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Experimental Signatures of Critically Balanced Turbulence in MAST

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Beam emission spectroscopy (BES) measurements of ion-scale density fluctuations in the MAST tokamak are used to show that the turbulence correlation time, the drift time associated with ion temperature or density gradients, the particle (ion) streaming time along the magnetic field, and the magnetic drift time are consistently comparable, suggesting a “critically balanced” turbulence determined by the local equilibrium. The resulting scalings of the poloidal and radial correlation lengths are derived and tested. The nonlinear time inferred from the density fluctuations is longer than the other times; its ratio to the correlation time scales as $\nu_{*i}^{-0.8 \pm 0.1}$, where ν_{*i} = ion collision rate/streaming rate. This is consistent with turbulent decorrelation being controlled by a zonal component, invisible to the BES, with an amplitude exceeding those of the drift waves by $\sim \nu_{*i}^{-0.8}$.

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Introduction.—Microscale turbulence hindering energy confinement in magnetically confined plasmas is driven by gradients of equilibrium quantities such as temperature and density. These gradients give rise to instabilities that inject energy into fluctuations (“drift waves”) at scales just above the ion gyroscale. The most effective of these is believed to be the ion-temperature-gradient (ITG) instability [1–3]. A turbulent state ensues, giving rise to the “anomalous transport” of energy [4]. It is of both practical and fundamental interest what the structure of this turbulence is and how it and the resulting transport depend on the local equilibrium parameters.

Fluctuations in a magnetized toroidal plasma are subject to a number of physical effects, which can be classified in terms of various time scales: the drift times associated with the temperature and density gradients, the particle streaming time along the magnetic field as it takes them around the torus toroidally and poloidally, the magnetic drift times of particles moving across the field, the nonlinear time of the fluctuations being advected across the field by the fluctuating $\vec{E} \times \vec{B}$ velocity, the collision time, and the shear time of the plasma rotation. There has been a growing understanding [5], driven largely by theory [6–9], observations [10–12], and simulations of magnetohydrodynamic [13–15] and kinetic [7,16] plasma turbulence in space, that, if a medium can support parallel (to the magnetic field) propagation of waves (and/or particles) and nonlinear interactions in the perpendicular direction, the turbulence in such a medium would normally be “critically

balanced,” meaning that the characteristic time scales of propagation and nonlinear interaction would be comparable to each other and (therefore) to the correlation time of the fluctuations. This means that the turbulence is *not* weak and *not* two dimensional, unless artificially constrained [9].

In this Letter, we use beam emission spectroscopy (BES) measurements [17–20] in the MAST tokamak [21], along with the local equilibrium parameters calculated by other diagnostics, to estimate and compare the characteristic time scales of the turbulent fluctuations in the energy-containing range. We obtain, for the first time, direct evidence that the correlation, drift, and parallel streaming time scales are indeed comparable across a range of equilibrium parameters (cf. Refs. [22,23]) and that the magnetic drift time is part of this “grand critical balance” as well. We also find indirect evidence that the decorrelation rate of turbulence is controlled by a zonal component whose relative importance to the drift-wave-like fluctuations scales with the ion collisionality.

Before presenting this evidence and its implications (e.g., dependence of the correlation lengths on equilibrium parameters), let us describe how it was obtained.

Experimental data and its analysis.—The data presented here were collected from 39 neutral-beam heated “double-null-diverted” discharges (including *L*- and *H*-modes and internal transport barriers), with no pellet injection and no resonant magnetic perturbations. The BES system on MAST [21] collects photons from a 2D array of

8 radial \times 4 vertical locations in the outboard midplane of the tokamak, with a 2 cm separation between adjacent channels in either direction. The detected intensity (mean + fluctuating, $I + \delta I$) is used to infer, at each location, the density fluctuation level $\delta n/n = (1/\beta)(\delta I/I)$ [17], where β depends on the mean density n and is estimated based on the Hutchinson model [24]. (Dependence on the mean temperature is weak.) As the BES array was moved radially for different discharges, our database contains cases with radial viewing positions $10 \text{ cm} < r < 50 \text{ cm}$ from the magnetic axis (the minor radius of the plasma is $\approx 60 \text{ cm}$).

The local equilibrium parameters are measured by standard diagnostics: mean electron densities n_e and temperatures T_e by the Thomson scattering system [25], impurity ion (c^{6+}) mean temperatures (assumed to equal the bulk ion temperature T_i) and toroidal flow velocity U_ϕ by the charge exchange recombination spectroscopy system [26], and local magnetic pitch angle α by the motional stark effect (MSE) system [27]; further equilibrium magnetic field information is obtained from pressure- and MSE-constrained EFIT equilibria [28].

We filter the BES data to the frequency interval [20, 100] kHz [29] and calculate the spatiotemporal correlation function

$$\begin{aligned} \mathcal{C}(\Delta x, \Delta Z, \Delta t) &= \frac{\langle \delta I(x, Z, t) \delta I(x + \Delta x, Z + \Delta Z, t + \Delta t) \rangle}{\sqrt{\langle \delta I^2(x, Z, t) \rangle \langle \delta I^2(x + \Delta x, Z + \Delta Z, t + \Delta t) \rangle}}, \quad (1) \end{aligned}$$

where x , Z , and t are the radial, vertical, and time coordinates, respectively, and Δx , ΔZ , and Δt are the corresponding channel separations and the time lag; the angled brackets represent the time average over 5 ms periods. At $\Delta x = \Delta Z = 0$, the autocovariances $\langle \delta I(x, Z, t) \delta I(x, Z, t + \Delta t) \rangle$ contain not only the physical signal but also photon and electronic noise. We remove this effect by applying LED light to the BES channels, obtaining 150 different dc levels of BES signal from 0 to 1.5 V, calculating the noise autocovariance $C_N(\Delta t)$ at each dc level with the same band frequency filter of [20, 100] kHz, and then finding $C_N(\Delta t)$, whose dc level of the signal matches the dc level of the BES data from the MAST discharges, and subtracting it from the calculated autocovariances. From the correlation function (1) (illustrated in Fig. 1), we calculate the local characteristics of the density fluctuations.

The fluctuation level at each radial location is obtained from the (noise-subtracted) autocovariance function $\delta n/n = (1/\beta) \sqrt{\langle \delta I^2(x, Z, t) \rangle} / I$ at all 32 locations and then averaged over the four poloidally separated channels at the same radial location.

The correlation length ℓ_y , perpendicular to the magnetic field within a flux surface is obtained from the vertical (poloidal) correlation length ℓ_Z via $\ell_y = \ell_Z \cos \alpha$, assuming that the parallel correlation length is sufficiently long:

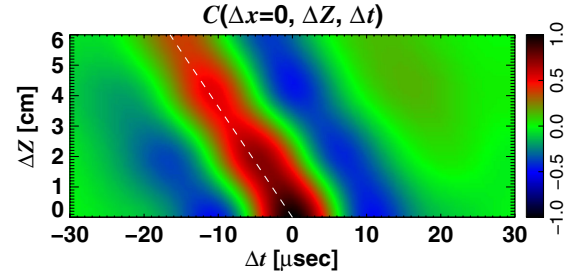


FIG. 1 (color online). An example of the correlation function in the poloidal-temporal plane $\mathcal{C}(\Delta x = 0, \Delta Z, \Delta t)$. These data were taken at $r = 30 \text{ cm}$; the toroidal rotation speed was $U_\phi = 10 \text{ km/s}$, and magnetic pitch angle was $\alpha = 20^\circ$. The direction of maximum correlation is the direction of the magnetic field (dashed line), as inferred from the MSE [27] by assuming that the toroidal coordinate is $U_\phi \Delta t$ [29].

$\ell_{\parallel} \gg \ell_y \tan \alpha$. The correlation length ℓ_Z is estimated using four poloidal channels at each radial location (the top channel is the reference channel) by fitting $\mathcal{C}(\Delta x = 0, \Delta Z, \Delta t = 0)$ to the function $f_Z(\Delta Z) = p_Z + (1 - p_Z) \cos[2\pi \Delta Z / \ell_Z] \exp[-|\Delta Z| / \ell_Z]$, where p_Z is a fitting constant that serves to account for global structures such as coherent magnetohydrodynamics modes [for which $\mathcal{C}(\Delta x = 0, \Delta Z = \infty, \Delta t = 0) = p_Z \neq 0$]. In choosing $f_Z(\Delta Z)$, we assumed wavelike fluctuations in the poloidal direction [18] (drift-wave turbulence), with the wavelength and correlation length comparable to each other. It is not possible to distinguish meaningfully between the two with only four poloidal channels. Assuming wavelike structure is essential, as, in most cases, we find that $\mathcal{C}(\Delta x = 0, \Delta Z, \Delta t = 0)$ goes negative and/or is nonmonotonic over the vertical extent of the BES array.

The radial correlation length ℓ_x is estimated using eight radial channels at each poloidal location (the reference channel is the fourth from the inside). The correlation function $\mathcal{C}(\Delta x, \Delta Z = 0, \Delta t = 0)$ is fitted to the function $f_x(\Delta x) = p_x + (1 - p_x) \exp[-|\Delta x| / \ell_x]$, where p_x plays the same role as p_Z did for f_Z . The values of ℓ_x from four poloidal locations are averaged, assuming that the radial correlations do not change significantly within the poloidal extent of the BES array. Because we have to use the entire array to estimate ℓ_x , the number of data points for ℓ_x is 8 times smaller than for ℓ_y .

To estimate the correlation time τ_c , we use the fact that the fluctuating density patterns are advected poloidally past the BES array with an apparent velocity $v_{\text{BES}} = U_\phi \tan \alpha$ due to the toroidal rotation velocity U_ϕ [29]. We fit $\mathcal{C}[\Delta x = 0, \Delta Z, \Delta t = \Delta t_{\text{peak}}(\Delta Z)]$ taken at the time delay $\Delta t_{\text{peak}}(\Delta Z)$ when the correlation function is maximum at a given ΔZ [30], to the function $f_\tau(\Delta Z) = \exp[-|\Delta t_{\text{peak}}(\Delta Z)| / \tau_c]$. This method relies on the temporal decorrelation dominating over the parallel spatial decorrelation, viz., we require $\tau_c \ll \ell_{\parallel} \cos \alpha / U_\phi$. Anticipating the critical balance assumption $\tau_c \sim \ell_{\parallel} / v_{\text{th}i}$

[5], where $v_{\text{th}i} = \sqrt{2T_i/m_i}$ is the ion thermal speed, and denoting the Mach number $\text{Ma} = U_\Phi/v_{\text{th}i}$, we estimate that the fractional error in τ_c is $\sim \text{Ma}/\cos\alpha$, which was never more than 20% in the discharges we used.

The quantities $\delta n/n$, ℓ_y , ℓ_x , and τ_c are calculated every 5 ms for all 39 discharges. All fitting is done using the MPFIT procedure [31]. We consider a data point unreliable and discard it if (i) $I < 0.3$ V (the signal-to-noise ratio is too low, i.e., $\text{SNR} < 150$), (ii) the estimated correlation lengths are smaller than the size of the point spread function [32] $\ell_x < 2$ or $\ell_y < 5$ cm, (iii) the assumption that plasma rotation is mostly toroidal is suspect, viz., $|(v_{\text{BES}} - U_\Phi \tan\alpha)/v_{\text{BES}}| \geq 0.2$ (see Ref. [29]), where v_{BES} is calculated at each radial location using the cross-correlation time delay method [30], (iv) the estimated error in the calculation of v_{BES} is $>20\%$, and (v) p_z or $p_x > 0.5$. The last two exclusion criteria pick out the cases when magnetohydrodynamics modes are too strong; they are known to degrade the reliability of the BES data [29]. The remaining database contains 448 points.

Correlation time vs drift time.—The turbulence can be driven by radial gradients in the mean ion and electron temperatures $T_{i,e}$ and density n . Denoting $L_{T_{i,e}}^{-1} = |\nabla \ln T_{i,e}|$ and $L_n^{-1} = |\nabla \ln n|$, the associated time scales are the inverse drift frequencies

$$\tau_{*i,e}^{-1} = \frac{\rho_{i,e}}{\ell_y} \frac{v_{\text{th},i,e}}{L_{T_{i,e}}}, \quad \tau_{*n}^{-1} = \frac{\rho_i}{\ell_y} \frac{v_{\text{th}i}}{L_n}, \quad (2)$$

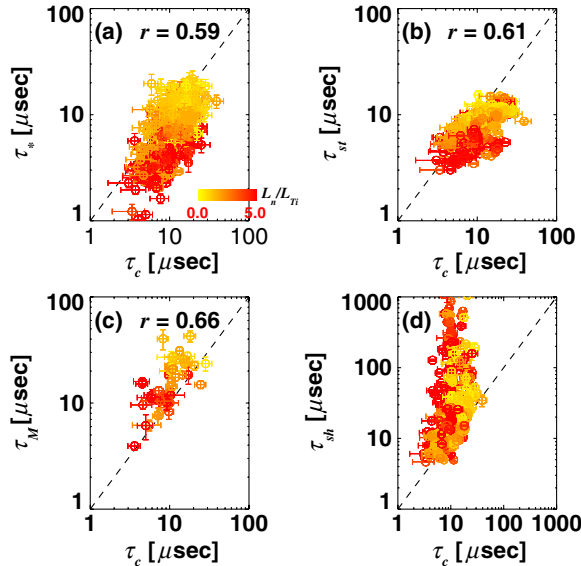


FIG. 2 (color online). (a) Drift time $\tau_* = (\ell_y/\rho_i)L_*/v_{\text{th}i}$ vs correlation time τ_c , (b) streaming time $\tau_{\text{st}} = \Lambda/v_{\text{th}i} = (B/B_p)\pi r/v_{\text{th}i}$ vs τ_c , (c) magnetic drift time $\tau_M = (\ell_x/\rho_i)R/v_{\text{th}i}$ vs τ_c , and (d) perpendicular velocity shear time $\tau_{\text{sh}} = [(B_p/B)dU_\Phi/dr]^{-1}$ vs τ_c . In all cases, the color of points represents $\eta_i = L_n/L_{T_i}$. r in (a)–(c) represents Pearson’s correlation coefficient between the logarithms of the respected quantities.

where $\rho_{i,e}$ are the ion (i) and electron (e) gyroradii and $v_{\text{th},i,e} = \sqrt{2T_{i,e}/m_{i,e}}$ the thermal speeds.

In Fig. 2(a), we compare the drift times with the correlation time τ_c . We find that $\tau_* = (0.7 \pm 0.3)\tau_c$, where $\tau_* = \min\{\tau_{*i}, \tau_{*n}\}$ and the spread is calculated as the root mean square deviation from the mean value. Thus, the turbulence appears to be driven by the larger of the ion temperature or density gradient [33]. We find no clear correlation of τ_{*e} with τ_c or with any of the other time scales discussed below.

Critical balance.—The standard argument behind the critical balance conjecture is causality [9]: Two distant points on a field line cannot stay correlated if information cannot be exchanged between them over a turbulence correlation time. Assuming information travels at $v_{\text{th}i}$ [1–5], one gets $\ell_{\parallel} \sim v_{\text{th}i}\tau_c$. This cannot be checked directly because there are no diagnostics capable of measuring ℓ_{\parallel} on MAST. Considering that the inboard side of the torus is a region of “good” (stabilizing) curvature, not much turbulence is expected there, so we assume that, at the energy injection scale, $\ell_{\parallel} \sim \Lambda$ [5], where the distance along the field line that takes a particle from the outer to the inner side of the torus is $\Lambda = \pi r B/B_p$ (r is the minor radius at the BES position on the outer side and B_p the poloidal component of the magnetic field) [34]. Then, critical balance means that τ_c should be comparable to

$$\tau_{\text{st}}^{-1} = \frac{v_{\text{th}i}}{\Lambda} = \frac{v_{\text{th}i}}{\pi r} \frac{B_p}{B} \sim \frac{v_{\text{th}i}}{\ell_{\parallel}}, \quad (3)$$

the ion streaming time (the first two equalities are its definition, the last an assumption). Indeed, we find $\tau_{\text{st}} = (0.8 \pm 0.3)\tau_c$ [see Fig. 2(b)].

The balance $\tau_{\text{st}} \sim \tau_*$ implies that the poloidal correlation scale is $\ell_y/\rho_i \sim \Lambda/L_*$, where $L_* = \min\{L_T, L_n\}$ [5]. This is tested in Fig. 3(a), showing that, while the two quantities are certainly of the same order, we do not have enough of a range of equilibrium parameters to state conclusively that this theoretically predicted scaling works.

Magnetic drift time and radial correlation scale.—The time scale of the magnetic (∇B and curvature) drifts is

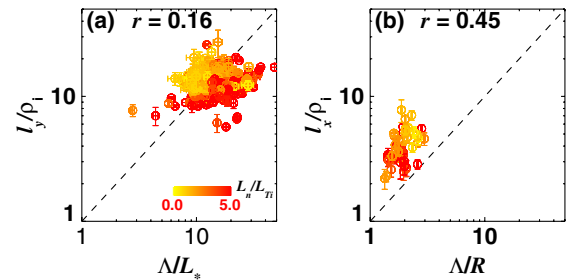


FIG. 3 (color online). (a) Poloidal correlation length ℓ_y/ρ_i vs Λ/L_* . (b) Radial correlation length ℓ_x/ρ_i vs Λ/R . The color and r are as in Fig. 2.

$$\tau_M^{-1} = \frac{\rho_i}{\ell_x} \frac{v_{thi}}{R}, \quad (4)$$

$$(\tau_{nl}^{NZ})^{-1} = \frac{v_{thi}\rho_i}{\ell_x\ell_y} \frac{T_e}{T_i} \frac{\delta n}{n}. \quad (5)$$

where we have assumed that the scale length of the background magnetic field is R (major radius at the viewing location) and $\ell_x < \ell_y$ (this will shortly prove correct). It is clear that this scale cannot be shorter than τ_c because damping due to the drift resonance would eliminate such fluctuations. While magnetic drift physics may matter (in a torus, curvature contributes to the ITG drive [4]), it does not have to affect scalings, as, e.g., it did not in the numerical simulations of [5]. In contrast, Fig. 2(c) shows that, in the MAST discharges we have analyzed, τ_M is not negligible and scales with τ_c , similarly to τ_* and τ_{st} . We find $\tau_M = (1.6 \pm 0.7)\tau_c$. Thus, a “grand critical balance” appears to hold in MAST, viz., $\tau_c \sim \tau_* \sim \tau_{st} \sim \tau_M$.

This suggests that the balance of relevant time scales determines correlation lengths of the turbulence in all three spatial directions. Indeed, balancing $\tau_M \sim \tau_{st}$, we find the radial correlation scale $\ell_x/\rho_i \sim \Lambda/R$, the scaling tested in Fig. 3(b), with a degree of success. Thus, density fluctuations in MAST are not isotropic in the perpendicular plane but elongated poloidally: $\ell_y/\ell_x \sim R/L_*$ (~ 5 in our data). Interestingly, this clashes with the approximate isotropy ($\ell_x \sim \ell_y$) reported in Cyclone Base Case simulations [5] and in measured DIII-D tokamak turbulence (where $\ell_y/\ell_x \sim 1.4$ [35] and ℓ_x appears independent of B_p [36]). Whether this is a difference between spherical and conventional tokamaks is not as yet clear (cf. Ref. [37]).

Nonlinear time.—Since we know the fluctuation amplitude, we can directly estimate the time scale associated with the advection of the fluctuations ($\delta \vec{u}_\perp \cdot \vec{\nabla} \delta n$) by the fluctuating $\vec{E} \times \vec{B}$ velocity $\delta u_\perp = c\vec{B} \times \vec{\nabla} \varphi / B^2$. The electrostatic potential φ is not directly measured but can be estimated by assuming the Boltzmann response of the electrons: $\delta n/n \approx e\varphi/T_e$. This estimate ignores trapped particles and, more importantly, as we are about to argue, also does not apply to ion-scale zonal flows (poloidally and toroidally symmetric perturbations of φ with $\delta n = 0$ [38,39]). Thus, the nonzonal nonlinear time is

Figure 4(a) shows that τ_{nl}^{NZ} is always larger than τ_c (or the other time scales discussed above) and, furthermore, observed to have an inverse rather than a direct correlation with it. Since turbulence clearly cannot be saturated by linear physics alone, this means that our estimate does not capture the correct nonlinear time. We conjecture that it is in fact the coupling to the zonal flows, invisible (directly) to BES (because their $\delta n = 0$), that dominates over the nonlinear interaction between the drift-wave-like fluctuations represented by τ_{nl}^{NZ} [38,40–46]. It has long been suspected that the relative amplitude of the zonal flows compared to that of the drift waves depends on the ion collisionality [38,47–50]. We can test this expectation by assuming that τ_c is the characteristic time associated with the coupling of the drift waves to the zonal flows and so depends on their amplitude. Figure 4(b) indeed shows a strong collisionality dependence: using the MPFIT procedure [31], $\tau_c/\tau_{nl}^{NZ} \sim \nu_{*i}^{0.8 \pm 0.1}$, where $\nu_{*i} = \nu_{ii}\tau_{st}$. (The ion collision time itself ν_{ii}^{-1} is at least an order of magnitude longer than the time scales that participate in the “grand critical balance.”) If $\tau_c^{-1} \sim (v_{thi}\rho_i/\ell_x\ell_y)e\varphi^{ZF}/T_i$, where φ^{ZF} is the amplitude of the zonal potential, this result implies that the ratio of the zonal to the nonzonal components of the turbulence is $\varphi^{ZF}/\varphi^{NZ} \sim \nu_{*i}^{-0.8 \pm 0.1}$ [51].

We note that this situation is qualitatively distinct from what is seen in numerical simulations of ITG turbulence far from the threshold [5], where the drift-wave nonlinearity appears to dominate ($\tau_{nl}^{NZ} \sim \tau_c$). However, the turbulence in a real tokamak is likely to be close to marginal and so possibly in the state of reduced transport controlled by weakly collisionally damped zonal flows [39] and usually associated with the so-called “Dimits upshift” of the stiff-transport threshold [43,44,50,52].

Discussion.—Our results support the notion that the statistics of turbulence are determined by the local equilibrium properties of the plasma [53]. We find little correlation between the quantities reported above and the radial location [54]. (Note that we have limited our consideration to temporal and spatial scales and did not touch on the fluctuation amplitudes or transport properties, which do of course depend on radius.) Our results also appeared insensitive to (i.e., not measurably correlated with) three other parameters that might in principle have proven important: T_i/T_e (varied between 0.5 and 2), the magnetic shear $\hat{s} = d\ln q/d\ln r$ (varied between -1 and 5), and the perpendicular component of the toroidal velocity shear $\tau_{sh}^{-1} = (B_p/B)dU_\Phi/dr$. In much of our data, $\tau_{sh} \geq \tau_c$, τ_{st} [see Fig. 2(d)], so it stands to reason that the statistics of the turbulence would not be dramatically affected; in the instances of $\tau_{sh} \sim \tau_{st}$, the effect of τ_{sh} could not be isolated [55]. It would be interesting to investigate higher-rotation plasmas, as τ_{sh}^{-1} , when sufficiently large, is expected to

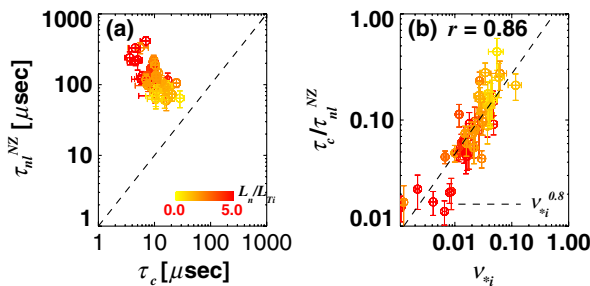


FIG. 4 (color online). (a) The nonlinear time associated with density fluctuations τ_{nl}^{NZ} vs the correlation time τ_c . (b) Their ratio vs the normalized ion collision rate $\nu_{*i} = \nu_{ii}\tau_{st}$. The color and r are as in Fig. 2.

have a dramatic effect on transport [41,56–62]; even in our database, there is in fact some evidence that velocity shear might raise the critical temperature gradients [63], but we see no signature of this effect in the correlation properties of the turbulence.

Conclusion.—We have presented experimental results that are statistically consistent with a turbulent state in MAST set by the local equilibrium and in which the time scales of the linear drive, turbulence decorrelation, ion streaming, and magnetic drifts are all similar and scale together as equilibrium parameters are varied. This “grand critical balance” implies a three-dimensionally anisotropic turbulence, with parallel, poloidal, and radial correlation lengths having different parameter dependences and $\ell_{\parallel} \gg \ell_y > \ell_x$. Our results also suggest the presence of a zonal component with an amplitude $\nu_{*i}^{-0.8 \pm 0.1}$ greater than the drift-wave density fluctuations. Note that these results are entirely consistent with a drift-wave turbulence obeying the gyrokinetic ordering [5,53,64] but provide a more detailed view of the dependence of this turbulence on the local equilibrium parameters.

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- [1] L. I. Rudakov and R. Z. Sagdeev, Dokl. Akad. Nauk SSSR **138**, 581 (1961).
- [2] B. Coppi, M. N. Rosenbluth, and R. Z. Sagdeev, *Phys. Fluids* **10**, 582 (1967).
- [3] S. C. Cowley, R. M. Kulsrud, and R. Sudan, *Phys. Fluids B* **3**, 2767 (1991).
- [4] W. Horton, *Rev. Mod. Phys.* **71**, 735 (1999).
- [5] M. Barnes, F. I. Parra, and A. A. Schekochihin, *Phys. Rev. Lett.* **107**, 115003 (2011).
- [6] P. Goldreich and S. Sridhar, *Astrophys. J.* **438**, 763 (1995).
- [7] J. Cho and A. Lazarian, *Astrophys. J.* **615**, L41 (2004).
- [8] A. A. Schekochihin, S. C. Cowley, W. Dorland, G. W. Hammett, G. G. Howes, E. Quataert, and T. Tatsuno, *Astrophys. J. Suppl. Ser.* **182**, 310 (2009).
- [9] S. V. Nazarenko and A. A. Schekochihin, *J. Fluid Mech.* **677**, 134 (2011).
- [10] T. S. Horbury, M. Forman, and S. Oughton, *Phys. Rev. Lett.* **101**, 175005 (2008).
- [11] J. J. Podesta, *Astrophys. J.* **698**, 986 (2009).
- [12] R. T. Wicks, T. S. Horbury, C. H. K. Chen, and A. A. Schekochihin, *Mon. Not. R. Astron. Soc.* **407**, L31 (2010).
- [13] J. Cho and E. T. Vishniac, *Astrophys. J.* **539**, 273 (2000).
- [14] J. Maron and P. Goldreich, *Astrophys. J.* **554**, 1175 (2001).
- [15] C. H. K. Chen, A. Mallet, T. A. Yousef, A. A. Schekochihin, and T. S. Horbury, *Mon. Not. R. Astron. Soc.* **415**, 3219 (2011).
- [16] J. M. TenBarge and G. G. Howes, *Phys. Plasmas* **19**, 055901 (2012).
- [17] R. J. Fonck, P. A. Duperrex, and S. F. Paul, *Rev. Sci. Instrum.* **61**, 3487 (1990).
- [18] R. J. Fonck, G. Cosby, R. D. Durst, S. F. Paul, N. Bretz, S. Scott, E. Synakowski, and G. Taylor, *Phys. Rev. Lett.* **70**, 3736 (1993).
- [19] G. McKee, R. Ashley, R. Durst, R. Fonck, M. Jakubowski, K. Tritz, K. Burrell, C. Greenfield, and J. Robinson, *Rev. Sci. Instrum.* **70**, 913 (1999).
- [20] G. R. McKee, C. Fenzi, R. J. Fonck, and M. Jakubowski, *Rev. Sci. Instrum.* **74**, 2014 (2003).
- [21] A. R. Field, D. Dunai, R. Gaffka, Y.-c. Ghim, I. Kiss, B. Meszaros, T. Krizanoczi, S. Shibaev, and S. Zoletnik, *Rev. Sci. Instrum.* **83**, 013508 (2012).
- [22] G. R. McKee, C. C. Petty, R. E. Waltz, C. Fenzi, R. J. Fonck, J. E. Kinsey, T. C. Luce, K. H. Burrell, D. R. Baker, E. J. Doyle, X. Garbet, R. A. Moyer, C. L. Rettig, T. L. Rhodes, D. W. Ross, G. M. Staebler, R. Sydora, and M. R. Wade, *Nucl. Fusion* **41**, 1235 (2001).
- [23] P. Hennequin, R. Sabot, C. Honore, G. T. Hoang, X. Garbet, A. Truc, C. Fenzi, and A. Quemeneur, *Plasma Phys. Controlled Fusion* **46**, B121 (2004).
- [24] I. H. Hutchinson, *Plasma Phys. Controlled Fusion* **44**, 71 (2002).
- [25] R. Scannell, M. J. Walsh, M. R. Dunstan, J. Figueiredo, G. Naylor, T. O’Gorman, S. Shibaev, K. J. Gibson, and H. Wilson, *Rev. Sci. Instrum.* **81**, 10D520 (2010).
- [26] N. J. Conway, P. G. Carolan, J. McCone, M. J. Walsh, and M. Wisse, *Rev. Sci. Instrum.* **77**, 10F131 (2006).
- [27] M. F. M. De Bock, N. J. Conway, M. J. Walsh, P. G. Carolan, and N. C. Hawkes, *Rev. Sci. Instrum.* **79**, 10F524 (2008).
- [28] L. L. Lao, H. St. John, R. D. Stambaugh, A. G. Kellman, and W. Pfeiffer, *Nucl. Fusion* **25**, 1611 (1985).
- [29] Y.-c. Ghim, A. R. Field, D. Dunai, S. Zoletnik, L. Bardoczi, A. A. Schekochihin, and the MAST Team, *Plasma Phys. Controlled Fusion* **54**, 095012 (2012).
- [30] R. D. Durst, R. J. Fonck, G. Cosby, H. Evensen, and S. F. Paul, *Rev. Sci. Instrum.* **63**, 4907 (1992).
- [31] C. B. Markwardt, in *Proceedings of the Astronomical Data Analysis Software and Systems XVIII, Quebec, Canada, 2008*, edited by D. Bohlender, P. Dowler, and D. Durand, ASP Conference Series Vol. 411 (Astronomical Society of the Pacific, San Francisco, 2009), p. 251.
- [32] Y.-c. Ghim, A. R. Field, S. Zoletnik, and D. Dunai, *Rev. Sci. Instrum.* **81**, 10D713 (2010).
- [33] However, for $\tau_c \lesssim 10 \mu \text{ sec}$, $\tau_{*e} \sim \tau_{*i}$, and, for $\tau_c \gtrsim 10 \mu \text{ sec}$, $\tau_{*n} \sim \tau_{*i}$, so we cannot rule out the ion-scale electron drive.

- [34] In a conventional tokamak, $\Lambda \approx \pi q R$, where q is the safety factor and R the major radius, but, in a spherical tokamak, our local estimate is more appropriate.
- [35] M. W. Shafer, R. J. Fonck, G. R. McKee, C. Holland, A. E. White, and D. J. Schlossberg, *Phys. Plasmas* **19**, 032504 (2012).
- [36] T. L. Rhodes, J.-N. Leboeuf, R. D. Sydora, R. J. Groebner, E. J. Doyle, G. R. McKee, W. A. Peebles, C. L. Rettig, L. Zeng, and G. Wang, *Phys. Plasmas* **9**, 2141 (2002).
- [37] S. Zweben, R. Maqueda, D. Stotler, A. Keesee, J. Boedo, C. Bush, S. Kaye, B. LeBlanc, J. Lowrance, V. Mastrocola, R. Maingi, N. Nishino, G. Renda, D. Swain, J. Wilgen, and (the NSTX Team), *Nucl. Fusion* **44**, 134 (2004).
- [38] P. H. Diamond, S.-I. Itoh, K. Itoh, and T.-S. Hahm, *Plasma Phys. Controlled Fusion* **47**, R35 (2005).
- [39] A. Fujisawa, *Nucl. Fusion* **49**, 013001 (2009).
- [40] G. W. Hammett, M. A. Beer, W. Dorland, S. C. Cowley, and S. A. Smith, *Plasma Phys. Controlled Fusion* **35**, 973 (1993).
- [41] R. E. Waltz, G. D. Kerbel, and J. Milovich, *Phys. Plasmas* **1**, 2229 (1994).
- [42] Z. Lin, T. S. Hahm, W. W. Lee, W. M. Tang, and R. B. White, *Science* **281**, 1835 (1998).
- [43] A. M. Dimits, G. Bateman, M. A. Beer, B. I. Cohen, W. Dorland, G. W. Hammett, C. Kim, J. E. Kinsey, M. Kotschenreuther, A. H. Kritz, L. L. Lao, J. Mandrekas, W. M. Nevins, S. E. Parker, A. J. Redd, D. E. Shumaker, R. Sydora, and J. Weiland, *Phys. Plasmas* **7**, 969 (2000).
- [44] B. N. Rogers, W. Dorland, and M. Kotschenreuther, *Phys. Rev. Lett.* **85**, 5336 (2000).
- [45] M. Nakata, T. H. Watanabe, and H. Sugama, *Phys. Plasmas* **19**, 022303 (2012).
- [46] K. D. Makwana, P. W. Terry, and J. H. Kim, *Phys. Plasmas* **19**, 062310 (2012).
- [47] F. L. Hinton and M. N. Rosenbluth, *Plasma Phys. Controlled Fusion* **41**, A653 (1999).
- [48] Y. Xiao, P. J. Catto, and K. Molvig, *Phys. Plasmas* **14**, 032302 (2007).
- [49] P. Ricci, B. N. Rogers, and W. Dorland, *Phys. Rev. Lett.* **97**, 245001 (2006).
- [50] Z. Lin, T. S. Hahm, W. W. Lee, W. M. Tang, and P. H. Diamond, *Phys. Rev. Lett.* **83**, 3645 (1999).
- [51] A scaling popular in theoretical models of zonal-flow ITG turbulence is $\nu_{*i}^{-1/2}$ [38].
- [52] D. R. Mikkelsen and W. Dorland, *Phys. Rev. Lett.* **101**, 135003 (2008).
- [53] I. G. Abel, G. G. Plunk, E. Wang, M. Barnes, S. C. Cowley, W. Dorland, and A. A. Schekochihin, [arXiv:1209.4782](https://arxiv.org/abs/1209.4782).
- [54] There is a slight bias in Fig. 4 for larger ν_{*i} to be found farther from the magnetic axis.
- [55] In general, we expect that a strong velocity shear would change ℓ_{\parallel} via a modified critical balance: If $\tau_{\text{sh}} < \tau_{\text{st}}$, then $\tau_c \sim \tau_{\text{sh}} \sim \ell_{\parallel} / v_{\text{th}i}$, so $\ell_{\parallel} \sim v_{\text{th}i} \tau_{\text{sh}} < \Lambda$.
- [56] P. C. de Vries, E. Joffrin, M. Brix, C. D. Challis, K. Crombé, B. Esposito, N. C. Hawkes, C. Giroud, J. Hobirk, J. Lönroth, P. Mantica, D. Strintzi, T. Tala, I. Voitsekhovitch, and (JET-EFDA) Contributors to the Work Programme, *Nucl. Fusion* **49**, 075007 (2009).
- [57] E. G. Highcock, M. Barnes, A. A. Schekochihin, F. I. Parra, C. M. Roach, and S. C. Cowley, *Phys. Rev. Lett.* **105**, 215003 (2010).
- [58] M. Barnes, F. I. Parra, E. G. Highcock, A. A. Schekochihin, S. C. Cowley, and C. M. Roach, *Phys. Rev. Lett.* **106**, 175004 (2011).
- [59] F. I. Parra, M. Barnes, E. G. Highcock, A. A. Schekochihin, and S. C. Cowley, *Phys. Rev. Lett.* **106**, 115004 (2011).
- [60] P. Mantica, C. Angioni, C. Challis, G. Colyer, L. Frassinetti, N. Hawkes, T. Johnson, M. Tsolas, P. C. de Vries, J. Weiland, B. Baiocchi, M. N. A. Beurskens, A. C. A. Figueiredo, C. Giroud, J. Hobirk, E. Joffrin, E. Lerche, V. Naulin, A. G. Peeters, A. Salmi, C. Sozzi, D. Strintzi, G. Staebler, T. Tala, D. Van Eester, and T. Versloot, *Phys. Rev. Lett.* **107**, 135004 (2011).
- [61] E. G. Highcock, A. A. Schekochihin, S. C. Cowley, M. Barnes, F. I. Parra, C. M. Roach, and W. Dorland, *Phys. Rev. Lett.* **109**, 265001 (2012).
- [62] C. M. Roach, I. G. Abel, R. J. Akers, W. Arter, M. Barnes, Y. Camenen, F. J. Casson, G. Colyer, J. W. Connor, S. C. Cowley, D. Dickinson, W. Dorland, A. R. Field, W. Guttenfelder, G. W. Hammett, R. J. Hastie, E. Highcock, N. F. Loureiro, A. G. Peeters, M. Reshko, S. Saarelma, A. A. Schekochihin, M. Valovic, and H. R. Wilson, *Plasma Phys. Controlled Fusion* **51**, 124020 (2009).
- [63] Y.-c. Ghim, A. R. Field, A. A. Schekochihin, E. G. Highcock, C. Michael, and the MAST Team, [arXiv:1211.2883](https://arxiv.org/abs/1211.2883).
- [64] E. A. Frieman and L. Chen, *Phys. Fluids* **25**, 502 (1982).