{L_n}\right)^2\right) |\phi_0|^2}. \quad (5)$$

Physically, this result is intuitive, as the $m=0$ mode has no radial velocity fluctuation associated with it, and so cannot tap the free energy gradient to extract energy. As mentioned above, this analysis does not include ion viscous damping and as a result the growth rates for the $k_{\theta} \rho_S > 1$ region will be lower than what is shown; however, this will not substantially influence the $k_{\theta} \rho_S \ll 1$ stability analysis. We therefore conclude based upon the linear eigenmodes analysis and upon this analytical argument that the $k_{\theta} \rho_S \ll 1$ portion of the spectrum is linearly stable. By the arguments outlined earlier, the power in this region (corresponding to the low-frequency 0–2 kHz fluctuating shear flow, $\bar{m}=0$) must therefore have been transferred into this range via a nonlinear process. Power must also be transferred nonlinearly into the higher wave number region as well where it is presumably dissipated by finite Larmor radius ion viscosity. The bicoherence analysis presented earlier indicates that finite phase coherence exists between the very low frequencies corresponding to the shear flow and the other frequencies corresponding to coherent fluctuations and broadband turbulence. Thus two key conditions for turbulence-driven shear flows—a net transfer of energy into large-scale shear flows and phase-coherent wave-wave interactions—appear to be satisfied in this experiment.

The shear flows must exceed minimum levels in order to have a substantive effect on the turbulent transport. From the spectra shown in Fig. 2(a), we estimate that the decorrelation rate of the broadband fluctuations $\nu_{corr} \sim \Delta\omega \sim 6 \times 10^4 \text{ s}^{-1}$ where $\Delta\omega$ is the spectral width of the fluctuations near the shear layer. The shearing rate of the mean flow is defined as⁹ $\omega_{sh} = k_{\theta}(\partial \bar{V}_{\theta}) / \partial r L_r^{corr}$ where $\bar{V}_{\theta}$ is the mean (i.e., steady-state) azimuthal $E \times B$ drift velocity determined from the mean potential profile shown in Fig. 1(c), and we have taken the average azimuthal wave number of the fluctuations to be about $k_{\theta} \sim 2 \text{ cm}^{-1}$ as measured by the azimuthally separated probe tips. From Fig. 1(c) we find $(\partial V_{\theta}) / \partial r \sim 10^5 \text{ s}^{-1}$ and we estimate that the radial correlation length of the turbulence is roughly $L_r^{corr} \approx L_n = 2 \text{ cm}$. We then estimate that $\omega_{sh} \approx 4 \times 10^5 \text{ s}^{-1}$ and the hybrid decorrelation rate of Biglari *et al.* (which is the net decorrelation rate in the presence of both turbulent mixing and shear flow) then becomes $(2\omega_{sh}^2 \nu_{corr})^{1/3} \approx 2-3 \times 10^5 \text{ s}^{-1}$, i.e., a few times higher than ν_{corr} . These estimates suggest that the mean shear flow would then be expected to impact the turbulent radial correlation lengths and transport rates via the usual arguments.⁹

The discussion about the significance of the fluctuating shear flow proceeds slightly differently. Recent theoretical

work argues that fluctuating shear flows decorrelate turbulent fluctuations via a series of random shearing events, each of which induces a small increase in the radial wave number of the turbulent fluctuations.^{41,42} Thus this series of shearing events is viewed as a diffusion of fluctuation energy in k_r with a diffusion rate $D_{k_r} = \sum_q (k_{\theta} \tilde{V}_{\theta})^2 \tau_{corr}$, where q represents the radial wave number of the fluctuating shear flow, $\tilde{V}_{\theta}$ represents the amplitude of the fluctuating shear flow, and k_{θ} (k_r) is the azimuthal (radial) wave number of the turbulent fluctuations. The change in mean-squared radial wave number then increases in time according to the relation $\langle \delta k_r^2 \rangle = D_{k_r} t$ for times of the order of the turbulence decorrelation time. As a result the fluctuating shear flow can then be expected to have a significant impact on the turbulence when $(D_{k_r} \tau_{corr}) / k_{r0}^2 > 1$ where k_{r0}^{-1} represents the radial correlation length of the turbulence in the absence of flow shear. This inequality provides the criteria that must be satisfied in order for the fluctuating shear flow to have a significant impact on the time scale of the turbulent fluctuation lifetime. Using the definitions above, the minimum random shear flow strength requirement necessary to decorrelate turbulent fluctuations becomes $[\sum_q (k_{\theta} q V_{\theta})^2 \tau_{corr}] / k_{r0}^2 > 1$. The fluctuating shear flow radial wave number was estimated earlier to lie in the range $q = 2-3 \text{ cm}^{-1}$. Using this value we then estimate that the fluctuating shear flow amplitude lies in the range $\tilde{V}_{\theta} \approx (q \tilde{\phi}_{m=0}^{\max}) / B \approx 6-9 \times 10^4 \text{ cm/s}$ at $r = 4-5 \text{ cm}$. If the un-sheared turbulence is isotropic we have $k_{\theta} = k_{r0}$ and thus the inequality above becomes roughly $q V_{\theta} \tau_{corr} > 1$. Using the previously determined value of $\nu_{corr} = \tau_{corr}^{-1} = 6 \times 10^4 \text{ s}^{-1}$ we estimate that $q V_{\theta} \tau_{corr} \approx 2-4$. Thus we conclude that, if the assumption of isotropic turbulence is not violated, then the fluctuating shear flow observed here should be strong enough to influence the turbulence dynamics.

Absent any external momentum input, the existence of a mean and fluctuating shear flow in the presence of flow damping implies that a radial transport of turbulent angular momentum via a turbulent Reynolds stress must occur. In the presence of a turbulent radial particle flux $\tilde{\Gamma}_r$, mean shear flow $\langle v_{\theta} \rangle$, finite turbulent Reynolds stress $\langle \tilde{v}_r \tilde{v}_{\theta} \rangle$, constant pressure on a magnetic flux surface, and flow damping due to ion-neutral collisions (with damping rate ν_{i0}) the time-stationary azimuthal momentum balance fluid equation averaged over timescales longer than the turbulent correlation times and magnetic flux surfaces (denoted by $\langle \cdots \rangle$) can be written as

$$n_0 \frac{\partial \langle \tilde{v}_r \tilde{v}_{\theta} \rangle}{\partial r} + \frac{\partial (\langle \tilde{\Gamma}_r \rangle \langle v_{\theta} \rangle)}{\partial r} = -\nu_{i0} n_0 \langle v_{\theta} \rangle.$$

Taking the radial scale $\partial / \partial r \approx 1 / L_r$ we can then estimate the peak magnitude of the mean shear flow as $|\bar{v}_{\theta}| = n_0 \langle \tilde{v}_r \tilde{v}_{\theta} \rangle / [\tilde{\Gamma}_r + L_r \nu_{i0} n_0]$. The Reynolds stress can be approximated as $\langle \tilde{v}_r \tilde{v}_{\theta} \rangle \approx \gamma_{v_R v_{\theta}} |k_{\perp}^2 \phi_k^2 / B^2|$, where $\gamma_{v_R v_{\theta}}$ represents the cross-coherence between radial and azimuthal velocity fluctuations and again we have invoked the approximation that the turbulent fluctuations are roughly isotropic. We then estimate from the measured values reported earlier that the Reynolds stress can range between $\langle \tilde{v}_r \tilde{v}_{\theta} \rangle$

$\approx 0-3 \times 10^{10} \text{ cm}^2/\text{s}^2$ for our measured fluctuation amplitude and for $\gamma_{v_R v_\theta}$ ranging between the allow values $0 < \gamma_{v_R v_\theta} < 1$. At $r=3 \text{ cm}$ (where the mean flow is largest) the results in Fig. 1 show that the average plasma density is 10^{13} cm^{-3} , the radial scale length of the gradient in turbulent flux and in azimuthal flow velocity $L_r \sim 1 \text{ cm}$, the ion-neutral collision frequency $\nu_{i0} \approx 2-4 \times 10^4 \text{ s}^{-1}$, and the particle flux $\Gamma_r = 5 \times 10^{16} \text{ cm}^{-2} \text{ s}^{-1}$. Using these values we estimate the peak azimuthal flow velocity that could be driven by the observed fluctuation levels as $|\bar{V}_\theta| = 0.4-1 \times 10^6 \text{ cm/s}$ for $0.4 < \gamma_{v_R v_\theta} < 0.8$. This velocity corresponds to a radial electric field of about $0.4-1 \text{ V/cm}$, close the value deduced from the data shown in Figs. 1(c) and 1(d). Thus the magnitude of the shear flow is consistent with the value of flow that could, in principle, be driven by the turbulent Reynolds stress with reasonable assumptions about the magnitude of $\gamma_{v_R v_\theta}$.

IV. CONCLUSIONS

In summary, we report evidence for the formation of a velocity shear layer that exists at the edge of a region of broadband collisional drift turbulence in a cylindrical helicon plasma device. This velocity shear layer consists of an azimuthally symmetric radially sheared $E \times B$ flow moving in the azimuthal direction that has both a time-stationary and a time-varying component. The shear layer also appears to sustain low-order azimuthal mode number coherent fluctuations that coexist with the broadband fluctuations. Turbulent mixing across the shear layer is inhibited, and the steepest plasma density gradients are observed just inside the shear layer. A small but nonzero phase coherence between the low-frequency fluctuating shear flow, coherent structures, and broadband fluctuations is observed, indicating that energy is exchanged nonlinearly between these three types of fluctuations. Both density and potential k_θ spectra were measured, and indicate the presence of significant fluctuation energy in the $k_{\theta s} \ll 1$ region of the spectrum. A linear stability analysis of the plasma equilibrium indicates that this portion of the spectrum is linearly stable. Taken together, the finite phase coherence, measured k spectra, and linear stability results thus suggest that a nonlinear transfer of fluctuation energy from the linearly unstable $k_{\theta s} \sim 1$ wave number region into the larger scaled $k_{\theta s} \ll 1$ shear flow occurs. The time-averaged potential profile corresponding to the mean shear flow is in excellent agreement with theoretical predictions for a cylindrical plasma device, and the amplitudes of both the mean and fluctuating shear flows are both estimated to be strong enough to influence the turbulence dynamics. The magnitude of the mean and fluctuating shear flow is large enough to affect the turbulent decorrelation process, and is also consistent with estimates of the turbulent Reynolds stress. Taken as a whole, the results are consistent with the theoretical expectations of low-frequency shear flows generated via the turbulent Reynolds stress, which then, in turn, reduce or eliminate turbulent transport across the shear layer.

Several follow-on experiments are suggested by these results, and preparations for these studies are underway. First, the radial scale length of the turbulent fluctuations will

be measured to validate the assumptions used in estimating the relative shear strength. Second, a direct measurement of the turbulent Reynolds stress will be undertaken, and the nonlinear energy transfer between different spatial and temporal scales will be measured. Both of these tasks require the construction of probe arrays capable of providing the correct spatiotemporal turbulence data. Finally, numerical simulations of the experiment are underway using the Hasegawa-Wakatani theoretical model of collisional drift turbulence, suitably modified to account for ion-neutral drag, and will be compared with these experimental results to provide a more complete picture of shear flow formation in this simple physical system.

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