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Short Wavelength Electron Temperature Gradient Driven Drift Wave Turbulence

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Short wavelength drift-wave turbulence driven by electron-temperature gradient

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A local, hydrodynamic model of the short wavelength electrostatic drift waves driven unstable by the electron temperature gradient in toroidal geometry is used to find the saturation level and the mode coupling to the longer wavelength collisionless skin depth c/ω_{pe} magnetic turbulence. For plasma with $\beta_e > 2\, m_e/m_i$ the k-spectrum is peaked in the c/ω_{pe} wavelength region at the mixing length amplitude and the magnetic turbulence is sufficient to produce the empirical neo-Alcator and Goldston type of confinement formulas.

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The increase of the electron confinement time with plasma density, major and minor radius described by a large quantity of tokamak data is given by the empirical confinement laws of either the neo-Alcator[1] or the Goldston[2] laws. Attempts to explain the observed confinement have largely centered on the dissipative trapped electron mode, $k_\perp \rho_i < 1$ scale turbulence. Here ρ_i and ρ_e are the ion and electron gyroradius, respectively. Ohkawa[3] first observed with a simple argument that electromagnetic fluctuations with characteristic scale of the collisionless skin depth $\delta = c/\omega_{pe}$ may be the relevant step size in an electron random walk at the rate associated with transit frequency of the electrons in the toroidal trap: the resulting diffusivity $\chi_e \sim \frac{c^2}{\omega_{pe}^2} \frac{v_{the}}{qR}$ leads to a reasonable first approximation for the magnitude of the empirical scaling. Furthermore, electron transport scaling[4] studies in reversed field pinches (RFP's) show reasonable agreement with T_e versus current scaling provided $k_\| \sim 1/L_s \sim B_\theta/rB_z \sim 1/a$ for the RFP. The amplitude of the magnetic fluctuations are given by $\delta B_r/B \simeq k_\| \delta$ where $\delta = c/\omega_{pe}$ with $k_\| = 1/qR$ for tokamaks and $k_\| \sim 1/a$ for the RFP's.

In the present Letter we develop for the tokamak system the suggestion of Guzdar *et al.*[5] that the nonlinear development of short wavelength drift wave instabilities driven by the electron temperature gradient η_e may be the source of the c/ω_{pe} turbulence. The instability is a short wavelength version of the toroidal ion temperature gradient instability[6] except that the roles of the ions and electrons are reversed so that the source of fluctuation energy is at $k_\perp \rho_e \gtrsim 0.5$. The linear toroidal kinetic stability theory giving the threshold value of η_e and the scaling of the growth rate and wavenumbers for both the $\eta_e \sim 1$ and the large η_e regimes are given by Horton *et al.*[7] In the limit of strong shear and weak toroidicity the instability reduces to the sheared slab form given by Lee *et al.*[8]

Here we show that the toroidal η_e instability evolves due to the influence of the $\mathbf{E} \times \mathbf{B}$ convective nonlinearity and the $\delta \mathbf{B} \cdot \nabla$ magnetic field line nonlinearity to a saturated state of electromagnetic drift wave turbulence. The amplitude of the fluctuations in the saturated state is given by the mixing length formulas obtained by balancing the mode coupling terms with the dominant linear terms driving the instability. We find that the three-dimensional form of the equations is important to obtain the correct level of the magnetic fluctuations. Although the magnetic fluctuations are small in the short wavelength region of the fluctuation spectrum near the maximum linear growth rate,[7,8] and thus the direct $\mathbf{E} \times \mathbf{B}$ transport is smaller by $(m_e/m_i)^{1/2}$ than the ρ_s scale drift wave transport, we find that as the nonlinear mode coupling terminates the exponential growth phase of the instability, there is a transfer of fluctuation energy to the wavelengths of order c/ω_{pe} as shown in Figs. 1 and 2. After the break of the exponential growth phase during which the small scale damping is unimportant,[9] there is a slow transfer of energy along the spectrum to shorter (ρ_e^{-1}) wavelengths. The spectral features that determine the anomalous transport, however, are established by the longer wavelength components produced during the breaking of the exponential growth phase.

The equations describing the toroidal η_e instability used in this Letter are based on the hydrodynamic electron equations assuming an adiabatic ion density response $(k_\perp \rho_i > 1)$, quasineutrality $\tilde{n}_e = \tilde{n}_i = -n_0(e\Phi/T_i)$, and the conservation of particles, parallel electron momentum and electron thermal energy in the region on the outside of the torus. The equations give a self-consistent field model for predicting the type of electromagnetic fluctuation spectrum of the type assumed in the earlier studies.[3–8] The derivation of the equations and comparison of the linear hydrodynamic modes with linear Vlasov theory is given in Horton *et al.*[7]

In the dimensionless variables of $(x,y) \to \rho_{ei}(x,y)$, $z \to r_n z$ and $t \to r_n t/v_{ei}$ with $v_{ei} = (T_i/m_e)^{1/2}$, $\rho_{ei} = v_{ei}/\omega_{ce}$ and the scaling of the amplitude of the fields by

$$\left(\frac{\tilde{T}_e}{T_e}, \frac{v_{\|e}}{v_{ei}}, \frac{e\Phi}{T_e}, \frac{2}{\beta_i}\frac{ev_{ei}}{cT_i} A_\|\right) = \frac{\rho_{ei}}{r_n}(\tilde{T}_e, v, \phi, A)$$

which makes the low-β dependence of $\delta B_x = i\, k_y\, A_\|$ explicit. The nonlinear equations for the coupling of the short wavelength electrostatic fluctuations to the longer wavelength electromagnetic fluctuations are as follows:

$$\left(1-\nabla_\perp^2\right)\frac{\partial\phi}{\partial t}=\left(1-2\epsilon_n(1+\tau)+\tau(1+\eta_e)\nabla_\perp^2\right)\frac{\partial\phi}{\partial y}+2\epsilon_n\frac{\partial\widetilde{T}_e}{\partial y}+\tau[\phi,\nabla_\perp^2\phi]$$
$$+\frac{1}{\tau}\left(\frac{\partial}{\partial z}\nabla_\perp^2 A-\frac{\beta_i}{2}[A,\nabla_\perp^2 A]\right)+d_c\nabla_\perp^2\phi \qquad (1)$$

$$\left(\nabla_\perp^2-\frac{\beta_i}{2}\right)\frac{\partial A}{\partial t}=\tau(1+\tau)\frac{\partial\phi}{\partial z}-\tau\frac{\beta_i}{2}[A,(1+\tau)\phi]-\tau\frac{\partial\widetilde{T}_e}{\partial z}+\tau\frac{\beta_i}{2}[A,\widetilde{T}_e]$$
$$+\tau\frac{(1+\eta_e)}{2}\beta_i\frac{\partial A}{\partial y}-\tau[\phi,\nabla_\perp^2 A]-\eta\nabla_\perp^2 A \qquad (2)$$

$$\frac{\partial\widetilde{T}_e}{\partial t}=-\tau\left[\eta_e-2\epsilon_n(1+\tau)(\Gamma-1)\right]\frac{\partial\phi}{\partial y}-2\tau\epsilon_n\left(2\Gamma-1\right)\frac{\partial}{\partial y}\widetilde{T}_e$$
$$-(\Gamma-1)\left(\frac{\partial}{\partial z}\nabla_\perp^2 A-\frac{\beta_i}{2}\left[A,\nabla_\perp^2 A\right]\right)$$
$$-\tau[\phi,\widetilde{T}_e]+\chi_\perp\nabla_\perp^2\widetilde{T}_e+\chi_\parallel(\partial_\parallel^{n\ell})^2\widetilde{T}_e \qquad (3)$$

where $\Gamma=5/3$ in fluid theory. In Eqs. (1)–(3) we include electron crossfield diffusion, resistivity and electron thermal conductivity to absorb energy transformed to $|k|=k_\perp\rho_{ei}\gg 1$ outside the range of validity of the fluid equations. Using the classical transport coefficients, the dimensionless dissipation coefficients are $d_c=(\nu_{ei}r_n/v_{ei})\,(T_e/T_i)$, $\eta=0.51\,(\nu_{ei}r_n/v_{ei})$, $\chi_\perp=4.67\,(\nu_{ei}r_n/v_{ei})\,(T_e/T_i)$ and $\chi_\parallel=3.16\,(v_{ei}/\nu_{ei}\,r_n)\,(T_e/T_i)$.

Including the energy conserving nonlinear FLR terms $[\widetilde{T}_e,\nabla^2\phi]+[\partial_x\widetilde{T}_e,\partial_x\phi]+[\partial_y\widetilde{T}_e,\partial_y\phi]$ in Eq. (1) gives a more complete nonlinear model of the turbulence with a slightly reduced transport rate. In Eq. (3) the parallel thermal flux requires the use of the nonlinear derivative $\partial_\parallel^{n\ell}=\partial_z-(\beta_i/2)[A,\quad]$ and $\chi_\parallel\gg 1\gg\chi_\perp$.

In deriving Eqs. (1)–(3) it is assumed that the dominant nonlinearities are the $\mathbf{E}\times\mathbf{B}$ convective derivative and the $\mathbf{B}\cdot\nabla$ magnetic gradient due to the perturbations in the magnetic field lines. In writing these two nonlinearities we use the Poisson bracket operator $[f,g]$ defined by

$$\delta\mathbf{B}\cdot\nabla f=-\hat{\mathbf{z}}\cdot\nabla A\times\nabla f\equiv-[A,f]\quad,\quad \mathbf{v}_E\cdot\nabla g=\hat{\mathbf{z}}\cdot\nabla\phi\times\nabla g\equiv[\phi,g] \qquad (4)$$

with the property $\overline{[f,g]}_{y,z} = -\partial_x\overline{\left(\frac{\partial f}{\partial y}g\right)} = \partial_x\overline{\left(f\frac{\partial g}{\partial y}\right)}$ where the bar denotes the y,z average. We define the volume average by $\langle F\rangle = V^{-1}\int d^3x\, F(x,y,z,t) = L_x^{-1}\int dx\, \bar{F}(x,t)$ and note the properties $\langle h[f,g]\rangle = \langle f[g,h]\rangle = \langle g[h,f]\rangle$.

The crossfield correlation functions,[9,10] $Q_{es}(x)$ and $Q_{em}(x)$ giving the transport of thermal energy are defined by

$$Q_{es}(x) = \frac{3}{2}\overline{\frac{\partial\phi}{\partial y}\tilde{T}_e}\left[\frac{cT_e}{eB}\frac{\rho_{ei}}{r_n}\frac{n_e}{\eta_e}\frac{dT_e}{dx}\right]$$

$$Q_{em}(x) = \frac{3}{2}\frac{\beta_i}{2\tau}\overline{\chi_{\|e}\left(\frac{\beta_i}{2}\eta_e\frac{\partial A}{\partial y} - \nabla_\|\tilde{T}_e\right)\frac{\partial A}{\partial y}}\left[\frac{cT_e}{eB}\frac{\rho_{ei}}{r_n}\frac{n_e}{\eta_e}\frac{dT_e}{dx}\right] \tag{5}$$

with $Q_{em}(x)$ analyzed by Hong & Horton[10] in the quasilinear limit. The three fields energy components are

$$E_\phi = \frac{1}{2}\left\langle\phi^2 + (\nabla\phi)^2\right\rangle \;,\; E_A = \frac{1}{2}\left\langle(\nabla^2 A)^2 + \frac{\beta_i}{2}(\nabla A)^2\right\rangle \;,\; E_T = \frac{1}{2}\left\langle(\tilde{T}_e)^2\right\rangle \;. \tag{6}$$

The total energy $E_{\text{Tot.}} = E_\phi + \frac{E_A}{\tau^2(1+\tau)} + \frac{E_T}{\tau(1+\tau)(\Gamma-1)}$ grows and decays according to

$$\frac{dE_{\text{Tot.}}}{dt} = \eta_e\langle Q_{es}\rangle - d_c\left\langle(\nabla_\perp\phi)^2\right\rangle - \frac{\eta}{\tau^2(1+\tau)}\left\langle(\nabla_\perp^2 A)^2\right\rangle - \frac{\chi_\perp}{\tau(1+\tau)(\Gamma-1)}\left\langle(\nabla_\perp\tilde{T}_e)^2\right\rangle$$
$$-\frac{\chi_\|}{\tau(1+\tau)(\Gamma-1)}\left\langle(\nabla_\|\tilde{T}_e)^2\right\rangle \;. \tag{7}$$

Since enstrophy is not conserved by the system there is a power transfer in the fluctuation spectrum to both high k-modes and low k-modes.

Now we give the magnitudes of the nonlinear terms and the mixing length level of saturation for the rms amplitudes of the fields. In the electrostatic limit $\beta_e \lesssim 2\,m_e/m_i$ the equations are the same structure as the toroidal η_i equations where the mixing length level of saturation was shown analytically[6] and numerically[9] and the resulting Q_{es} is greater than the neoclassical plateau value but small compared with the ρ_s scale χ_e.

In the nonlinear regime there are two mixing rates: (1) the $\mathbf{E}\times\mathbf{B}$ mixing frequency Ω_E determined by the time for convection around the vortex given by $\Phi_{\mathbf{k}}(x,y)$ and (2) the spatial rate of mixing $k_{\|}^{n\ell}$ determined by the distance along $\mathbf{B}_0$ required to go around the magnetic vortex given by $A_{\|\mathbf{k}}(x,y)$. The $\mathbf{E}\times\mathbf{B}$ mixing rate and $k_{\|}^{n\ell}$ magnetic mixing rate are given by

$$\Omega_E = \frac{ck_x k_y}{B}\Phi_{\mathbf{k}} = \left(k_x r_n \frac{e\Phi_{\mathbf{k}}}{T_e}\right)\omega_{*e} \qquad \text{and} \qquad k_{\|}^{n\ell} = \frac{k_x k_y}{B} A_{\mathbf{k}} \ . \tag{8}$$

The mixing produces[10,11] stochastic $\mathbf{E}\times\mathbf{B}$ transport at the saturation level $\Omega_E \sim |\omega|$ and which the stochastic magnetic transport when $k_{\|}^{n\ell} \sim k_{\|}^{0}$. For MHD-like fluctuations with $E_{\|} \simeq 0$ the mixing rates are related by $k_{\|}^{n\ell} \simeq k_{\|}^{0}(\Omega_E/\omega)$. The end of the linear regime and "wave breaking" occurs at the saturation levels $|\omega| \sim \Omega_E$ and $k_{\|}^{0} \sim k_{\|}^{n\ell}$ yielding the mixing length fluctuation amplitudes

$$\frac{e\Phi}{T_e} \sim \frac{1}{k_x r_n}\left|\frac{\omega}{\omega_*}\right| \qquad ; \qquad \frac{\delta B_x}{B} \simeq \frac{k_{\|}^{0}}{k_x} \ . \tag{9}$$

Test electron orbits in tokamaks with such levels of electromagnetic drift wave fluctuations shows global stochasticity with transport well described by the diffusion approximation.[11]

For η_e, ϵ_n well past their threshold values[5,7,8] the linear modes grow until the mixing rates are sufficiently rapid to essentially eliminate the electron pressure gradient over the width π/k_x of the fluctuation. Thus we estimate for the saturated state that

$$\mathbf{v}_E \cdot \nabla(\overline{T}_e + \widetilde{T}_e) \simeq 0 \qquad \text{and} \qquad \mathbf{B}\cdot\nabla(\overline{T}_e + \widetilde{T}_e) \simeq 0 \ , \tag{10}$$

in the strongly turbulent state. Both conditions in Eq. (10) determine the same mixing length level of the temperature fluctuation

$$\left\langle \widetilde{T}_e^2 \right\rangle^{1/2} = \widetilde{T}_e = \frac{1}{|k_x|}\left|\frac{dT_e}{dx}\right| \tag{11}$$

eliminating the driving mechanism. Using $\Omega_E \sim |\omega_k| \sim \gamma_k$ and the rate of magnetic mixing when $k_{\|}^{n\ell} \simeq k_{\|}^{0}$, then the ratios of the nonlinear fluctuations are consistent with the linear fluctuations[7] equations of Eqs. (1)–(3). For example, using the dominant terms in Eq. (3)

and Eq. (11) yields

$$\frac{e\Phi_k}{T_e} \simeq \frac{\omega}{\omega_{*Te}} \frac{\widetilde{T}_e}{T_e} = \frac{\gamma_k}{\omega_{*e}} \frac{1}{k_x r_n} \cong \frac{1}{|k_x|} \left(\frac{2}{R r_{T_e}} \right)^{1/2} \tag{12}$$

consistent with $\Omega_E \simeq \gamma_k$. Using level (12) for Φ_k and Eq. (2) gives for $A_{\|k}$

$$\frac{eA_{\|k}}{T_e} \simeq \left| \frac{1 - \omega_{*pe}/\omega}{1 - \frac{\omega_{*pe}}{\omega} + \frac{c^2 k_\perp^2}{\omega_{pe}^2}} \right| \left| \frac{c k_\|}{\omega_{*e}} \right| \frac{1}{k_x r_n} \simeq \frac{c k_\|^0}{\omega_{*e}\, k_x r_n} \frac{1 + \eta_e}{\left[(\eta_e / 2\epsilon_n)^{1/2} + c^2 k_\perp^2 / \omega_{pe}^2 \right]} \tag{13}$$

which is consistent with $k_\|^{n\ell} \simeq k_\|^0$.

From numerous simulations on the CRAY 2 and the Fujitsu VP200, the x-y contour plots of $\phi(\mathbf{x}, t), A(\mathbf{x}, t)$ and $\widetilde{T}_e(\mathbf{x}, t)$ in the turbulent state are shown in Fig. 1 for the parameter values $\epsilon_n = 0.1$, $\eta_e = 1.0$, $\beta_e = 0.01$, $\tau = 1$ and $d_c = \eta = \mathcal{X}_\perp = 0.1$ and $\mathcal{X}_\| = 10$. The contours (Fig. 1(b)) of potential fluctuation, which are the streamline of $\mathbf{E} \times \mathbf{B}$ drift motion, show that the $\mathbf{E} \times \mathbf{B}$ flows are chaotic. The magnetic fluctuations, which are initially small and random, saturate into large scale ($\gtrsim c/\omega_{pe}$) magnetic vortex structures formed in the steady state (Fig. 1(c)). The linear growth rate has a maximum at $k_x \rho_{ei} \sim k_z r_n \cong 0$ and $k_y \rho_{ei} \cong 0.8$. During saturation the spectrum evolves into the isotropic state peaked at $k_x \rho_{ei} \sim k_y \rho_{ei} \sim (\beta_i/2)^{1/2} \cong 0.05$ to 0.1.

In Fig. 2 the wavenumber spectrum of $E_{\text{Tot.}}$ as a function of k_x summed over k_y, k_z in panel (a) and as a function of k_y summed over k_x, k_z, is shown for two time values in the steady state. The high $k_\perp$ contributions is larger than that given for the 2D η_i-mode theory ($k_\perp^{-3}$) in Refs. 6,9 due to the stronger 3D effects in these η_e simulations. The isotropic, long wavelength fluctuation spectrum shown in Fig. 2 has direct implications for the $\mathcal{X}_e$ formulas due to the stochastic diffusion of electrons over the correlation scale c/ω_{pe}. Turbulent energy on space scales larger than that shown in Fig. 1 falls into the $k_\perp \rho_i < 1$ regime of η_i-modes where nonadiabatic ion behavior applies, which modifies Eq. (1)–(3).

The electrostatic component of the energy flux shown in Fig. 3 follows approximately from Eqs. (5), (11), and (12)

$$Q_{es} \simeq -\frac{3}{2} n_e \frac{\rho_{ei}}{r_n} \frac{cT_e}{eB} \frac{\langle k_y \rho_{ei} \rangle \left[2\varepsilon_n \eta_e\right]^{1/2}}{\langle k_x^2 r_n^2 \rangle} \frac{dT_e}{dx} \tag{14}$$

which is smaller by $(m_e/m_i)^{1/2}$ than the $\mathbf{E} \times \mathbf{B}$ transport produced by ρ_s scale turbulence. Here $\langle k_y \rangle$ and $\langle k_x^2 \rangle$ are averages over the fluctuation spectrum.

In terms of the dimensionless A field we have $B_x^2/B_0^2 = (\beta_i \rho_{ei}/2r_n)^2\, k_y^2 |A_k|^2$. Now using Eq. (13) to evaluate $|A_k|^2$ and from Eq. (5) for $Q_{em}(x)$ we find that for a turbulent spectrum with $c^2 \overline{k}_\perp^2/\omega_{pe}^2 < |1 - \omega_{*pe}/\omega| \approx (\eta_e/2\varepsilon_n)^{1/2}$

$$\chi_e^M \cong \frac{3}{2} \left[\varepsilon_n\, \eta_e\, \beta_e\right]^{1/2} \frac{v_e}{r_n} \frac{c^2}{\omega_{pe}^2} \tag{15}$$

for $T_e \simeq T_i$. As noted earlier, the simple model (1)–(3) used here does not give an accurate threshold for η_e [see Ref. 7] so that Eq. (15) applies for $3 \gtrsim \eta_e > \eta_{e,\text{crit}} \approx 1$. For larger η_e the rescaled, large η_e equations[7] are required and the higher k_z slab modes[8] become important.

The calculation shows that the condition on the plasma pressure $\beta_e \gtrsim 2(m_e/m_i)$ is required for the magnetic transport to compete with electrostatic $\mathbf{E} \times \mathbf{B}$ transport. The Guzdar *et al.*[5,8] formula for the sheared ($s = rq'/q$) slab analysis is $\chi_e = 0.13 c^2 s\, \eta_e (1 + \eta_e)/\omega_{pe}^2\, qR$.

Studies with c/ω_{pe} transport formulas show that the scaling and magnitude of electron thermal confinement in both the Ohmic heated tokamak[5] and RFP.[4] The auxiliary heated tokamaks can be interpreted with formula (15) or closely related formulas. Taking into account[11] both the ρ_s scale turbulence and the c/ω_{pe} turbulence Yu *et al.*[12] are able to reproduce several (five) features of the observed dependence of the energy replacement time τ_E on the system parameters. Eliminating all long wavelength ($> \rho_s$) turbulence still leaves a substantial anomalous χ_e according to Eqs. (14) and (15).

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Figure Captions

1. Contours of electrostatic potential $\phi(\mathbf{x}, t)$, magnetic flux function $A(\mathbf{x}, t)$, and temperature fluctuation $\widetilde{T}_e(\mathbf{x}, t)$. (a) The linear regime potential and (b)–(d) fields in the turbulent state.

2. Steady state wavenumber spectrum. (a) The k_x spectrum and (b) the k_y spectrum.

3. The steady-state $\mathbf{E} \times \mathbf{B}$ heat flux component versus the temperature gradient parameter η_e.

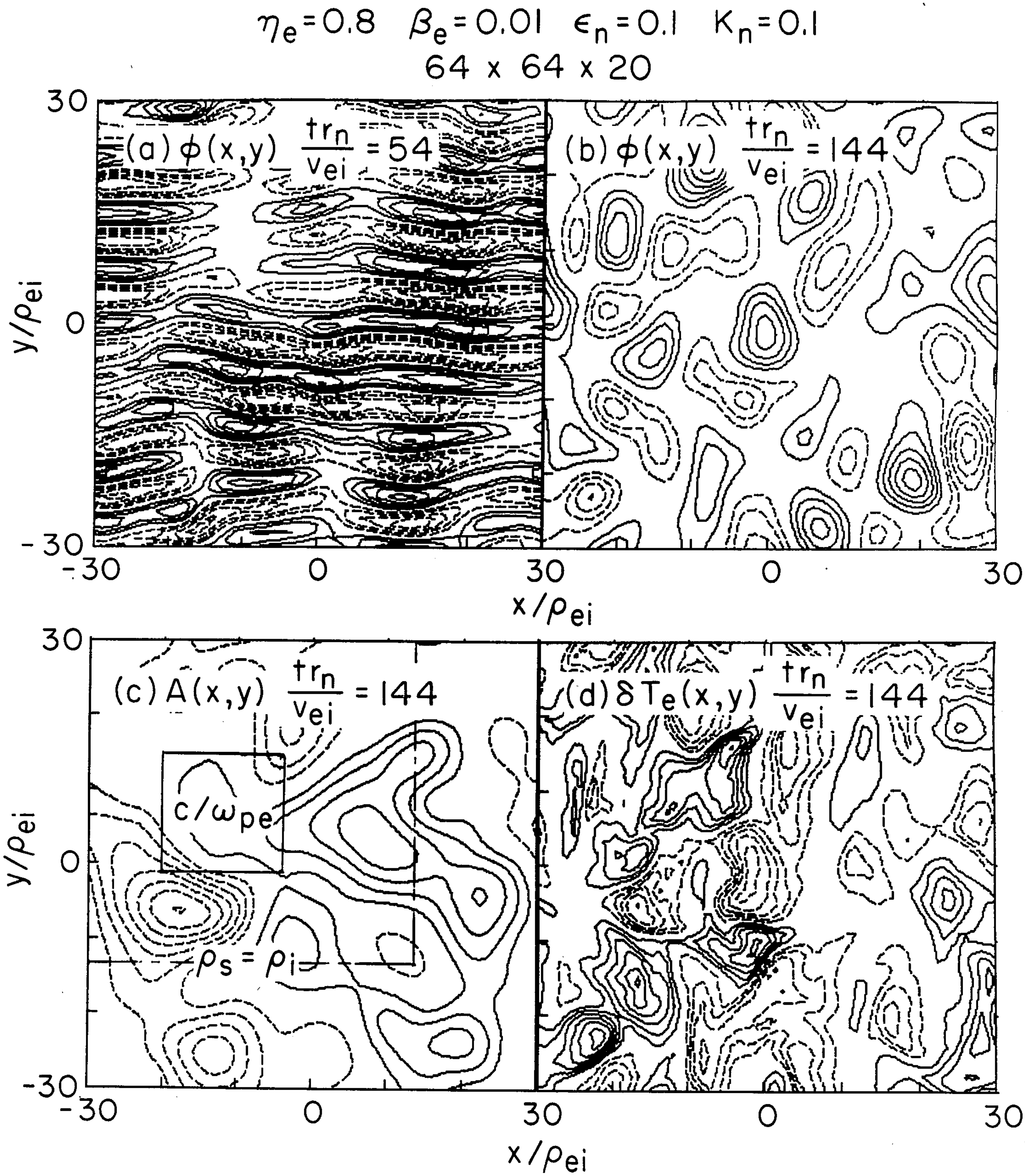
$\eta_e = 0.8$ $\beta_e = 0.01$ $\epsilon_n = 0.1$ $K_n = 0.1$
64 x 64 x 20
(a) $\phi(x,y)$ $\frac{tr_n}{v_{ei}} = 54$
(b) $\phi(x,y)$ $\frac{tr_n}{v_{ei}} = 144$
(c) $A(x,y)$ $\frac{tr_n}{v_{ei}} = 144$
c/ω_{pe}
$\rho_s = \rho_i$
(d) $\delta T_e(x,y)$ $\frac{tr_n}{v_{ei}} = 144$
y/ρ_{ei}
x/ρ_{ei}
30
0
-30

Fig. 1

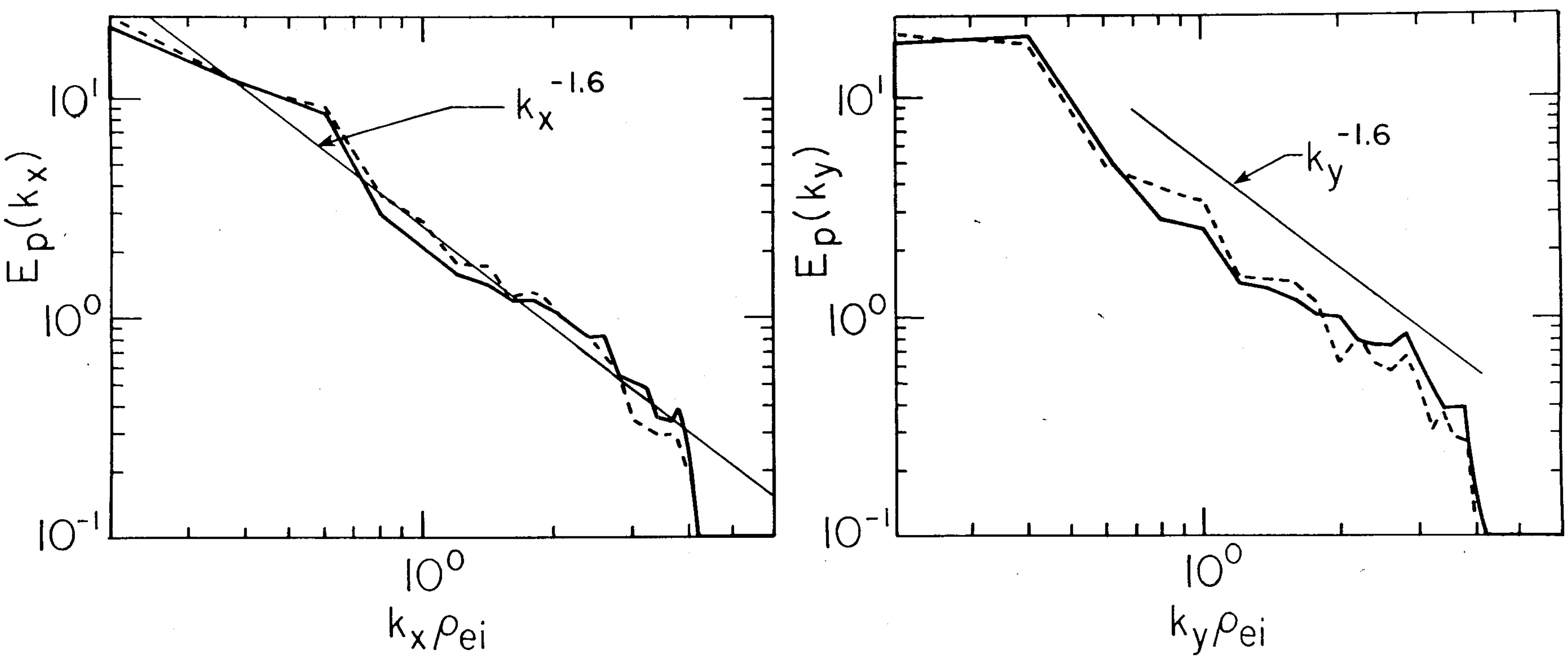
$E_p(k_x)$
10^1
10^0
10^{-1}
$k_x^{-1.6}$
10^0
$k_x \rho_{ei}$
$E_p(k_y)$
10^1
10^0
10^{-1}
$k_y^{-1.6}$
10^0
$k_y \rho_{ei}$

Fig. 2

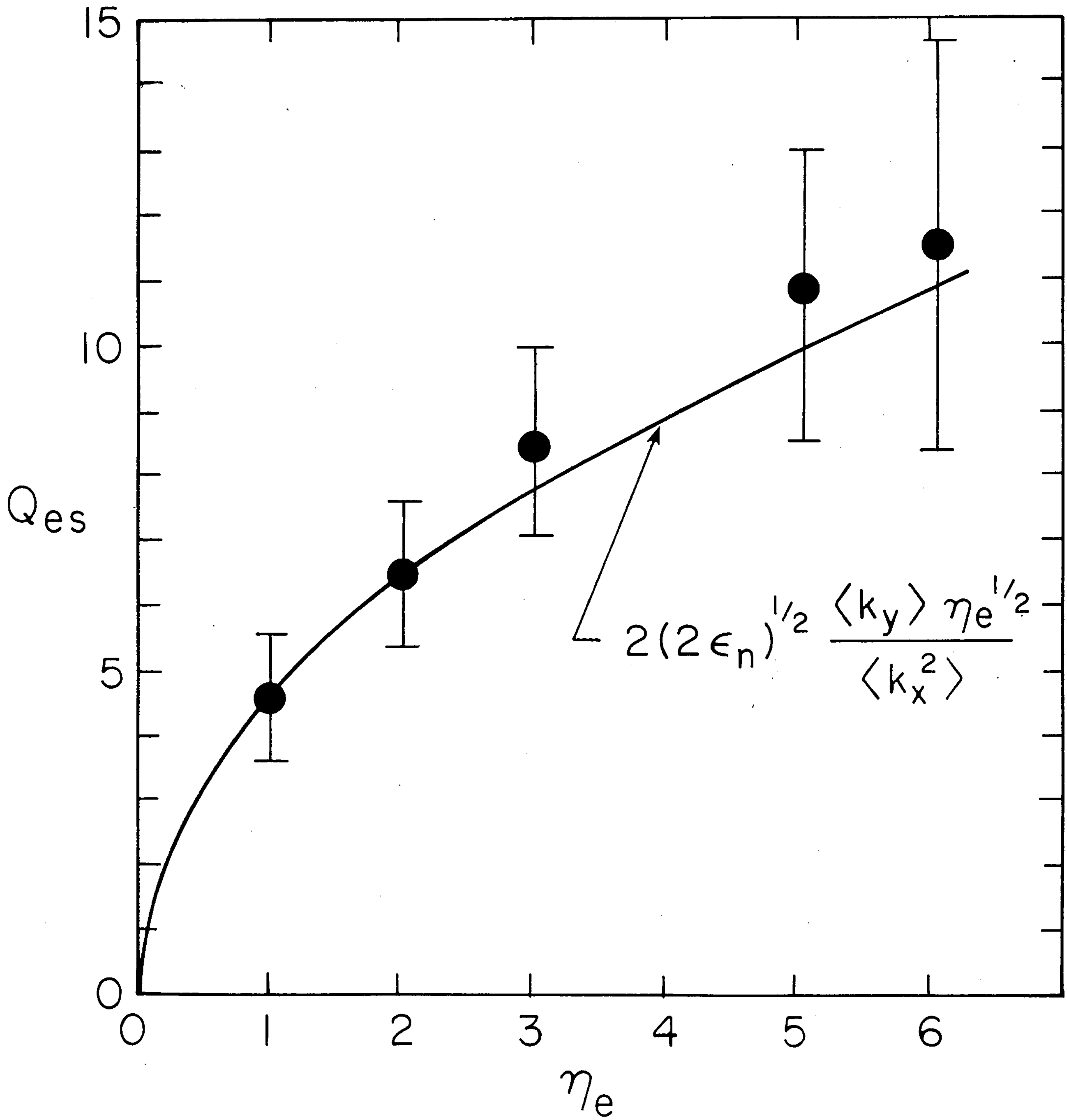
15
10
5
0
0
1
2
3
4
5
6
Q_{es}
η_e
$2(2\epsilon_n)^{1/2} \frac{\langle k_y \rangle \eta_e^{1/2}}{\langle k_x^2 \rangle}$

Fig. 3