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THE MODIFICATION OF TURBULENT TRANSPORT BY
ORBIT AVERAGING

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The Modification of Turbulent Transport by Orbit Averaging

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ABSTRACT

The effect on plasma turbulence of orbit averaging by thermal ions is considered, and illustrated for two modes of potential importance for tokamaks. The effect can reduce the ion response below that in earlier treatments, modifying the predicted mode growth rate, which in turn modifies the turbulent transport. For both modes, the effect modifies earlier transport expressions with a 'neoclassical factor,' which makes the scalings of the resultant transport coefficients with plasma current and magnetic field closer to those found experimentally. Additionally, for the trapped electron mode, this mechanism provides a potential explanation of the observed more favorable scaling of χ_i with T_i in supershots than in L-modes.

For some time, it has been appreciated theoretically¹ that the net effect of field fluctuations on particles with appreciable gyroradii ρ_g and drift excursions ρ_b can be substantially different from (usually, smaller than) their effect in the $\rho_{b,g} \rightarrow 0$ limit, due to an averaging over the fluctuation phase

MASTER

which a particle with finite Larmor radius and finite drifts performs in the course of a gyroperiod τ_g or bounce (or transit) period τ_b . This reduction in interaction strength causes a corresponding reduction in the transport levels induced by a given spectrum of fluctuations. This theoretical expectation is supported by experimental observations of transport rates of energetic electrons³ and ions,^{4,5} as well as by numerical simulations.⁶

Here, we consider a complementary effect less widely appreciated, namely, that not only does orbit averaging diminish the extent to which a specified spectrum of fluctuations affects particles (as manifested by the transport coefficients D or χ),^{1,2} but also it diminishes the extent to which particles affect fluctuations (as manifested by the linear susceptibility K), and hence modifies the stability of modes which cause turbulent transport. We apply this basic idea to two illustrative modes which may be important in causing anomalous transport in tokamaks, the trapped electron mode (TEM)⁷ and the ion temperature gradient (ITG) mode,⁸ and obtain estimates for the modification by orbit averaging of tokamak turbulent transport coefficients. The averaging due to finite ρ_s introduces an extra q -dependence into the transport coefficients, similar to the one encountered in going from classical to neoclassical transport, which brings the resultant scalings with current I_p and magnetic field B into closer agreement with those observed experimentally. Moreover, the modified expressions for the TEM predict an important effect observed on TFTR,^{9,10} namely, the decrease in χ_i with increased T_i for supershots, for which $\tau \equiv T_e/T_i \ll 1$), in contrast to the usual increase in χ_i with increased T_i in L-mode plasmas (where $\tau \sim 1$) (see Fig.1, from 10).

The formal theory permitting study of both types of orbit-averaging

effects was developed by Kaufman,¹¹ who used an action-angle framework to develop expressions for both the quasilinear diffusion tensor $\mathbf{D}$ as well as the linear susceptibility K , valid for a range of unperturbed magnetic geometries (including slabs and tokamaks), and for arbitrary mode structure. Application of the formal expression for $\mathbf{D}$ to tokamak transport problems was made in Refs. 1,2,12-14. The structure of K , of principal interest here, is quite similar. Both are given by a sum over all harmonics $\mathbf{l} \equiv (l_g, l_b, l_\zeta)$ of the gyro-, bounce-, and toroidal precession frequencies ($\Omega_g, \Omega_b, \Omega_\zeta$), with an 'orbit-averaging factor' $G_{\mathbf{l}}^2$, whose effect is central here, weighting the summand. Axisymmetry requires that $l_\zeta = n$, and one may take only $l_g = 0$ components for the low frequency modes considered here. As usual, we parametrize the tokamak by the radial, poloidal, and toroidal coordinates (r, θ, ζ) . Simplifying somewhat for purposes of brevity and physical clarity, we model the mode structure as locally a plane wave with wave vector $\mathbf{k} = \hat{r}k_r + \hat{\theta}m/r + \hat{\zeta}n/R$. Then for species s , $g_s \equiv k^2 \lambda_s^2 K_s$ is approximately given by (λ_s is the Debye length)

$$g_s(\mathbf{k}, \omega) = 1 - \sum_l \langle J_0^2(z_g) J_l^2(z_b) \frac{\omega - \omega_{*s}^l}{\omega - \sigma k_{\parallel} \bar{v}_{\parallel} - (1 - \sigma)\omega_{ds} - i\Omega_b + i\nu_{fs}} \rangle. \quad (1)$$

Here, $l \equiv l_b - \sigma m$, with trapping state index $\sigma = 0$ for trapped particles, and $\sigma = 1$ for passing particles, $\langle \dots \rangle \equiv \int d\mathbf{v} (f_0/n_0) \dots$ is the velocity space average over the unperturbed distribution function f_0 (species label suppressed), $\omega_{*s}^l \equiv \omega_* [1 + \eta(u^2 - 3)/2]$, with $\omega_* \equiv -k_y c T / (e B L_n)$ the diamagnetic drift frequency, $\eta \equiv d \ln T / d \ln n$, $u \equiv v/v_s$ the particle velocity, normalized to the thermal speed v_s , $k_y \equiv \hat{b} \times \hat{r} \cdot \mathbf{k}$, and $L_n^{-1} \equiv -\partial \ln n_0 / \partial r$. $\bar{v}_{\parallel}$ is the bounce-averaged parallel velocity, $\omega_d = n \Omega_\zeta (\sigma = 0)$ is the toroidal

drift frequency, and $\nu_f \simeq \nu/\epsilon$ is the effective collision frequency, with ν the 90-degree scattering frequency. For the chosen model of mode structure, G_l^2 is given by the product of the squares of Bessel functions J^2 appearing in (1), with one ($J_g^2 \equiv J_0^2(z_g)$) due to gyro-orbit-averaging, with argument $z_g = k_\perp \rho_g$, and the other ($J_b^2 \equiv J_l^2(z_b)$) with argument z_b , due to averaging over the drift orbit. The full expression^{2,14} for z_b is slightly more complicated than that for z_g , but the essential physics is captured by approximating it as $z_b \simeq k_\perp \rho_b$. The principal modifying effect on the dispersion relation held in Eq. (1) arises from the finite-drift factor J_b^2 for ions.

For electrons, the treatment is standard. One has $z_{g,b} \rightarrow 0$, so that $G^2 \rightarrow \delta(l)$, as usual. Making the standard assumption $\xi_e \equiv \omega/k_\parallel v_e \ll 1$, the nonadiabatic portion of the electron response is dominated by trapped particles: $g_e = 1 - i\alpha_e$, where $i\alpha_e \equiv \epsilon^{1/2} \langle (\omega - \omega_{ce}^I)/(\omega - \omega_{ce} + i\nu_{fe}) \rangle_t$, with $\langle \dots \rangle_t$ the average over the trapped portion of the distribution.

For ions, if one makes the conventional assumption that $\omega/\Omega_b \gg 1$, then over the l -range $\Delta l \sim z_b$ over which J_l^2 in Eq. (1) is appreciable, the denominator there does not change greatly, so that one can perform the summation, using the identity $\sum_l J_l^2 = 1$, which eliminates the J_b^2 factor, leaving only the usual factor J_g^2 . In actuality, however, for the modes considered here, one has $\omega/\Omega_b \gtrsim 1$, so that only a small number of bounce harmonics, having $\omega/|l|\Omega_b \gtrsim 1$, contribute appreciably to the sum,¹⁵ and the effect of J_b^2 survives. For simplicity, we assume that only the $l = 0$ term satisfies this for passing particles, which is often the case, and the resultant scaling is similar to the more general situation where a few low- l harmonics must be kept. When this is the case, since $\Omega_b(\sigma = 0)/\Omega_b(\sigma = 1) \simeq \epsilon^{1/2}$, a few low- l harmonics for trapped particles contribute, reducing the drift averaging effect

for $\sigma = 0$. Thus, here we neglect the drift averaging for $\sigma = 0$, noting that at somewhat lower frequencies ($\omega/\Omega_b(\sigma = 0) \sim 1$) it will be present for $\sigma = 0$ as well as for $\sigma = 1$, somewhat strengthening the overall effect. For the single $\sigma = 1$ term we assume that the usual condition $\xi \equiv \xi_i \equiv \omega/k_{\parallel}v_i > 1$ is (marginally) satisfied, and neglect ω_{di} and ν_{fi} in (1). Expanding expression (1) to third order in ξ^{-1} , one may evaluate g_i in terms of velocity space moments of G^2 , such as $\Lambda_0 \equiv \langle J_g^2 J_b^2 \rangle$ and $\Lambda_2 \equiv \langle (u_{\perp}^2/2) J_g^2 J_b^2 \rangle$. In the driftless (slab) limit $z_b \rightarrow 0$ (with z_g fixed), $J_b^2 \rightarrow 1$, the parallel and perpendicular integrations in Eq. (1) decouple, and can be exactly evaluated, in terms of $\Lambda_0^{(0)}(b_g) \equiv \langle J_g^2 \rangle = I_0(b_g)e^{-b_g^2}$, $\Lambda_2^{(0)}(g_g) \equiv \langle (u_{\perp}^2/2) J_g^2 \rangle = (I_0 + b_g I_0')e^{-b_g^2}$, and the plasma dispersion function $\langle (u_{\parallel} - \xi)^{-1} \rangle$. Here, $I_0(b)$ is the modified Bessel function of the first kind, and $b_g \equiv k_{\perp}^2 \rho_{gi}^2 \equiv k_{\perp}^2 v_i^2 / \Omega_{gi}^2$.

While ρ_g and ρ_b have the same energy scaling ($\propto E^{1/2} \propto v$), the pitch-angle dependence of ρ_b is more complicated than that for ρ_g , coupling the parallel and perpendicular integrations. Therefore, to capture the essential effect of finite z_b , we adopt an approximate form for Λ_0 , $\Lambda_0(b_g, b_b) \simeq (1 + b_c)^{-1} \equiv (1 + k_{\perp}^2 \rho_c^2)^{-1}$ and analogous approximations for the higher moments such as Λ_2 . Here, $b_c \equiv b_g + b_b$, with $b_b \equiv k_{\perp}^2 \rho_{bi}^2 \simeq k_{\perp}^2 (q\rho_{gi})^2$ the bounce-analogue of b_g . This choice properly models the slab behavior $\Lambda_0(b_g \ll 1, b_b = 0) = \Lambda_0^{(0)}(b_g) \simeq 1 - b_g$, with a similar linear behavior for $b_g = 0, b_b \ll 1$, and the limit $\Lambda_0(b_g, b_b \rightarrow \infty) \rightarrow 0$. Our treatment of g_i will thus be to evaluate it exactly in the driftless limit $b_b \rightarrow 0$ (but b_g arbitrary), and then to replace the $\Lambda^{(0)}$ s which appear by the corresponding Λ s, which incorporate the effect of finite b_b in the approximate fashion just indicated.

Thus, expanding the dispersion function to third order in ξ^{-1} , quasineu-

trality yields the dispersion equation

$$\begin{aligned} 0 &= g_e + \tau g_i \\ &= (1 - i\alpha) - \frac{\omega_{se}}{\omega} + \tau(1 - \Lambda_0)\left(1 + \frac{\omega_{se}}{\omega} K_3\right) - \frac{k_{\parallel}^2 c_s^2}{\omega^2} \Lambda_0 \left(1 + \frac{\omega_{se}}{\omega} K_4\right), \end{aligned} \quad (2)$$

where $c_s^2 \equiv \tau v_i^2$ as usual, use has been made of $\tau\omega_{si} = -\omega_{se}$, and $\alpha \equiv \alpha_e + \tau\alpha_i$, with α_e and $\alpha_i \simeq (\omega_{si}/\omega - 1)\Lambda_0\sqrt{\pi/2}\xi \exp(-\xi^2/2)$ containing the resonant contributions from electrons and ions. α_i is usually neglected under the $\xi \gg 1$ assumption. Eq. (2) is written in a form which transparently reduces to the small- b_g limit usually considered. The first term is $g_e - i\tau\alpha_i$, and the remaining three are the rest of τg_i . In the third and fourth terms, $K_3 \equiv [1 + \eta_i(-\delta\Lambda_2/(1 - \Lambda_0))]/\tau$ and $K_4 \equiv [1 + \eta_i(\Lambda_2/\Lambda_0)]/\tau$, where $\delta\Lambda_2 \equiv \Lambda_2 - \Lambda_0$. For $b_g \ll 1$, $K_3, K_4 \rightarrow K \equiv (1 + \eta_i)/\tau$, and $(1 - \Lambda_0), \delta\Lambda_2 \rightarrow b_g$, and one recovers the standard dispersion relation.⁸

Eq. (2) contains both the TEM and the (slablike branch) of the ITG mode, both being drift waves, destabilized by different driving mechanisms.¹⁶ For the simplest version of the TEM, one sets $\eta_i = 0 = k_{\parallel}$, resulting in a mode with frequency $\omega = \omega_r + i\gamma$, with real and imaginary parts

$$\omega_r = \omega_{se}\Lambda_0/[1 + \tau(1 - \Lambda_0)], \gamma = \alpha\omega_{se}\Lambda_0/[1 + \tau(1 - \Lambda_0)]^2. \quad (3)$$

For the fluidlike limit ($\tau > 1, b_g \ll 1$) commonly considered, the denominator in γ governs the falloff of γ with $k_{\perp}$. Taking $b_b = 0$ and $\Lambda_0 \simeq 1 - b_g$, as is conventional, one has $\gamma \simeq \alpha\omega_{se}/(1 + b_g^2)^2$, where $b_g \equiv \tau b_g \equiv k_{\perp}^2 \rho_s^2$. (Fig. 2.) Thus, the range in $k_{\perp}$ of γ (and accordingly, one expects, of the spectrum) is given by $\Delta k_{\perp} \simeq \rho_s^{-1}$ which scales as $T_e^{-1/2}$:

$$\gamma(\tau \gtrsim 1, b_b = 0) \propto k_y/(1 + k_{\perp}^2 \rho_s^2)^2.$$

In the same fluidlike limit, but with b_b finite, $\Lambda_0 \simeq 1 - b_c$, so that b_s in the above formula should be approximately replaced by $b_t \equiv \tau b_c \equiv k_\perp^2 \rho_i^2$. $\Delta k_\perp \simeq \rho_i^{-1}$ thus still scales as $T_e^{-1/2}$.

However, for $\tau \ll 1$, as is the case in supershots, $\Delta k_\perp$ is instead determined by the factor Λ_0 in the numerator of γ , whose falloff depends upon T_i , and is more rapid for $\tau < 1$ (Fig.2):

$$\gamma(\tau < 1) \propto k_y \Lambda_0 \simeq k_y / (1 + k_\perp^2 \rho_c^2).$$

The corresponding transport is given by the expression for the particle flux $\Gamma_s = \sum_{\mathbf{k}} (ck_y/B) I_{\mathbf{k}} (e_s n_{0s}/T_s) \text{Im}(g_s)$, with a similar expression for the heat flux. Here, $I_{\mathbf{k}} \equiv \langle |\phi_{\mathbf{k}}|^2 \rangle$ is the intensity of the turbulent spectrum. Using g_i in Eq. (2), this gives, for the ions, $\Gamma_i = -D_i n'_{0i}$, where

$$D_i \simeq \sum_{\mathbf{k}} (ck_y/B)^2 I_{\mathbf{k}} \Lambda_0 \gamma / \omega_r^2. \quad (4)$$

From Eq. (3), the factor $\Lambda_0 \gamma / \omega_r^2$ here equals α / ω_{*e} , and as a result orbit-averaging does not enter the shown expression for D_i explicitly, but only through its modification of $I_{\mathbf{k}}$, and perhaps of α .

Obtaining $I_{\mathbf{k}}$ correctly through a full nonlinear calculation is difficult, and is not attempted here. The usual shortcut is to invoke the 'mixing-length prescription' $(e/T)^2 \sum_{\mathbf{k}} I_{\mathbf{k}} \simeq 1/(\Delta k_\perp L_n)^2$, which applied to Eq. (4) for the conventional $\tau \gtrsim 1$ theory, for which $\Delta k_\perp \simeq \rho_s^{-1}$, yields $D \simeq \gamma / \Delta k_\perp^2 \simeq \alpha D_0$, where $D_0 \equiv \rho_s^2 c_s / L_n$ is the 'gyro-reduced Bohm' expression, scaling as $T_e^{3/2}$. Including nonzero b_b in the $\tau \gtrsim 1$ theory, for which $\Delta k_\perp \simeq \rho_i^{-1}$, results in $D \simeq \alpha D_0 (\rho_i / \rho_s)$. The factor $(\rho_i / \rho_s) = (\rho_c / \rho_{gi}) \simeq (1 + q^2)^{1/2}$ introduced in going from the slab theory to its toroidal modification is reminiscent of a similar factor introduced in going from classical to neoclassical

transport, and introduces an additional dependence $\simeq q \propto B/I_p$ into D_i or χ_i . As is well known, the scaling $D, \chi \propto T_e^{3/2}/B^2$ characteristic of conventional drift-wave theories (of which the $b_i = 0, \tau \gtrsim 1$ TEM theory just discussed is an example) conflicts with the experimentally observed scaling with I_p and B , and must be supplemented by a factor of roughly $q^{2.5}$ to be brought into agreement.¹⁹ The additional factor of q which orbit averaging introduces into the $\tau \gtrsim 1$ result thus partially remedies one of the major failings to date of drift wave theories. A similar modifying factor has been found by Guo, *et al.*²⁰ for the ITG mode in going from slab to toroidal geometry, due not to the drift-averaging considered here (which is not present, due to the more standard frequency ordering made there), but to toroidal effects in the ballooning equation, which are not considered here.

To consider the modification of this result in the $\tau < 1$ regime, we assume that the saturated amplitude of I_k is proportional to a positive power p of γ , reflecting the intuition that the saturation amplitudes should become smaller as the term driving them diminishes. While the power implied by the simplest mixing-length argument just noted is $p = 0$, contradicting this intuition, weak turbulence calculations suggest that p should instead equal unity¹⁷, while scaling arguments for strong turbulence suggest¹⁸ $p = 2$. Here we take $p = 1$. Then one has

$$\frac{D(\tau < 1)}{D(\tau > 1)} \simeq \frac{\int dk k_y / (1 + k_\perp^2 \rho_e^2)}{\int dk k_y / (1 + k_\perp^2 \rho_i^2)^2} \simeq \frac{\rho_i^{2r}}{\rho_e^{2r}} = \tau^r, \quad (5)$$

where the exponent r , lying between 1 and 2, depends upon the scaling one assumes for the spectral widths in $k_\parallel$ and k_r . The scaling from the simple analytic theory ($k_\parallel \sim k_v, k_r \sim k_v^0$) gives $r = 3/2$. Thus, while the TEM scaling for L-mode discharges is $D(\tau > 1) \propto T_e^{3/2}$, for supershots,

one expects $D(\tau < 1) \propto T_e^{3/2} \tau^\tau$. Therefore, if T_e increases only weakly as T_i increases, as is the case for the supershot data in Fig.1, one expects the inverse scaling of D_i or χ_i with T_i .

For the ITG mode, one may neglect α in Eq. (2), but now finite $k_{\parallel}$ is essential, and must be determined by converting the local dispersion relation (2) into the appropriate nonlocal mode equation, as in Ref. 8. As discussed in Ref. 8, the largest growth rate occurs when the second term in Eq. (2) about cancels the term in K_3 , hence for $k_{\perp}$ satisfying $1 \simeq \tau(1 - \Lambda_0)K_3$, reducing to $1 \simeq k_{\perp}^2 \rho_s^2 K$ in the usual slab treatment, or $1 \simeq k_{\perp}^2 \rho_i^2 K$ for $b_b \neq 0$. Thus, the growth rate γ peaks around $k_{\perp} \rho_i \simeq K^{-1/2}$ (sketched in Fig.2), with value given by

$$\Omega \equiv \omega/\omega_{se} \simeq (\frac{\rho_i}{\rho_s})^{1/2} (\frac{\Lambda'_0}{\Lambda_0})^{1/4} \left[\frac{\Lambda_0 K_4 \tilde{s}}{[1 + \tau(1 - \Lambda_0)]} \right]^{1/2},$$

with $\tilde{s} \equiv (2l + 1)s$, and l the radial mode number, where for the fastest growing mode l decreases with increasing s so as to keep $\tilde{s}$ about constant.⁸

Since $\Delta k_{\perp} \simeq K^{-1/2} \rho_i^{-1} = (1 + \eta_i)^{-1/2} \rho_e^{-1}$ is small, one may take $\Lambda_0 \simeq 1$, $K_4 \simeq K$ in Ω . And since it depends only on T_i for all values of τ , the transition from a large- τ to a small- τ regime evidenced by the TEM mode does not occur for the ITG mode. Using the given expressions for Ω or γ and $\Delta k_{\perp}$ for the single ITG regime, we summarize the overall discussion, by displaying the expressions for χ_i including the orbit-averaging effects, for the two TEM regimes, and the single ITG regime, followed by the approximate scalings of these results with T_e , T_i , B , and I_p :

$$\begin{aligned} \chi_i(TEM, \tau \gtrsim 1) &\simeq (\rho_i/\rho_s) \alpha D_0 \sim T_e^{3/2} B^{-1} I_p^{-1} \alpha \\ \chi_i(TEM, \tau < 1) &\simeq (\rho_i/\rho_s) \alpha \tau^\tau D_0 \sim T_e^{3/2+\tau} T_i^{-\tau} B^{-1} I_p^{-1} \alpha \quad (6) \\ \chi_i(ITG) &\simeq (\rho_i/\rho_s)^{3/2} K \tilde{s}^{1/2} D_0 \sim T_e^{1/2} T_i B^{-1/2} I_p^{-3/2}. \end{aligned}$$

For all three results, one notes that the 'neoclassical factor' $(\rho_i/\rho_s) \simeq (1+q^2)^{1/2} \sim q$ improves the agreement with the experimental scaling¹⁹ $\chi_{exp} \sim T_e^{3/2} B^{-2} q^{5/2}$ over the scaling $\chi(b_0 = 0) \sim T_e^{3/2} B^{-2}$ characteristic of earlier drift-wave theories, with the ITG mode showing the strongest improvement. The factor is a generic aspect of the theory, not specific to the particular modes considered here. As already noted, the two TEM regimes display the scaling with τ needed to explain the difference between L-mode and supershots exhibited in Fig.1, while the single ITG scaling does not. A preliminary TFTR database test of the expected TEM scaling for $r < 1$, made using the supershot discharges of Fig.1, yields a T_i dependence $\chi_i \sim T_i^{-r}$, with $r \simeq 1.44$. However, the remaining T_e dependence appears weaker than the $T_e^{3/2+1}$ in Eq. (7), suggesting that the dependence $D_0 \sim T_e^{3/2}$ from the usual TEM theory should be weaker, or that transport is controlled by some other mode with a two τ -regime behavior similar to that manifested by the TEM here, but with weaker T_e scaling.

We have included in the TEM scalings the factor α . The scaling of this factor remains to be more carefully explored, but further refinements of existing TEM theory have been suggested¹⁹ which might, along with the ballooning effect pointed out in Ref. 20 noted above, introduce the needed remaining q dependence into a more complete theory.

These results suggest a reinterpretation of the cause of the good confinement and the very high ion temperatures in TFTR supershots, which had been previously attributed to the stabilization of ITG due to peaked density profiles.²¹ The present result implies the possibility that the supershot central ion temperatures 'run away,' (at least to the equipartition limit of the average beam energy) due to the $\chi_i \propto T_i^{-r}$ scaling of the orbit-averaged

TEM transport for $\tau < 1$, and perhaps also that the peaked density profiles are caused by the high- T_i orbit-averaging effect on D . These conjectures still need to be further tested against the experimental databases, and against numerical solutions containing the orbit-averaging effect.¹⁶

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Figure Captions

Figure 1. Experimental measurements of χ_i versus T_i on TFTR, for L-mode and supershot discharges.

Figure 2. Sketch of growth rate γ versus $k_{\perp}\rho_s$, for the single- τ regime of π mode, for the usual TEM theory ($\tau \gtrsim 1, b_b = 0$), and for the modification of that result by orbit averaging, in the $\tau \gtrsim 1$ and $\tau < 1$ regimes.

ION TRANSPORT IN DIFFERENT ENERGY REGIMES

Fig. 1

Fig. 2