Collisional Damping of ETG-Mode-Driven Zonal Flows

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We study collisional damping of electron zonal flows in toroidal electron temperature gradient (ETG) turbulence due to the friction between trapped and untrapped electrons. With the assumption of adiabatic ions, the collisional damping is shown to occur on fast time scales $\sim 0.24\epsilon^{1/2}\tau_e$. The comparison with the growth rate of electron zonal flows indicates that the shearing by electron zonal flows is unlikely to be a robust mechanism for regulating ETG turbulence. This finding vitiates the claims of several simulation studies that have ignored the effects of collisional damping of electron zonal flows and offers a possible partial explanation of the high levels of electron thermal transport observed in the National Spherical Torus Experiment.

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One of the greatest challenges in magnetic confinement fusion is to understand both the origin and mechanisms for suppression of anomalous electron heat transport [1]. In the past few decades, several candidate mechanisms for anomalous electron heat transport, such as the trapped electron modes (TEM) [2], the current-diffusive ballooning mode [3], and the electron temperature gradient (ETG) model [4–6], have been proposed, spanning the entire range of relevant scales from ρ_e to a (minor radius). Among these, the ETG model appears to be promising, with one of its merits being the natural separation of χ_e from χ_i and D. However, since ETG excites fluctuation energy at small scales ($\sim \rho_e$) with high frequency ($\sim v_{Te}/Rq$ in torus and $v_{Te}/L_\perp$ in slab), a naive mixing length estimate based on ρ_e yields $\chi_e^{\rm ETG} \sim \chi_i^{\rm ITG}/60$ (ITG: ion temperature gradient), which is too small to be consistent with experimental observations ($\chi_e \sim \chi_i$). Thus, a successful ETG model *must* provide a physical mechanism which significantly enhances χ_e over mixing length estimate levels, possibly giving an (radial) effective transport scale length significantly larger than ρ_e. An important issue in the determination of χ_e is to determine whether formation of extended structure is possible. As zonal flow shearing naturally inhibits formation of such structures, we thus are intensely interested in zonal flow damping mechanisms. Indeed, a proper understanding and representation of such a flow damping mechanism is essential for meaningful computer simulations of ETG turbulence. While TEM is another strong candidate for electron thermal transport, in this Letter, we examine the ETG mode as one of the plausible theoretical candidates.

One mechanism for regulating structure scale in ETG turbulence is the shearing by self-generated electron zonal flows [7]. Electron zonal flows are axisymmetric poloidal flows ($k_\theta = 0$ and $k_\phi = 0$) [8], generated by modulational instability, for instance, and are the analogs of ion zonal flows in ion temperature gradient turbulence [9]. While the shearing by ion zonal flows is now well-known to regulate and thus stabilize ITG turbulence, this shearing is unlikely to affect ETG turbulence which has very small scales ($\sim \rho_e$). In comparison, electron zonal flows have much smaller radial scales than ion zonal flows, and thus their shearing can be potentially effective in regulating ETG turbulence. As such, the elucidation of the dynamics of electron zonal flows, i.e., both their growth and damping, is critical to assess their role in the regulation of ETG turbulence.

Electrostatic ETG has a strong resemblance to ITG with the role of electrons and ions reversed (when the Debye screening effect is negligible, i.e., $\lambda_{\rm De} \ll \rho_e$). *Note that the difference between the two is that ions can be treated as adiabatic for electron zonal flows (as $\rho_i k \gg 1$), while electrons are nonadiabatic for ion zonal flows.* It is well-known that in a toroidal ITG, ion zonal flow undergoes both collisionless and collisional damping, due to the charge screening by both trapped and untrapped ions and pitch angle scattering between the two [10,11]. In view of the similarity between ITG and ETG, a similar damping of electron zonal flows is expected to occur in ETG due to the friction between trapped and untrapped electrons. In fact, since $\nu_{ee}, \nu_{ei} \gg \nu_{ii}$, the collisional damping of electron zonal flows is likely to be much stronger than ion zonal flows damping.

In the present Letter, we limit ourselves to the case where the Debye screening effect is negligible ($\lambda_{\rm De} \ll \rho_e$) and assume adiabatic ions for electron zonal flows. We consider an axisymmetric magnetic field $\mathbf{B} = \nabla\xi \times \nabla\psi + I\nabla\xi$ and axisymmetric potential perturbation $\phi(\mathbf{x}, t) = \phi_k(\mathbf{x}) \exp[iS(\mathbf{x}_\perp)]$ with wave number $\mathbf{k}_\perp = S'(\psi)\nabla\psi$. Here, $I = RB_\phi$, $S' = \partial_\psi S$, ψ is the poloidal flux function ($\partial_r\psi = RB_p$), ξ is the toroidal angle variable, and $\phi_k(\mathbf{x})$ is a slowly varying function of space. For electron zonal flows, the scale of potential lies between ρ_e and ρ_i, i.e., $k_\perp\rho_e < 1 < k_\perp\rho_i$ [$\rho_e = (T_e/m_e)^{1/2}/|\Omega_e|$ and $\rho_i = (T_i/m_i)^{1/2}/\Omega_i$ are electron and ion gyroradii].

We assume ϕ_k to be constant on the flux surface, i.e., $\phi_k = \langle\phi_k\rangle$, since it survives Landau damping. Here, angular brackets denote flux surface average

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$\langle A \rangle = \oint (dl/B_p) A / \oint (dl/B_p)$. Then, the quasineutrality condition for the k zonal mode takes the following form:

$$n_0 \frac{e}{T_{e0}} k_\perp^2 \rho_e^2 \phi_k + \langle n_{ek} \rangle = -\frac{e\phi_k}{T_{io}} n_0. \quad (1)$$

Here, n_0 is the background density associated with Maxwellian distribution F_0, while $n_{ek} = \int d^3 v f_{ek}$ is the perturbed electron guiding center density with distribution function f_{ek}; the first term in Eq. (1) is the electron polarization density.

The perturbed electron guiding center distribution function f_{ek} satisfies the drift kinetic equation. Following the analysis of Ref. [8], we obtain the solution for f_{ek} by a perturbation expansion in terms of the small parameter ω/ω_b (ω_b bounce frequency) as $f_{ek} = i(e/T_e) F_0 Q \phi_k + h_k$, where $\hat{b} \cdot \nabla h_k = 0$, $\hat{b} = \mathbf{B}/B$, and $Q = I \partial_\psi S v_\parallel / \Omega_e$. Note that ω^{-1} is a time scale for zonal flows and that the geodesic acoustic mode is not considered in this Letter. h_k is then determined from the solubility condition obtained by averaging

$$\partial_t h_k - \overline{(C_e h_k)} = -i \frac{e}{T_{e0}} F_0 \overline{Q} \partial_t \phi_k + \overline{S}_{ek}. \quad (2)$$

Here, C_e is the collision term due to electron-electron and electron-ion collisions, S_{ek} is the source term, and the average $\overline{A} = \oint (dl/v_\parallel) A / \oint (dl/v_\parallel)$ is taken along a closed orbit for trapped particles and along one poloidal circumference for untrapped particles.

To obtain the potential to second order in banana width, we first compute the surface averaged radial current $\langle \mathbf{J} \cdot \nabla\psi \rangle = -e \langle \int d^3 v \mathbf{v}_D \cdot \nabla\psi f_{ek} \rangle = eI \langle \int d^3 v (v_\parallel / \Omega_e) v_\parallel \hat{b} \cdot \nabla f_{ek} \rangle$, and then the electron density on the flux surface $\langle n_{ek} \rangle = -\langle \rho_k \rangle / e$, via $\partial_t \langle \rho_k \rangle = -\partial_\psi \langle \mathbf{J} \cdot \nabla\psi \rangle$. The quasineutrality condition then gives the potential in terms of the electric susceptibility ζ_k as

$$\tilde{\zeta}_k(p) \tilde{\phi}(p) = \frac{-1}{p(n_0 e / T_{e0})} \left\langle \int d^3 v \tilde{S}_{ek} \right\rangle, \quad (3)$$

$$\tilde{\zeta}_k(p) = \tau + (k_\perp \rho_e)^2 + \left\langle \frac{(m_e c I S')^2}{n_0 e^2} \int d^3 v \frac{v_\parallel}{B} \left(\frac{v_\parallel}{B} F_0 - g \right) \right\rangle. \quad (4)$$

Here, $\tau = T_e/T_i$, $\tilde{A}(p) = \int dt A(t) e^{-pt}$ is the Laplace transform of A, and $g = -i(T_{e0} \tilde{h}_k / I S' m_e c \tilde{\phi}_k)$ is the Laplace transform of h_k up to a normalization factor, which satisfies

$$g - \frac{1}{p} \overline{C_e g} = \overline{\left(\frac{v_\parallel}{B} \right)} F_0. \quad (5)$$

Since we are interested in the damping of the potential, we consider an initial value problem by assuming $S_{ek} = f_{ek}(0) \delta(t)$. This initial electron density perturbation is accompanied by rapid potential adjustment due to classical polarization, in order to satisfy the quasineutrality condition (1) on the order of electron gyroperiod ($\sim 1/|\Omega_e|$). If an adiabatic ion response is established within this time, the evolution of the potential for $t > 1/|\Omega_e|$ can be described with the initial electron density perturbation

$$\langle n_{ek}(0) \rangle = -n_0 \frac{e}{T_{e0}} [k_\perp^2 \rho_e^2 + \tau] \phi_k(0). \quad (6)$$

Then, the time evolution of the potential takes the following simple form:

$$\phi_k(t) = K(t) \phi_k(0), \quad (7)$$

where $K(t)$ is the normalized response kernel with Laplace transform $\tilde{K}(p) = (k_\perp^2 \rho_e^2 + \tau) / p \tilde{\zeta}(p)$.

Let us first consider the collisionless case. In this case, the solution to Eq. (5) is simply given by $g = \overline{(v_\parallel / B)}$. By using this, and assuming a large-aspect ratio tokamak ($\epsilon = r/R \ll 1$) with circular cross section, we can compute $\tilde{\zeta}(p)$ and $K(t)$ in Eqs. (4) and (7), respectively, by doing standard integrals [12] to obtain

$$\phi(t) = \frac{\tau + k_\perp^2 \rho_e^2}{\tau + k_\perp^2 \rho_e^2 [1 + 1.6 q^2 / \sqrt{\epsilon}]} \phi(0). \quad (8)$$

The second term in the square brackets comes from the neoclassical polarization due to finite banana orbits. Note that the ITG case is recovered by putting $\tau = 0$ and $\rho_e \to \rho_i$ in Eq. (8). If $(k_\perp \rho_e)^2 < \tau$, the residual potential in the ETG is larger than in the ITG case. In the opposite short wavelength case where $(k_\perp \rho_e)^2 > \tau$, the residual potentials for ETG and ITG become comparable.

It is interesting to note that $g \propto \overline{v}_\parallel$ reduces the charge screening [see Eq. (4)]. Since g vanishes for trapped electrons, there is a discontinuity in g at the trapped-untrapped boundary. With collisions, g becomes a smooth function of pitch angle λ through pitch angle scattering. In fact, in the long time limit, the asymmetry in g (with $v_\parallel$) is wiped out via pitch angle scattering, and g asymptotically approaches to zero value. That is, passing electrons with $v_\parallel > 0$ get converted to those with $v_\parallel < 0$ by scattering. Therefore, the overall screening due to passing electrons is more significant in the collisional case, and this screening is the cause of the larger damping of the potential. In more physical terms, this is because of the damping of current carried by passing electrons as the asymmetry in the distribution of passing electrons in $v_\parallel$ decreases via pitch angle scattering.

The collisional damping of the potential may be seen from the following two steps. The initial stage for $t < \epsilon \tau_e$ consists of rapid adjustment of g through pitch angle scattering (primarily near the trapped-untrapped boundary) leading to the algebraic decay of potential [13]. Here, τ_e is the microscopic collision time of electrons. This initial algebraic damping rate depends on the initial distribution of g, becoming very rapid for a distribution of g localized near the trapped-untrapped boundary. At later time $t > \epsilon^{1/2}\tau_e$, g damps exponentially, and asymptotically approaches zero. It is easy to see that the damping of g (or h_k) is directly related to the damping of poloidal flows u_p as follows [13]:

$$u_p = \frac{\mathbf{u}\cdot\nabla\theta}{|\nabla\theta|} = u_\parallel \frac{B_p}{B} + \frac{cB_\phi}{B^2}\left(\partial_r\phi + \frac{B}{en_e}\partial_r p\right), \quad (9)$$

where $u_\parallel = (1/n_0)\int d^3v v_\parallel f_{ek}$. By assuming that pressure gradient is negligible compared to the electric field in Eq. (9) (as we did in obtaining the solution f_{ek}), and then by substituting f_{ek}, one can show that $u_p \simeq (B_p/Bn_0)\int v_\parallel h_k$ for a large-aspect ratio tokamak ($B \simeq B_\phi \gg B_p$). Therefore, poloidal flows damp out due to collisions, while the potential approaches a finite value. The value of this residual potential can be computed from Eq. (4) as

$$\phi(t) \simeq \frac{\tau + k_\perp^2\rho_e^2}{\tau + k_\perp^2\rho_{ep}^2}\phi(0), \quad (10)$$

where $\rho_{ep} = \rho_e(B/B_p) > \rho_e$ is the electron gyroradius for poloidal magnetic field B_p. Thus, for $(k_\perp\rho_{ep})^2 < \tau$, $\phi(t)$ does not damp significantly, while for $(k_\perp\rho_{ep})^2 > \tau$, $\phi(t) \sim (\tau/k_\perp^2\rho_e^2)(B_p/B)^2\phi(0)$. Of course, retaining residual ion finite Larmor radius effects (for $k_\perp\rho_i > 1$) changes the condition on $k_\perp^2\rho_i^2$ from $k_\perp^2\rho_i^2 \gg 1$ to $k_\perp^2\rho_e^2 > 1/\sqrt{2\pi}k_\perp\rho_i$ [14]. For $\tau \sim 1$ and deuterium, this is equivalent to $k_\perp\rho_e > 0.2$ and $k_\perp\rho_{ep} > 0.2B/B_\theta$. Note for ITG, while the zonal flow spectrum peaks at $k_\perp\rho_i \sim 0.1$, analogous significantly higher $k_\perp$ content *does* exist [15]. Thus, the condition for the applicability of the long wavelength neoclassical polarization density is marginally satisfied. Also, it is interesting to note that the simulations in [6,8,16,17] assume purely adiabatic ions, resulting in the reduction of the ETG zonal flow growth due to enhanced inertia.

The time scale over which this collisional damping occurs is much shorter than that for the ITG case since $\nu_{ee}, \nu_{ei} \gg \nu_{ii}$. In order to compute this damping rate, we assume $\tau = T_{i0}/T_{e0} = 1$ and use a momentum conserving pitch angle scattering operator for the collision term, including both electron-electron and electron-ion collisions [18],

$$C_e g = \left(\frac{T_e}{E}\right)^{3/2}\left[\frac{2\nu_e}{B}\sqrt{1-\lambda B}\,\partial_\lambda(\lambda\sqrt{1-\lambda B}\,\partial_\lambda g) + \bar{\nu}_e v_\parallel F_0 \frac{\int d^3v \nu_e v_\parallel g}{\int d^3v \nu_e v_\parallel^2 F_0}\right]. \quad (11)$$

Here, $E = m_e v^2/2$, $\nu_e = \bar{\nu}_e[1 + H(\sqrt{E/T_e})]$, $\bar{\nu}_e = (\nu_{ee}/m_e^{1/2}(2T_e)^{3/2})$, $\nu_{ee} = 4\pi n_0 e^4 \ln\Lambda$, Λ is the Coulomb logarithm, $H(z) = E'(z)/2z + (1 - 1/2z^2)E(z)$, $E(z)$ is the error function, and λ is the pitch angle. Note that ν_e includes both electron-ion (the first term in the square brackets) and electron-electron collisions (the second term in the square brackets). By using (11) and following a similar analysis as in Ref. [8], we obtain

$$\phi(t) \simeq \frac{\tau + k_\perp^2\rho_e^2}{\tau + k_\perp^2\rho_{ep}^2}\left[1 + \frac{\tau + k_\perp^2\rho_{ep}^2}{k_\perp^2\rho_{ep}^2}\frac{1.4}{\epsilon}\exp(-t/t_d)\right]\phi(0), \quad (12)$$

with the damping rate $\tau_d \simeq 0.24\epsilon^{1/2}\tau_e$, where $\tau_e = 3\sqrt{\pi}/4\nu_{ee}$ is the collision time. As expected, poloidal flows damp exponentially fast on a time scale $1/\tau_d \propto \nu_{ee}$, which is much shorter than that for the ITG case (by a factor of $\sqrt{m_i/m_e}$). It is important to note though that in normalized units ($1/\Omega_i$ and $1/|\Omega_e|$ for ITG and ETG, respectively), the zonal flow damping rate in ETG becomes comparable to that in the ITG case. Therefore, in order to determine whether electron zonal flows can regulate ETG turbulence, we need to compare their damping rate with their growth rate (for instance, by modulational instability). Because of the adiabatic ion response (which enhances zonal flow inertia), the growth rate of electron zonal flows $\{\propto (\rho_e k)^4/[\tau^{-1} + (\rho_e k)^2] \simeq (\rho_e k)^4\tau$ with $\rho_e k < \tau^{-1/2}\}$ is, however, much smaller than that of ion zonal flows [$\propto (\rho_i k)^2$] in the same normalized units. This result casts a rather strong doubt on the conclusion reached in [19] regarding the effects of electron zonal flows on ITG turbulence.

To be more quantitative, we ignore the presence of streamers and estimate the saturation level of ETG turbulence on the ρ_e scale by balancing the growth rate of ETG zonal flows against their collisional damping. By using $p\rho_e \sim 1$ and $k\rho_e \ll 1$ (p and k are the characteristic wave numbers of turbulence and zonal flows) in the growth rate of ETG zonal flow [7], we obtain

$$\gamma_q \sim \frac{v_{Te}}{L_n}(k\rho_e)^4 \frac{1}{\sqrt{\epsilon(\eta - \eta_c)/\tau}}\frac{I}{I_{ML}}. \quad (13)$$

Here, $I = \int d^3p|e\tilde{\phi}/T_e|^2$ and $I_{ML} = (\rho_e/L_n)^2$ is the mixing length estimate. By balancing Eq. (13) against $\gamma_d \sim \nu_{ee}/\epsilon^{1/2}$, we find that

$$\frac{I}{I_{ML}} \sim \frac{\nu_*}{\tau^{1/2}}\frac{(\eta - \eta_c)^{1/2}}{(k\rho_e)^4}. \quad (14)$$

Here, $\nu_* = \nu_{ee}L_n/v_{Te}$ is the normalized collisionality. Thus, the key parameters in setting the turbulence intensity level are the characteristic scale of zonal flows, deviation from the marginality, and collisionality. The scaling of $I \sim (k\rho_e)^{-4}$ in Eq. (14) suggests a large saturation level, on the order of $100I_{ML}$, assuming typical core tokamak parameters and $(k\rho_e) \sim 1/10$. Of course, retaining the residual nonadiabatic ion response, as discussed above, will weaken the dependence of I/I_{ML} on $k\rho_e$ and constrain $k\rho_e > 0.2$ for $\tau = 1$. Even with this caveat, the predicted saturation level, which is rather large, highlights the inefficiency of zonal flow shearing in saturating ETG turbulence and suggests that another saturation mechanism must be operating.

It is instructive to consider the implication of collisional damping in the context of "pattern selection"; namely, what determines if zonal flows or streamers are formed? This question is particularly important as it is argued that it is the occurrence of large amplitude streamers which allows ETG to drive experimentally relevant transport. Thus, determining the parameter space for streamer formation is key to understanding ETG-driven transport. Simulations in [6,16] indicate that linear streamers are more likely to form for positive shear and lower α, while weak/negative shear and/or large α are expected to give lower transport and saturation levels. The latter regime supports enhanced zonal flow formation, similar to sheared slab simulations. Note that with collisions included, the regime of ETG zonal flow relevance is likely to be narrowed in [6] as the zonal flow damping would enhance the thermal transport [20] and the formation of streamers. Zonal flow formation in slab geometry with a negative shear has also been observed in [8,17]. Weak or negative shear and large α correspond to the case of internal transport barrier regions of advanced tokamak operating regimes, where anomalous electron heat transport is sometimes observed even if ion thermal transport and particle transport are suppressed. These observations are seemingly inconsistent with the aforementioned simulation results. However, *noting that these simulations were performed without collisions, these conflicting results may be reconciled* by the incorporation of collisional damping of ETG zonal flows described in this Letter. Note also that in a 2D slab geometry, there is no equivalent damping of electron zonal flows due to the absence of trapped electrons. Thus, the dynamics of zonal flows and ETG turbulence, and thus pattern selection, is quite sensitive to the geometry of the problem. It should also be noted that a recent gyrofluid simulation of ETG turbulence [21] found the result that streamers were only slightly anisotropic in the turbulent state, so that heat transport was only weakly enhanced over the (rather insignificant) mixing length level (i.e., $\chi_e \sim 2\chi_{e,ML}$). In this simulation, zonal flows were artificially suppressed by the imposition of an unrealistically large damping. As discussed above, an accurate zonal flow damping decrement implies (when balanced against production) modest levels of electron flow $\mathbf{E} \times \mathbf{B}$ shear which, however, are not necessarily negligible. Thus, results of [21] may be viewed as an *upper bound* on the true fluctuation levels. Some caution should, however, be taken in interpreting the former results since they are based on very long wavelength modes with $(k_\perp \rho_e)^2 \simeq 10^{-4}$. Further experimental studies quantifying the dependence of χ_e on magnetic shear and α would also be very helpful.

An interesting direct application of these results is the question of the high level of heat transport observed in neutral beam heated National Spherical Torus Experiment (NSTX) plasmas [22]. While the flow damping calculation presented here is limited formally to the regime of large-aspect ratio, it is certainly reasonable to expect the flow damping to increase for smaller aspect ratios, such as those found in a spherical torus. Thus, ETG zonal flows are likely to be heavily damped in NSTX, thereby removing an important self-regulation mechanism for the turbulence, and naturally implying higher levels of electron heat transport. Further investigation of this suggestion requires a more general form of the neoclassical density [23] to extend the ETG zonal flow damping theory to regimes of low aspect ratio, and of course, quantitative comparisons with experiment. Finally, we remark that due to the neglect of finite alpha and magnetic shear effects, our results cannot be directly applied to electron transport in internal transport barrier plasmas.

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