Multiscale Studies of ETG and Drift Wave Turbulence and Transport Bifurcation Dynamics

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Abstract. Transport barriers are of fundamental interest in fusion research because of the improved confinement they provide, and because they can uncover intrinsic electron transport physics, such as electron temperature gradient (ETG) turbulence. We present recent developments on the generation and stability of extended structures in ETG turbulence. These structures allow the electron thermal transport to exceed the ETG mixing length level, which is too small to be physically relevant. We also present work relevant to the case in which the ρ_i-scale turbulence has not been completely suppressed, via a self-consistent theory of interactions between ETG and drift-ion temperature gradient (DITG) turbulence.

1. Introduction

Understanding the underlying physics mechanism of turbulent transport in magnetic confinement devices is crucial to developing confidence in transport models used to predict the behavior of current and next-step devices. Two areas of intense investigation in recent years has been the possibility of electron temperature gradient (ETG) driven electron thermal transport due to the formation of large-scale structures such as streamers, and the physics of transport barrier formation and dynamics. In this paper we overview some recent progress in the area of ETG dynamics; issues related to transport barriers will be discussed elsewhere.

2. Structure Formation and Collapse in ETG Turbulence

Central to the question of the role of ETG turbulence is understanding the physics of streamers [1-4] and other large-scale coherent structures, particularly what sets their spatial dimensions. The question of streamer formation also represents a specific example of pattern selection and formation in nonlinear systems. In order to investigate this question, the stability of small-scale ETG turbulence to different modulations has been examined. The ETG turbulence is described by a simple fluid model [5,6]

$$\partial_t\left(1-\nabla_\perp^2\right)\phi+\partial_y(\phi+p)-\left\{\phi,\nabla_\perp^2\phi\right\}\approx 0 \qquad (1)$$

$$\partial_t p - r\partial_y\phi+\left\{\phi,p\right\}\approx 0 \qquad (2)$$

where standard normalizations have been used [7], and the Poisson brackets $\{f,g\}=\left(\vec{\nabla}f\times\vec{\nabla}g\right)\cdot\hat{b}$ denote $\vec{E}\times\vec{B}$ convection. A scale separation analysis [8] is employed, in which we restrict ourselves to consideration of modulations of a monochromatic waves i.e. waves of the form $\tilde{\phi} = \Phi\left(\vec{X},T\right)e^{i\left(\vec{k}\cdot\tilde{x}-\omega\tilde{t}\right)}$ + c.c. (where tildes denote fast scales and X, T represent large-scale variations), and derivation of evolution equations for a slowly-varying potential and pressure fields

$$\left(\varepsilon\partial_\tau - \vec{v}_g\cdot\vec{\nabla}\right)\overline{\Phi} + \partial_Y\left(\overline{\Phi}+\overline{P}\right) = 2\left(\hat{z}\times\vec{k}\cdot\vec{\nabla}\right)\left(\vec{k}\cdot\vec{\nabla}\right)\left|\Phi\right|^2 \tag{3}$$

$$\left(\varepsilon\partial_\tau - \vec{v}_g\cdot\vec{\nabla}\right)\overline{P} - r\partial_Y\overline{\Phi} = \frac{rk_y}{\omega}\left(\hat{z}\times\vec{k}\cdot\vec{\nabla}\right)\left(\vec{v}_g - \frac{\omega}{k_y}\hat{y}\right)\cdot\vec{\nabla}\left|\Phi\right|^2 \tag{4}$$

Taking the streamer limit $\partial_X \to 0$, and neglecting the slow time-dependence terms (similar to sub-sonic Langmuir turbulence), one can combine the above equations with the Eqns. 1 and 2 to obtain an evolution equation for the complex wave amplitude

$$i\partial_\tau\Phi + \beta\partial_{YY}\Phi + \alpha\partial_{YY}\left(\left|\Phi\right|^2\right)\Phi = 0 \tag{5}$$

where $\beta = \left(1/2\right)\partial^2\omega/\partial k_y^2$ and α is defined as

$$\alpha = \frac{k_x^2}{\left(2\omega\left(1+k_\perp^2\right)-k_y\right)}\left[\frac{\left(\omega\left(1+2k_\perp^2\right)-k_y v_{gy}\right)\left(r\left(\omega-k_y v_{gy}\right)\right)-2\omega^2 k_y v_{gy}}{\left(r-v_{gy}\left(1-v_{gy}\right)\right)\omega^2} - 2k_y^2\right] \tag{6}$$

Eqn. 5 is a derivative nonlinear Schrödinger equation (DNLS) [9], with cubic second derivative nonlinearity. The nonlinearity is can be attractive (self-focusing) or repulsive depending upon the sign of α/β. Solutions of Eqn. 5 for different parameter regimes have been examined, and are discussed in more detail in Ref. [7].

It is particularly useful to observe that Eqn. 5 can be rescaled and shown to have a Hamiltonian structure of the form

$$H = \int\left[\left|\partial_Y\Phi\right|^2 - \frac{1}{2}\left(\partial_Y\left(\left|\Phi\right|^2\right)\right)^2\right]dY \tag{7}$$

(for $\alpha/\beta > 0$). Defining the variance $V = \int dY\left|\Phi\right|^2 Y^2$, the virial theorem [9] can be used to show

$$\frac{d^2V}{dt^2} = 8H - 2\int dY\left(\partial_Y\left(\left|\Phi\right|^2\right)\right)^2 \tag{8}$$

Noting that the total Hamiltonian H is conserved, Eqn. 8 predicts that the width of the streamer vanishes infinite time (for negative Hamiltonians). Combined with the fact that the total intensity $I = \int dY\left|\Phi\right|^2$ will also be conserved, one can conclude that as the width of the streamer approaches zero, the amplitude must diverge! *Thus, there exists an inherent tendency for streamers to collapse to extended singular shear flow layers.*

The analysis above assumes purely one-dimensional poloidal variations. More directly relevant is the general two-dimensional modulation instability of a small-scale ETG mode with predominately poloidal wavenumber (e.g. the modulational stability of linear ballooning mode with $k_y >> k_x$). More importantly, if one is interested in determining the radial extent of a streamer, there should be no *a priori* assumptions about scale separation in the radial direction. An analysis analogous to the one-dimensional case can be undertaken for the case of a two dimensional modulation, assuming a wave train of the form $\tilde{\phi} = \Phi(x, Y, \tilde{t}, T) e^{i(k_y \tilde{y} - \omega \tilde{t})}$ + c.c., and an isotropic modulation ($\partial_X \sim \partial_Y$), which leads to a two-dimensional amplitude equation

$$i\partial_\tau \Phi + \beta_Y \partial_{YY}\Phi + \beta_X \partial_{XX}\Phi + \alpha k_y^2 \partial_{XX}\left(|\Phi|^2\right)\Phi = 0 \tag{9}$$

$$\alpha = \frac{\left[\left(\omega\left(1+2k_y^2\right)-k_y\right)\left(r\left(v_g - v_\phi\right)/\omega^2 + 2v_g\right) - r\left(\left(v_g - 1\right)\left(v_g - v_\phi\right)/\omega^2 - 2\right)\right]}{\left(v_g\left(v_g - 1\right)+r\right)\left(k_y - 2\omega\left(1+k_y^2\right)\right)} k_y \tag{10}$$

with $\beta_y = \beta$ from the one dimensional case, and $\beta_x = (1/2)\partial^2\omega/\partial k_x^2 = -2\omega^2 / \left(2\left(1+k_y^2\right)\omega - k_y\right)$. Eqn. 9 can be rescaled via $X \to X/\sqrt{\beta_X}$, $Y \to Y/\sqrt{\beta_Y}$, $\Phi \to \left(k_y\sqrt{\alpha/\beta_X}\right)\Phi$, and $\Phi \to \Phi^*$ if $\beta_y <$ 0, and then shown to have a total conserved Hamiltonian of the form

$$H = \int\left[|\partial_X\Phi|^2 + |\partial_Y\Phi|^2 + \frac{\sigma}{2}\left(\partial_X\left(|\Phi|^2\right)\right)^2\right]dXdY \tag{11}$$

where σ = sign(α/β). σ = -1 corresponds to an attractive nonlinearity, and σ = +1 to a repulsive one. To study the elongation and anisotropy of the modulation dynamics, we define the variances $V_X = \int dXdY|\Phi|^2 X^2$ and $V_Y = \int dXdY|\Phi|^2 Y^2$. Equations for the total variance and anisotropy can the be derived, and shown to be

$$\frac{d^2V}{dt^2} = \frac{d^2(V_x + V_Y)}{dt^2} = 8H + 4\sigma\int dXdY\left(\partial_X\left(|\Phi|^2\right)\right)^2 \tag{12}$$

$$\frac{d^2(V_x - V_Y)}{dt^2} = 8H - 16\int dXdY|\partial_Y\Phi|^2 \tag{13}$$

If $H < 0$ (possible for σ = -1), then the total variance V goes to zero in a finite time, and so collapse remains a possibility in the two-dimensional case. However, in this case, the anisotropy measure $V_x - V_y$ is also negative. Since the analysis assumed an initially isotropic modulation, this results suggests that in the final state, $V_x << V_y$! *Therefore, it appears that in the two-dimensional case, collapse to zonal flows over streamers is preferred in ETG turbulence!* Recent numerical work [10] provides evidence for this conclusion. While zonal flows are known to grow more slowly in ETG than ion-temperature gradient (ITG) turbulence [11,12], they are observed to continuously grow in amplitude, while streamer effects are not observed to significantly enhance transport in the time-asymptotic state. Note that this result implies that one should not take the linear mode structure (assumed to be radially extended) as a good indicator of the final (an)isotropy of the system. Ongoing work is investigating saturation scales and levels of the collapse, and suggests that the final radial scale of the zonal flows may be $\Delta X \sim \rho_e^{3/4} L_n^{1/4}$ [13].

3. ETG Interaction with Drift-ITG Turbulence Dynamics

One of the key recent developments in studies of electron thermal transport has been the invocation of ETG turbulence as the source of anomalous electron thermal transport ***outside of ion transport barriers*** [14,15], where drift-ion temperature gradient (DITG) turbulence [16,17] is believed to be unstable and to be the primary driver of all anomalous transport, including electron thermal transport. Here, we term this larger-scale (ρ_i vs. ρ_e) turbulence as drift-ion temperature gradient rather than simply ion temperature gradient to make explicit that we are including both the long-wavelength ($k\rho_i < 1$) curvature-driven component, as well as the somewhat shorter wavelength component ($k\rho_i \sim 1$) of the spectrum which is heavily influenced by kinetic effects such as trapped electrons. Taking ETG turbulence as the dominant driver of electron thermal transport in the bulk of the plasma where DITG turbulence is present represents a fundamental conceptual shift, as it implicitly suggests that ETG turbulence not only drives experimentally relevant levels of electron thermal transport, but dominates over the DITG turbulence which was previously presumed to drive sufficient electron thermal transport to explain the observed levels. Given that the proposed mechanisms by which ETG turbulence generates transport at levels greater than the expected gyroBohm level involve formation of structures much greater than the characteristic ρ_e scale (e.g. streamer formation or inverse cascade to the collisionless skin depth $\delta_e = c/\omega_{pe}$), it is important to investigate the dynamics of ETG turbulence, and in particular, these large-scale structures, in the presence of DITG turbulence.

3.1 Random Shear Suppression of ETG Turbulence by DITG Turbulence

In order to elucidate the effects of DITG turbulence on ETG turbulence, we have studied the suppression of ETG turbulence due to random ***two-dimensional*** shearing by DITG turbulence. We exploit the difference in space and timescales between the ETG and DITG (ρ_e vs. ρ_i and v_{Te}/L_n vs. v_{Ti}/L_n, respectively) to describe the evolution of the ETG turbulence in terms of an adiabatic invariant (the potential enstrophy) $N(\vec{x},\vec{k}) = (\tau + k_\perp^2\rho_e^2)|\phi_k^{ETG}|^2 + |T_k^{ETG}|^2$ [18], which evolves via a wave-kinetic equation [19]

$$\frac{\partial N}{\partial t} + \frac{\partial}{\partial \vec{k}}\left(\omega_k + \vec{k}\cdot\vec{V}_{DITG}\right)\cdot\frac{\partial N}{\partial \vec{x}} - \frac{\partial}{\partial \vec{x}}\left(\vec{k}\cdot\vec{V}_{DITG}\right)\cdot\frac{\partial N}{\partial \vec{k}} = 2\gamma_k N - \Delta\omega_k N^2 \qquad (14)$$

Here $\vec{V}_{DITG} = v_{Te}\hat{b}\times\rho_e\vec{\nabla}\phi^{DITG}$ is the velocity field due to the (ρ_i-scale) DITG turbulence. We emphasize that this flow field corresponds to the *entire* DITG velocity field, not just the dominant zonal flow component, and in particular includes the slightly smaller-scale fluctuations due to trapped electrons and other effects. We then use a quasi-linear approach to rewrite the effects of the DITG flow field on the ETG turbulence in terms of a k-space diffusion tensor

$$\frac{\partial\langle N\rangle}{\partial t}=\frac{\partial}{\partial x_\alpha}D_{\alpha\beta}\frac{\partial\langle N\rangle}{\partial x_\beta}+2\gamma_k\langle N\rangle-\Delta\omega_k\langle N\rangle^2 \tag{15}$$

$$D_{\alpha\beta}=\rho_e^2 v_{Te}^2\sum_q q_\alpha q_\beta\left|\left(\vec{k}\times\vec{q}\right)\cdot\hat{b}\right|^2 R\left(\Omega_q\right)\left|\phi_q^{DITG}\right|^2 \tag{16}$$

$$R\left(\Omega_q\right)=\frac{1}{2\gamma_k-i\left(\Omega_q-\vec{q}\cdot\vec{v}_g\right)}\approx\frac{1}{2\gamma_k} \tag{17}$$

The tensor structure of $D_{\alpha\beta}$ reflects the fact that the DITG flow field is a two-dimensional shear field, such that radial shear leads to diffusion in k_r, and poloidal shear leads to diffusion in k_θ; this tensor form has also been discussed by Hahm and Burrell [20]. The brackets denote the quasi-linear averaging, meaning that Eqn. 15 describes the evolution of the ETG intensity *on DITG space and timescales*.

The importance of DITG shearing can be estimated by comparing the k-space diffusion rate $\gamma_D\approx D_{\alpha\beta}/k_\alpha k_\beta$ against the self-damping rate $\gamma_{self}=\Delta\omega\langle N\rangle\approx 2\gamma_k=\gamma_{lin}$ (the equivalence between γ_{self} and γ_{lin} reflects the saturation level in the absence of DITG turbulence). If diffusion is much more rapid than linear growth $\left(\gamma_D>>\gamma_{lin}\right)$, then the DITG turbulence will rapidly take energy from linearly unstable wavenumbers to larger, stable wavenumbers, and thereby suppress the ETG turbulence. Conversely, if diffusion is much smaller than the linear growth rate $\left(\gamma_D<<\gamma_{lin}\right)$, then the ETG "self-saturation" mechanism dominates and the ETG turbulence is insensitive to the presence of the DITG turbulence. Estimation of the diffusion rate requires an estimation of the DITG spectrum, for which we use a mixing-length estimate which represents an upper bound for DITG intensity (since it does not reflect the suppression of DITG mode by ρ_i-scale zonal flows). Taking $\left|\phi_q^{DITG}\right|^2=1/q^2L_{DITG}^2$, $L_{DITG}=\sqrt{L_{Ti}L_B}$ (representing curvature-driven modes), we have

$$\gamma_D=\frac{D_{\alpha\beta}}{k_\alpha k_\beta}\approx\frac{\rho_e^2 v_{Te}^2}{\gamma_{lin}}\sum_q q^4\left|\phi_q^{DITG}\right|^2\approx\frac{\rho_e^2 v_{Te}^2}{\gamma_{lin}}\frac{\overline{q^2}}{L_{DITG}^2} \tag{18}$$

$$\frac{\gamma_D}{\gamma_{lin}}=\left(\frac{v_{Te}}{\gamma_{lin}L_{DITG}}\right)^2\overline{q^2}\rho_e^2=\frac{m}{M}\frac{\overline{q^2}\rho_i^2}{k_\theta^2\rho_e^2}\frac{\tau^2\eta_i}{\eta_e-\eta_e^c} \tag{19}$$

, where we have used $\gamma_{lin}\approx\left(v_{Te}/L_n\right)\left(k_\theta\rho_e\right)\sqrt{\varepsilon\left(\eta_e-\eta_e^c\right)/\tau}$ ($\varepsilon=L_n/L_B$, $\tau=T_{e0}/T_{io}$, $\eta_e=L_n/L_{Te}$, and $\eta_e{}^c$ is the critical value of η_e needed for ETG instability), and $\overline{q^2}$ is the mean square wavenumber of the DITG turbulence. In general, the diffusive rate will be much smaller than the linear growth rate, as evidenced for by the electron to ion mass ratio m/M which physically represents the scale separation of the ETG and DITG modes. However, for the cases of streamers or skin depth fluctuations, one has a more interesting result:

1. Streamers with $k_\theta \rho_e \approx 0.1$ have a ratio

$$\frac{\gamma_D}{\gamma_{lin}} \approx 100 \frac{m}{M}\left(\overline{q^2}\rho_i^2\right)\frac{\tau^2 \eta_i}{\eta_e - \eta_e^c} \tag{20}$$

2. Collision skin depth scale fluctuations with $k_\theta \rho_e \approx \rho_e / \delta_e = \sqrt{\beta_e}$, $\beta_e = 8\pi n_0 T_e / B_0^2$ (β_e is generally between 0.01 and 0.1 for current high-power tokamaks) have a ratio

$$\frac{\gamma_D}{\gamma_{lin}} \approx \frac{m}{M}\frac{\overline{q^2}\rho_i^2}{\beta_e}\frac{\tau^2 \eta_i}{\eta_e - \eta_e^c} \tag{21}$$

In either case, one finds that these large-scale components of the ETG turbulence may be susceptible to shearing by the short-wavelength component of the DITG turbulence, particularly if the ETG turbulence is close to marginality. In essence, the large-scale ETG structures become sufficiently close in size to the short-wavelength DITG turbulence to experience serious interactions. *These scalings therefore suggest that streamers or δ_e-scale fluctuations, which are the proposed mechanisms for experimentally relevant ETG-driven transport, may be significantly impacted by DITG shearing.* This finding not only calls into question the validity of assuming ETG driven electron thermal transport in the bulk of the plasma, but also in the transport barriers where DITG is believed to be suppressed, since the short-wavelength component of the DITG turbulence is more likely to survive in the shear region which is believed to suppress the long-wavelength component of the DITG, and it is precisely this component which is relevant for shearing the ETG turbulence. Taken with the possibility that TEM modes (or at least their short-wavelength component) may be able to survive in regions of finite shear and still produce significant levels of electron thermal transport, as well as the results of Section 2 which raise serious questions about the viability of streamers in ETG turbulence, ***one is naturally lead to the conclusion that the relevance of ETG turbulence for explaining tokamak transport levels highly questionable, and that there are serious conceptual and physical issues which much be addressed before a true understanding of the physics of electron thermal transport can be claimed!***

3.2 Effects of DITG Modulations of η_e

In addition to direct shearing of ETG turbulence, DITG modes can also affect ETG turbulence via modulations and distortions of other relevant fields, such as the electron temperature field T_e. The impact of DITG modulations of T_e will be qualitatively different, since the ρ_i-scale modulations of T_e will appear as modulations of the equilibrium T_e gradient L_{Te}, or equivalently $\eta_e = L_n/L_{Te}$ on ρ_e scales. Although we do not consider modulations of L_n here, they could in principle be treated in a manner analogous to the analysis below. To lowest order, modulations of η_e represent modulations of $\gamma_{ETG} \propto \sqrt{\eta_e - \eta_e^c}$, where η_e^c represents a critical value of he needed for ETG instability. Writing $\eta_e = \eta_e^0 + \delta\eta_e$, we again exploit the separation of space and timescales to use a wave-kinetic description of the ETG turbulence, which is linearized to provide the response of the ETG turbulence to the modulation $\delta\eta_e$

$$R^{-1}\left(\Omega_q\right)\delta N_q = 2\left.\frac{\partial\gamma_k}{\partial\eta_e}\right|_{\eta_e=\eta_e^c}\delta\eta_q\langle N\rangle \tag{22}$$

where $R(\Omega_q)$ is defined in Eqn. 17, and we have expressed $\delta\eta_e = \sum_q \delta\eta_q \exp\left(i\left(\vec{q}\cdot\vec{x} - \Omega_q t\right)\right)$. Ignoring the shearing effects of the DITG turbulence, we again use a quasilinear closure to describe the dynamics of $\langle N\rangle$ as

$$\frac{\partial\langle N\rangle}{\partial t} = 2\left(\gamma_k + \gamma_{NL}\right)\langle N\rangle + O\left(\langle N\rangle^2\right) \tag{23}$$

$$\gamma_{NL} = 2\left(\left.\frac{\partial\gamma_k}{\partial\eta_e}\right|_{\eta_e=\eta_e^c}\right)^2 \sum_q R\left(\Omega_q\right)\left|\delta\eta_q\right|^2 \tag{24}$$

$$\Rightarrow \frac{\gamma_{NL}}{\gamma_k} = \frac{\sum_q\left|\delta\eta_q\right|^2}{4\left(\eta_e^0 - \eta_e^c\right)^2} \sim \left(\frac{\delta\eta^{RMS}}{\eta_e^0 - \eta_e^c}\right)^2 \tag{25}$$

where we have used $\gamma_k \propto \sqrt{\eta_e - \eta_e^c}$ and $R\left(\Omega_q\right) \approx 1/2\gamma_k$. Therefore, when the RMS amplitude of the η_e modulations is comparable to the deviation of η_e from its critical value (a condition which is readily achievable for realistic DITG intensities), this modulation effect will be important, and the effective growth rate of the ETG turbulence will be enhanced from its "bare" value. It is important to note that the analysis presented here is a mean-field treatment which implicitly assumes $\eta_e^0 + \delta\eta_e(t) > \eta_e^c$ i.e. the modulation of η_e never brings η_e below the critical value, which would stabilize the ETG instability. However, it is clear that such a situation could easily occur (particularly if the system was close to ETG marginality), in which case describing the dynamics becomes much more complex. It is therefore easy to see that a much richer range of dynamics is possible than described here, including possibilities such as submarginal ETG turbulence, or complex spatio-temporal behavior of the ETG turbulence *on DITG space and time scales* are two obvious possibilities. Also note that this interaction does not depend explicitly upon the electron-ion mass ratio, in contrast to the shearing effects discussed in Section 3.1 (although it does implicitly rely upon separation of scales). *Thus, numerical investigations of ETG-DITG interactions which use artificial mass ratios will not correctly calculate the relative importance of flow shearing vs. η_e modulations!* Finally, we observe that while the η_e modulation effect described here represents a new avenue for investigating interactions between ETG and DITG turbulence, it is not yet clear what impact it would have upon the large scale structures in ETG turbulence described in Section 2, and how it would affect their relevance in explaining the net electron thermal transport level.

6. Acknowledgements

We thank T. S. Hahm, Z. Lin, and W. Horton for collaboration and useful discussions. This research was supported by D.O.E. Grant No. DE-FG03-88ER53275 and FG02-04ER54738.

C. H. performed this research under an appointment to the Fusion Energy Sciences Fellowship Program, administered by Oak Ridge Institute for Science and Education under a contract between the U.S. Department of Energy and the Oak Ridge Associated Universities.

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