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LETTERS

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Influence of sheared poloidal rotation on edge turbulence

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The impact of radially sheared poloidal flows on ambient edge turbulence in tokamaks is investigated analytically. In the regime where poloidal shearing exceeds turbulent radial scattering, a hybrid time scale weighted toward the former is found to govern the decorrelation process. The coupling between radial and poloidal decorrelation results in a suppression of the turbulence below its ambient value. The turbulence quench mechanism is found to be insensitive to the sign of either the radial electric field or its shear.

It is by now generally accepted that the most efficient way to enhance the prospects for ignition in fusion devices is to improve the quality of plasma confinement in auxiliary-heated discharges. To this end, a number of confinement regimes, enhanced beyond the so-called L regime of degraded confinement, have now been experimentally identified and phenomenologically characterized. While theoretical predictions of confinement characteristics in the improved confinement regimes have been borne out to a fair extent, theoretical insight into the mechanism of the transition from the state of degraded to enhanced confinement has so far proven elusive. The establishment of poloidal flows, possibly a result of radial electric fields in response to non(intrinsic) ambipolarity, appears to be a common enough feature in a variety of auxiliary-heated scenarios associated with enhanced confinement regimes to warrant special scrutiny.

Indeed, interest in the possible role of a radial electric field in triggering the H mode has been sparked recently by the observation on DIII-D of a change in the poloidal rotation velocity about the time of the L to H transition, from which it was inferred that the radial electric field had become more negative.[1] In important, recent work, Shaing *et al.*,[2] have proposed a theory that appears to be consistent with the DIII-D observation (also see the work of Itoh *et al.*[3]). The novel observation made in this work is that the presence of a *constant* radial electric field would, in addition to the usual Doppler-shifted center-of-mass $\mathbf{E}\times\mathbf{B}$ rotation (which, by Galilean invariance, does not affect the dynamics of the turbulence), give rise to differential rotation as a result of geometric effects associated with curvature. By studying the evolution and saturation of resistivity-gradient-driven turbulence under the influence of a radial electric field in cylindrical geometry, Shaing and co-workers were led to (i) conclude that the turbulent decorrelation time was not modified to leading order from its ambient value for $\omega_s \lesssim \Delta\omega_t$, where $\Delta\omega_t$ is the decorrelation time associated with turbulent radial diffusion of fluctuations by the ambient turbulence and ω_s is the shearing frequency, and (ii) ascribe the reduction in the saturation level of the turbulence to the sign of the electric field that enters because of diamagnetic effects. In the present Letter, we reconsider the same question, focusing instead on the *shear* in the electric field as the dominant decorrelation mechanism (although our results do not depend on this feature in any way), and come to different conclusions about the physics of the decorrelation process. In particular, we show analytically that the *coupling* between poloidal shearing and turbulent radial scattering of fluctuations by itself can account for the quenching of the turbulence without the need to invoke *additional* model-dependent symmetry-breaking effects such as diamagnetic rotation. Our results, which are general in the sense that they do not assume any particular model for the background turbulence, suggest that if a sufficiently strong shear in the radial electric field (of either sign) can be built up in response, say, to nonambipolarity instigated by the interaction of an external agent with the plasma, then this will act to suppress the turbulence. Significantly, *the turbulence quench mechanism is insensitive to the sign of either the radial electric field, or its shear.* Before proceeding with the analysis, it should be pointed out that both our work and that of Shaing *et al.*, have, as their antecedent, the work of Chieuh *et al.*,[4] which represents the first attempt to consider the influence of the radial electric field on plasma turbulence (collisional drift waves in that instance). The recent works differ from the earlier reference in that they both consider perturbations that are flutelike as opposed to ballooning in character, so that the details of the underlying physical decorrelation process are different. However, the basic notion of a competition between the rates of radial scattering and poloidal shearing is the same in both cases.

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We adopt as our working prototype a generic fluid model in which fluctuations dynamically evolve according to

$$\left(\frac{\partial}{\partial t} + (\mathbf{v}_0 + \tilde{\mathbf{v}})\cdot\nabla + \mathscr{L}_d\right)\tilde{\xi} = \tilde{S}, \qquad (1)$$

where $\tilde{\xi}$ is the fluctuating field (e.g., $\tilde{T}$ in the case of ion temperature-gradient-driven turbulence, or $\tilde{p}$ in the case of resistive pressure-gradient-driven turbulence). In Eq. (1), $\mathbf{v}_0 = v_\vartheta(r)$ is taken to be an equilibrium poloidal flow, which, for concreteness, we imagine as coming from $\mathbf{E}_0 \times \mathbf{B}$ rotation, i.e., $v_\vartheta(r) = -cE_r(r)/B$. Here $\mathbf{v}\cdot\nabla$ is the advective nonlinearity, which can be either electrostatic (perpendicular $E\times B$ advection) or magnetostatic (parallel flow along stochastic field lines). In Eq. (1) $\tilde{S}$ represents a source of free energy driving the turbulence, and $\mathscr{L}_d$ is an operator responsible for dissipation of that energy. In order to focus on the physics of the decorrelation process, we need to construct an equation describing the evolution of the two-point correlation function. Assuming flutelike fluctuations, the two-point evolution equation can be cast, after standard diffusive renormalization,[5] in the form

$$\left[\partial_t + \left(v'_\vartheta - \frac{v_\vartheta}{r_+}\right) r_- \partial_{y_-} - \partial_{r_-}\mathscr{D}_-(r_-,y_-)\partial_{r_-} + \mathscr{L}_d\right]\langle\tilde{\xi}(1)\tilde{\xi}(2)\rangle = \mathscr{S}, \qquad (2)$$

where a transformation has been effected to relative ($-$) and center-of-mass ($+$) coordinates: $(r_\pm, \vartheta_\pm) = (r_1 \pm r_2, \vartheta_1 \pm \vartheta_2)$ (the indices "1" and "2" refer to the fluid coordinates at two different points) in order to best manifest the physics of decorrelation, $y_- = r_+\vartheta_-$, and $\mathscr{D}(r_-,y_-)$ characterizes the predominantly radial turbulent diffusion of fluctuations. There are two physical processes, represented by the second and third terms of Eq. (2), which physically characterize the fissure of a fluid element: poloidal decorrelation resulting from rotational shear and radial decorrelation resulting from turbulent scattering. Each of these processes can be associated with a rate: $\omega_s = (k_{0y}\,\Delta r_t)|v'_\vartheta - v_\vartheta/r_+|$, which can be thought of as the rate at which two fluid elements separated radially by Δr_t become separated poloidally by k_{0y}^{-1}, and $\Delta\omega_t = 4D/\Delta r_t^2$, which is the random, diffusive scattering rate of the ambient turbulence (Δr_t and k_{0y}^{-1} characterize the spatial correlation lengths of the ambient turbulence in the radial and poloidal direction, respectively). That ω_s has the form shown is manifested by the fact that there can be no rotational shearing of the fluctuations if the plasma were rotating as a rigid body. Thus, even in the absence of flow shear, i.e., $v'_\vartheta(r) = 0$, fluctuations can still be shorn apart as a result of curvature effects associated with curvilinear geometry. As a preliminary observation, however, note that if the shear in the poloidal flow is steeper than the radius of the minor cross section, i.e., $L_v = |d \ln v_\vartheta/dr|^{-1} < r_+ \sim a$, then it is the former that dominates rotational shearing. To our knowledge, the only available measurement of the shear flow profile on the edge of tokamaks is that of TEXT,[6] where it was observed that $L_v \sim 1$ cm $\ll a = 26$ cm. Indeed, if the $\mathbf{E}\times\mathbf{B}$-induced poloidal rotation is built up because of nonambipolarity, then intuition from neoclassical theory would suggest that the scale size of the radial electric field would not be much greater than a poloidal ion Larmor radius, which for the well-confined orbits that is desired in tokamaks, is much less than the minor radius. Consequently, it appears that *in most cases the shear flow term will dominate*, and so without loss of generality we focus on this contribution to ω_s for the balance of this Letter. More importantly, it is neither turbulent radial scattering nor poloidal shearing that determines the physical decorrelation process, but rather a hybrid of the two weighted toward the latter. Physically, this arises because the effect of poloidal flow is to enhance poloidal decorrelation by *coupling* radial scattering to sheared poloidal streaming. This has important consequences for the spatial correlation length, and consequently on the turbulent diffusion level, as we will now proceed to show.

The decorrelation rates are obtained by taking various moments of the lhs of Eq. (2),[5] which yields differential equations governing the stochastic evolution of fluid blob positions. Neglecting the dissipation operator, which does not participate in the physics of the decorrelation beyond smoothing out the peaking of the two-point correlation function at very short separation,[7] the coupled set of differential equations becomes

$$\partial_t\langle r_-^2\rangle = \Delta\omega_t[3\langle r_-^2\rangle + (k_{0y}\,\Delta r_t)^2\langle y_-^2\rangle], \qquad (3a)$$

$$\partial_t\langle y_-^2\rangle = 2(k_{0y}\,\Delta r_t)^{-1}\omega_s\langle r_-y_-\rangle, \qquad (3b)$$

$$\partial_t\langle r_-y_-\rangle = (k_{0y}\,\Delta r_t)^{-1}\omega_s\langle r_-^2\rangle + \Delta\omega_t\langle r_-y_-\rangle. \qquad (3c)$$

Here $\langle\cdots\rangle = \int dr_-\,dy_-(\cdots)G(r_-y_-)$, where $G(r_-,y_-)$ is the two-point Green's function. In the limit of negligible rotational shear, i.e., $\omega_s \ll \Delta\omega_t$, the equations decouple and one recovers the expected characteristics associated with ambient turbulence, i.e., decorrelation time on the order of $\Delta\omega_t^{-1}$, and decorrelation scales on the order of $(\Delta r_t, k_{0y}^{-1})$. The regime of interest here is the opposite limit of strong shear: $\omega_s > \Delta\omega_t$. Then there is strong coupling between radial and poloidal decorrelation, and the rates are determined by

$$\partial_{ttt}\langle r_-^2\rangle - 2\omega_s^2\,\Delta\omega_t\langle r_-^2\rangle = 0, \qquad (4)$$

with initial conditions $\langle r_-^2\rangle|_{t=0} = \langle r_-^2\rangle$, $\partial_t\langle r_-^2\rangle|_{t=0} = \Delta\omega_t[3r_-^2 + (k_{0y}\,\Delta r_t)^2 y_-^2]$, and $\partial_{tt}\langle r_-^2\rangle|_{t=0} = 3\,\Delta\omega_t^2[3r_-^2 + (k_{0y}\,\Delta r_t)^2 y_-^2] - 6\omega_s\,\Delta\omega_t(k_{0y}\,\Delta r_t)^{-1}r_-y_-$, and

$$\partial_{ttt}\langle y_-^2\rangle - 2\omega_s^2\,\Delta\omega_t\langle y_-^2\rangle = 6\omega_s^2\,\Delta\omega_t(k_{0y}\,\Delta r_t)^{-2}\langle r_-^2\rangle, \qquad (5)$$

subject to initial conditions $\langle y_-^2\rangle|_{t=0} = y_-^2$, $\partial_t\langle y_-^2\rangle|_{t=0} = 2\omega_s(k_{0y}\,\Delta r_t)^{-1}r_-y_-$, and $\partial_{tt}\langle y_-^2\rangle|_{t=0} = 2\omega_s^2(k_{0y}\,\Delta r_t)^{-2}r_-^2$. As is clear from Eqs. (4) and (5), the characteristic spectral decorrelation time is a hybrid of poloidal rotational shear and turbulent radial diffusion, i.e., $\tau_c = (2\omega_s^2\,\Delta\omega_t)^{-1/3}$. Note that (i) rotational shear weighs in more heavily than radial diffusion in detuning the fluctuations and (ii) the sign of ω_s is irrelevant. This feature is entirely analogous to the situation encountered in the nonlinear dynamics of the universal instability,[8] and the stochastic divergence of magnetic field lines.[9] The dominant solutions for $t < \tau_c$ are given by

$$\langle r_-^2 \rangle \simeq \tfrac{1}{3}\big[\,(1 + 3\sqrt[3]{2}\alpha^2) r_-^2 + \sqrt[3]{2}\alpha^2 (k_{0y}\,\Delta r_t)^2 y_-^2 - (6/\sqrt[3]{2}) k_{0y}\,\Delta r_t r_- y_-\,\big] \exp(t/\tau_c), \tag{6a}$$

$$\langle y_-^2 \rangle \simeq \tfrac{1}{3}\big[\,(2/\alpha^2)(k_{0y}\,\Delta r_t)^{-2} r_-^2 - (2/\alpha)(k_{0y}\,\Delta r_t)^{-1} r_- y_- + y_-^2\,\big] \exp(t/\tau_c), \tag{6b}$$

where $\alpha = (\Delta\omega_t/2\omega_s)^{1/3} < 1$. The rate at which a fluid blob fissures, τ_f^{-1}, is defined as the rate at which the relative separation between two adjacent fluid elements becomes comparable to fluid blob dimensions. This yields

$$\tau_f \simeq \tau_c \ln\left[\left(\frac{r_-}{\alpha\,\Delta r_t}\right)^2 - \frac{2}{3}(1 + \sqrt[3]{2})\,\frac{r_-}{\alpha\,\Delta r_t} k_{0y} y_- + \frac{1 + 6\sqrt[3]{2}}{3}(k_{0y} y_-)^2\right]^{-1}. \tag{7}$$

The characteristic spatial correlation lengths in the limit $\omega_s > \Delta\omega_t$ are then given by $(\Delta r_c, \Delta y_c) \simeq [\,(\Delta\omega_t/\omega_s)^{1/3}\Delta r_t, k_{0y}^{-1}\,]$. Note that the radial correlation length has been *reduced* by rotational shearing relative to its value as determined by ambient turbulence alone. It is useful to rewrite the turbulence suppression criterion, $\Delta\omega_t < (\omega_s^2\,\Delta\omega_t)^{1/3} \equiv \tau_c^{-1}$, in terms of radial scales, which yields $\Delta r_t/L_v > \Delta\omega_t/\omega_\vartheta$, where $\omega_\vartheta = k_{0y} v_\vartheta$ is the poloidal rotation frequency. Since all the quantities appearing above are, at least in principle, measurable experimentally, the expression provides a simple test of the importance of rotational shearing.

The reduction in the radial correlation length results in the quenching of fluctuations and a reduction in the transport level, as we now demonstrate with a simple dynamical model. The dynamic evolution of fluctuations is characterized by two competing processes represented by the left- and right-hand sides of Eq. (2). Fluctuations are driven by the free energy associated with the relaxation of macroscopic field gradient L_ξ and decay by turbulent decorrelation, characterized by τ_f [Eq. (7)]. Heuristically, then, we have $(\partial_t + \tau_f^{-1})\langle|\tilde{\xi}/\xi_0|^2\rangle \simeq D/L_\xi^2$. In the absence of poloidal shearing, $\tau_f^{-1} \sim D/\Delta r_t^2$. It follows that for stationary turbulence, $\langle|\tilde{\xi}/\xi_0|^2\rangle_{\omega_s=0} \simeq (\Delta r_t/L_\xi)^2$, which is the standard mixing length result. For strong poloidal shearing, on the other hand, $\tau_f^{-1} \simeq (\omega_s^2\,\Delta\omega_t)^{1/3}$, from which follows $\langle|\tilde{\xi}/\xi_0|^2\rangle_{\omega_s} \simeq (\Delta r_c/L_\xi)^2$, where Δr_c is given below Eq. (7). Thus we have

$$\langle|\tilde{\xi}/\xi_0|^2\rangle_{\omega_s} / \langle|\tilde{\xi}/\xi_0|^2\rangle_{\omega_s=0} \sim (\Delta\omega_t/\omega_s)^{2/3} < 1, \tag{8}$$

i.e., shear decorrelation reduces the fluctuation level from its ambient value. Moreover, since the turbulent diffusion coefficient is given by $D = \Sigma_{\mathbf{k}} F[\langle|\tilde{\xi}/\xi_0|^2\rangle_{\mathbf{k}}]\tau_f$, where F is a functional, the transport level is also reduced. The Fourier transform of Eq. (7) provides information about the spectral signature of the fluctuations for large k. Analysis shows that the spectral fall-off rates are not affected by the inclusion of rotational shearing; however, the spectral amplitude level, $\langle|\tilde{\xi}/\xi_0|^2\rangle_{\mathbf{k}}$ drops by an amount $\omega_s/\Delta\omega_t > 1$.

Finally, the extent to which a strong shear in the poloidal flow, such as what we envision here, would drive a Kelvin–Helmholtz type instability needs to be considered. The mechanism by which shear flow can destabilize fluctuations depends crucially on how efficiently the vorticity maximum (hence the free energy source) is localized relative to the region of minimum dissipation, i.e., a $\mathbf{k}\cdot\mathbf{B} = 0$ surface.[4] In the case of modes with a definite parity, the relevant comparison is between the width of the shear flow layer and the resistive layer [i.e., L_v vs $L_J = (\eta\,\Delta\omega_t\,L_s^2 c^2/k_{0y}^2 v_A^2)^{1/4}$, where $L_s = qR/\hat{s}$ is the magnetic shear length, $\hat{s} = rq'/q$, η is the resistivity, and $v_A = B/(4\pi\rho)^{1/2}$ is the Alfvén speed]. In general, $L_v > L_J$ for physically meaningful parameters, and thus a Kelvin–Helmholtz instability is unlikely.

To summarize, we have shown that the establishment of rotational shear has a quenching influence on the ambient turbulence. Of the two processes contributing to poloidal shearing, i.e., flow shear and geometric effects associated with curvature, the former is argued to be the dominant one in most cases of interest, although the results of this Letter do not depend on this point in any way. Coupling between poloidal shearing and turbulent radial scattering of fluctuations is shown to result in a hybrid decorrelation process that is weighted more strongly toward the former. Finally, the turbulence quench mechanism discussed is insensitive to the sign of either the radial electric field or its shear. An interesting consequence of the theory is that is may explain the scaling of the H-mode power threshold with magnetic field, as observed on many tokamaks. Since the radial correlation length for a variety of turbulence models scales inversely with magnetic field, a higher magnetic field implies that it would be more difficult to satisfy the turbulence suppression criterion. Finally, with respect to the relevance of the presence theory to various enhanced confinement scenarios, it must be stressed that we are not advocating this work as an explanation for the transition mechanism *per se* to enhanced confinement. What we have addressed is the manner in which poloidal flow, once established, affects the ambient turbulence.

Note: After the submission of this work for publication, we were made aware of a new calculation by Shaing, dated 8 June 1989,[10] which now appears to arrive at conclusions similar to ours.

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Measurements of classical deceleration of beam ions in the DIII-D tokamak

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Studies of the slowing down of short pulses of deuterium beam ions in the DIII-D tokamak [Nucl. Fusion **28**, 1897 (1988)] are extended to the regime where ion drag dominates electron drag. The deceleration of the beam ions is consistent with Sivukhin's classical theory [*Reviews of Plasma Physics* (Consultants Bureau, New York, 1966), Vol. 4, p. 93] to within 25%.

The deceleration of test particles through Coulomb collisions is a fundamental process in plasma physics. For such a fundamental topic, there have been relatively few direct experimental tests of the theory.[1] In tokamak research, energy transfer is almost invariably assumed to be classical. Transport coefficients inferred from neutral-beam heating experiments are determined by assuming that the beam power is partitioned classically between electrons and ions. Conceivably, however, other processes such as enhanced scattering associated with plasma turbulence[2] or the excitation of beam-driven instabilities[3] could lead to enhanced deceleration rates and might partition the energy differently between electrons and ions. The goal of the present research is to confirm that classical drag dominates beam deceleration in a tokamak plasma when the density of fast ions is well below the threshold for collective instability.

The experiment reported here is an extension of an earlier study.[4] Both studies use the same diagnostic technique. Short (2 msec) pulses of deuterium neutral beams are injected into deuterium plasmas in the DIII-D tokamak.[5] After ionization, the beam ions produce 2.5 MeV neutrons in beam–plasma fusion reactions. Because the $d(d,n)$ ^{3}He fusion cross section falls rapidly with decreasing energy, the neutron emission decays as the beam ions decelerate so the decay time of the neutron emission is simply related to the beam-ion slowing-down time. In the previous study, the neutron decay time τ_n was found to agree with classical theory (to within 30%) for an order of magnitude variation in τ_n.[4] The study was restricted to the regime $E_b > E_{\rm crit}$, where E_b is the beam energy and $E_{\rm crit}$ is the critical energy at which ion drag equals electron drag. The present work extends the previous study into the regime where beam ions slow down primarily upon thermal ions ($E_b < E_{\rm crit}$). In both experiments, the beam velocity is intermediate between the ion and electron thermal velocities ($v_e \gg v_b \gg v_i$). For comparison with classical theory, short beam pulses are superior to neutron measurements following beam heating,[6,7] because the beam ions are nearly monoenergetic and because the background plasma parameters are barely perturbed by the beam pulse.

The regime $E_b < E_{\rm crit}$ was accessed by reducing the beam voltage from 76 to 39 kV and by employing electron cyclotron heating (ECH) to raise the electron temperature (which raises $E_{\rm crit}$). The neutral-beam accelerating voltage rose in approximately 0.15 msec (to within 2 kV of its final value) and turned off following the 2 msec pulse in <0.05 msec. Approximately 1.5 MW were injected at 76 kV and 0.4 MW at 39 kV. Inside launch ECH with primarily extraordinary mode polarization (>1 MW for 500 msec) was used to heat the electrons. The plasma density ($\bar{n}_e = 1.1$–4.5×10^{13} cm^{-3}), toroidal field ($B_t = 1.7$–2.0 T), and plasma current ($I_p = 0.6$–1.2 MA) were selected for good ECH heating. The ionic composition was primarily deuterium (spectroscopic measurements found $<1\%$ hydrogen) but $Z_{\rm eff}$ (from visible bremsstrahlung) was 2 to 4 (presumably, principally carbon impurity). (Theoretically, τ_n is practically independent of $Z_{\rm eff}$ for fully stripped impurities in deuterium.) The plasma was run in the single-null divertor configuration. The electron temperature (0.9–2.5 keV) was measured using the electron cyclotron emission (ECE) profile diagnostic[8] corroborated by Thomson scattering.[9] Within $\sim 15\%$ errors, the measurements were consistent. Electron density was measured by four CO_2 interferometer chords. The Thomson scattering density measurements were typically 20% lower than the density inferred from the CO_2 measurements. The perturbations in T_e and $\bar{n}_e$ associated with the beam pulse were undetectable. The fast-ion density was well below the thresholds for beam driven instabilities ($n_b/n_e < 0.1\%$). The neutron emission was measured using an uncollimated plastic scintillator with a temporal resolution of <0.1 msec.[6,10]