any 2^1S-state population in the case of cesium, in obvious disagreement with our result.

We conclude that, in addition to the practical interest of knowing the composition of metastable He beams, the present results should challenge theorists to improve the description of the collision mechanisms involved in these processes.

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Finite-β_e Universal-Mode Turbulence and Alcator Scaling

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A self-consistent theory of finite-β_e universal-mode turbulence is developed. Saturation results from resonance broadening of the electron response due to magnetic shear. Electron diffusion, for $\beta_e > m_e/m_i$, is due to the magnetic part of the fluctuations. The diffusion coefficient, $D = 0.1\,[T_e/(T_i+T_e)]^4\,(m_e/m_i\beta_e)(L_s/L_n)^2\,v_i\rho_i^2/L_n$, scales inversely with density, is independent of magnetic field, and is in excellent quantitative agreement with observations on the Alcator tokamak.

One of the principal theoretical goals in tokamak research is the development of a self-consistent turbulence theory for the short-wavelength fluctuations through to be responsible for anomalous transport. This paper presents a nonlinear turbulence theory for the finite-β_e universal instability. An approximate analytic solution of the coupled, nonlinear, eigenmode equations is obtained. The accuracy of this solution has been verified by numerical computations. The resulting formula for the anomalous electron thermal conductivity, Eq. (9), has many similarities with experimental observations, including absolute magnitude, and scaling with density, temperature, and ion mass. For typical tokamak regimes where $\beta_e > m_e/m_i$, the calculation constitutes an example of a self-consistent theory of stochastic[1] magnetic fluctuations.

Until recently,[2] most turbulence theories ignored shear in the equilibrium magnetic field. Without shear, turbulence mainly affects the ions. The basic saturation picture, as developed by Dupree,[3] balanced linear electron growth, γ_e^{L}, against nonlinear (turbulent) ion damping, γ_i^{NL}. Taking $\gamma_i^{\mathrm{NL}} = k_\perp^2 D$ is the basis for the $\gamma_e^{\mathrm{L}}/k_\perp^2$ estimates of the anomalous diffusion coefficient. However, recent theory[4] has found that $\gamma_i^{\mathrm{NL}} \ll k_\perp^2 D$, because the ion-wave interaction is weak for low-frequency modes, and consequently, for tokamak parameters, that ion nonlinearity is not a viable saturation mechanism.

Because of their rapid mobility along the field lines, electrons are the species most strongly affected by turbulence in a sheared magnetic field.[5] Here, the ions cause linear damping, due to the shear, at a rate γ_i^S. Electron growth is modified by shear-induced resonance broadening, to γ_e^{NL}. Saturation, $\gamma_e^{NL}=\gamma_i^S$, can occur at low turbulence levels, consistent with observations.

We do not consider the existence of instability to be in doubt, but there has been some confusion on this point. All recent linear theories have been unanimous in predicting stability and our recent paper[2] showed instability through nonlinear effects. The key physical effect uncovered by the nonlinear theory, however, as we have emphasized recently,[6] is spatial diffusion in combination with streaming along the lines in the presence of shear. The amount of diffusion necessary for destabilization is so small that it would be unlikely not to be present in practice. Specifically, classical collisional scattering of the guiding centers, $D\sim\nu_{ei}\rho_e^2$, which is three orders of magnitude below *observed* electron diffusion, destabilizes the mode.

In a sense, the main point of this paper is to apply shear-induced resonance broadening to the electromagnetic problem and show that it provides an effective mechanism for the saturation of the instabilities. The specific example considered here is basically a drift wave, although finite β_e does modify the shear damping in an important way, and when $\beta_e>m_e/m_i$, most of the transport is due to the magnetic part of the fluctuation.

We consider a cylindrical tokamak, with equilibrium field components $\vec{B}=B_0\vec{e}_z+B_\theta(r)\vec{e}_\theta$. Since $\beta_e\equiv 4\pi nT_e/B^2\ll 1$, the compressional mode of fluctuations may be neglected, and it suffices to consider the parallel component of the vector potential $\tilde{A}_\parallel$. Expanding the modes in a Fourier series in poloidal and toroidal angle $\sim\exp(im\theta-in\varphi-i\omega t)$, the parallel wave vector is given by $k_\parallel\equiv(m-nq)/Rq$, where $q=rB_0/RB_\theta$ is the safety factor. For each mode, the rational surface is at r_{mn}, such that $q(r_{mn})=m/n$, and we let $x=r-r_{mn}$ be the distance from the rational surface. Then $k_\parallel=k_\parallel'x\equiv k_\theta x/L_s$. As in linear theory the modes have a definite parity with respect to x. The drift-wave parity considered here is $\tilde{\Phi}$ even, $\tilde{A}_\parallel$ odd, so that $\hat{E}=-ik_\parallel(\tilde{\Phi}-\tilde{\Psi})$ is odd, where $\tilde{\Psi}$, defined as $\tilde{A}_\parallel\omega/k_\parallel c$, is even in x.

The nonlinear electron response to $\tilde{\Phi}$ and $\tilde{\Psi}$ may be computed in the manner indicated previously.[2] Writing the electron fluctuation as

$$\tilde{f}_e=(e/T_e)F_M[\tilde{\Phi}-(1-\omega_{*e}/\omega)\tilde{\Psi}]+\tilde{h}_e, \tag{1}$$

where the first term is the adiabatic response (or the $k_\parallel v_\parallel\to\infty$ limit of the drift kinetic equation), the result of this "renormalization" is equivalent to computing $\tilde{h}_e$ from the diffusion equation[6]

$$[-i(\omega-k_\parallel v_\parallel)-Dd^2/dx^2]\tilde{h}_e = (e/T_e)F_M\, i(\omega-\omega_{*e})(\tilde{\Phi}-\tilde{\Psi}). \tag{2}$$

Now D contains magnetic as well as electrostatic contributions,

$$D\simeq\sum_{m,n}\frac{c^2k_\perp^2}{B^2}\left|\tilde{\Phi}_{mn}-\frac{v_\parallel}{c}\tilde{A}_{\parallel mn}\right|^2 R(\omega-k_\parallel v_\parallel), \tag{3}$$

where $R=\int_0^\infty dt\,\exp[i(\omega-k_\parallel v_\parallel)t-\frac{1}{3}(k_\parallel'v_\parallel)^2Dt^3]$ is a turbulently broadened resonance function. Since $\tilde{A}_\parallel\simeq(c/v_A)\tilde{\Phi}$ [see Eq. (10)], where $v_A=(B^2/4\pi nm_i)^{1/2}$, note that when $v_e^2/v_A^2>1$ or $\beta_e>m_e/m_i$, diffusion due to the magnetic part of the fluctuations is dominant. The physical properties of shear-induced resonance broadening can be inferred directly from Eq. (2).[7]

Equation (2) is appropriate when the particle orbits in the presence of the fluctuations exhibit the stochasticity property, thus justifying diffusive behavior on the spatial scale of a wavelength. This will happen when the island-overlap (or Chirikov) condition is satisfied. For poloidal mode numbers in the 10^2 range which, we find, dominate the spectrum, island overlap occurs for electrostatic fluctuations of order $e\tilde{\varphi}_{mn}/T_e\sim10^{-4}$, or magnetic fluctuations of $\tilde{B}_{rmn}/B\sim10^{-7}$. Thus for these high-mode-number fluctuations the rational surfaces are packed very densely together and the stochasticity condition is, for all practical purposes, always satisfied.

Turning to the eigenmode problem, the electron density and current fluctuations can be computed from Eq. (2). Finite-gyroradius ion fluctuations are computed from linear theory. Combining these, using quasineutrality $\tilde{n}_e=\tilde{n}_i$ and Ampere's law, $-\nabla_\perp^2 \tilde{A}_\parallel (4\pi/c)(\tilde{J}_{\parallel e}+\tilde{J}_{\parallel i})$, leads to

$$\frac{d'}{d}\frac{d^2\Phi}{dx^2}-\xi\Phi=(\Phi-\psi)\left(\Lambda-\xi-\mu^2x^2+\frac{\sigma_e(|x|)}{|x|}\right), \tag{4}$$

$$\left(\frac{d^2}{dx^2}-b\right)\frac{k_\parallel c}{\omega}\psi=(\Phi-\psi)\frac{c^2}{v_A{}^2}\frac{\omega d}{k_\parallel c}\left[\Lambda-\xi-\mu^2x^2+\frac{\sigma_e(|x|)}{|x|}\left(1+i\frac{\omega_c}{\omega}\right)\right]. \tag{5}$$

We have retained the standard notation wherever possible, defining

$$b=k_\perp{}^2\rho_i{}^2,\quad \Gamma_n(b)=e^{-b}I_n(b),\quad \tau=T_e/T_i,\quad x_e=\omega/k_\parallel{}'v_e,\quad x_c=\omega_c/k_\parallel{}'v_e,\quad d=(\Gamma_0-\Gamma_1)(\tau+\omega_{*e}/\omega),$$

$$\xi=(1-\Gamma_0)/(\Gamma_0-\Gamma_1),\quad \Lambda=d^{-1}[1+\tau(1-\Gamma_0)-\Gamma_0\omega_{*e}/\omega],\quad \mu^2=d^{-1}(\tau+\omega_{*e}/\omega)\Gamma_0(\omega_{*e}/\omega\tau)^2(L_n/L_s)^2,$$

$$k_\parallel{}'\equiv k_\perp/L_s,\quad d'=d+1.2i(1-\omega/\omega_{*e})(\omega_{*c}/\omega_c)x_c{}^2,\quad \sigma_e(|x|)/|x|=(1-\omega_{*e}/\omega)d^{-1}(x_e/|x|)Z((x_e+ix_c)/|x|).$$

Note that the underlying electromagnetic modes persist in the presence of stochasticity, at least for the odd parity of $A_\parallel$.

To simplify the solution to Eqs. (4) and (5), we use an approximate algebraic relation between Φ and ψ (obtained by interpolating between the asymptotic relations) in Eq. (4) to give a self-adjoint equation for Φ. The approximation is based on the following argument.[8] We are concerned with what is basically a drift wave, and the principal dynamics are described by Eq. (4). The primary finite-β effect (as is known from previous work[9]) is the reduction of the parallel electric field, or $\Phi-\psi$, within the Alfvén layer, $x_A=\omega/k_\parallel{}'v_A$, due to inductive effects. To represent this algebraically, note that $\psi\to\Phi$ follows from the $x\to 0$ limit of Eq. (5). The large-x behavior is found by combining (4) and (5) to give

$$d^2\Phi/dx^2-\xi\Phi=(k_\parallel c/\omega d)(v_A{}^2/c^2)(d^2/dx^2-b)\,k_\parallel c\psi/\omega.$$

As $x\to\infty$, the derivative terms are dominant, implying $\psi\sim\Phi x_A{}^2 d/x^2$, and so by a simple interpolation, $\psi\cong\Phi x_\beta{}^2/(x^2+x_\beta{}^2)$, where $x_\beta{}^2=x_A{}^2d$. The approximate eigenmode equation is then obtained from Eq. (4):

$$\frac{d^2\Phi}{dx^2}-\Phi\frac{x^2(\Lambda-\mu^2x^2+\alpha_e)+\xi x_\beta{}^2}{x^2+x_\beta{}^2}=0. \tag{6}$$

Equation (6) passes to the usual electrostatic equation[2] as $x_\beta\to 0$.

A quadratic form

$$S=\int dx\,\{(d\varphi/dx)^2+\varphi^2(d'/d)[x^2(\Lambda-\mu^2x^2+\alpha_e)+\xi x_\beta{}^2]/(x^2+x_\beta{}^2)\}$$

whose first variation gives Eq. (1) may be constructed in the usual way. Since we are interested in the dispersion relation near saturation, we pass to the limit $\omega\tau_c\equiv\omega[\frac{1}{3}(k_\parallel{}'v_e)^2D]^{-1/3}<1$ of the electron response,[2] or $\alpha_e=i\frac{1}{3}\Gamma(\frac{1}{3})d^{-1}(\omega-\omega_{*e})\tau_c$. This limit, as confirmed by the numerical solutions, gives the upper (saturated) marginal stability point. Taking the normalized trial function to be $\Phi=(\alpha/\pi)^{1/4}\exp(-\alpha x^2/2)$, with α the variational parameter, and doing the integrals for $\alpha^{1/2}x_\beta<1$, yield $S=\frac{1}{2}\alpha-\mu^2/2\alpha+\pi^{1/2}\xi x_\beta\alpha^{1/2}+\Lambda+\alpha_e$, which again reproduces the electrostatic form when $x_\beta\to 0$. The parameter α is determined by

$$\delta S/\delta\alpha=0=\tfrac{1}{2}+\mu^2/2\alpha^2+\pi^{1/2}x_\beta\xi\alpha^{-1/2}/2.$$

Dominance of the first two terms gives the Weber trial function. However, when $\pi x_\beta{}^2\xi^2>\mu$, or $\beta_e>(L_n/L_s)^3(1+\tau)^2/\tau\pi^{1/2}$, which is typically satisfied for tokamaks, the last two terms dominate and we obtain $\alpha^{1/2}=e^{i\pi/3}(\mu^2/\pi^{1/2}\xi x_\beta)^{1/3}$ as the variational parameter. Here $e^{i\pi/3}$ is chosen uniquely as the cube root of -1, by applying the outgoing-wave boundary condition. The dispersion relation now follows as

$$S=0=\tfrac{3}{4}(\pi^{1/2}\xi x_\beta\mu)^{2/3}+\Lambda+i\left[\tfrac{3}{4}^{3/2}(\pi^{1/2}\xi x_\beta\mu)^{2/3}+\tfrac{1}{3}\Gamma(\tfrac{1}{3})d^{-1}(\omega-\omega_{*e})\tau_c\right]. \tag{7}$$

Note that the primary finite-β_e effect is to modify the shear damping. Finite x_β pushes out the eigenmode, enhancing the outward convection of wave energy; the resultant damping involves a geometric mean of x_β and μ.

At marginal stability, the frequencies are real, and the real and imaginary parts of Eq. (7) give the frequency, $\omega=\omega(k_\perp)$, and diffusion coefficient, $D=D(k_\perp)$, respectively. Here $D(k_\perp)$, which is interpreted as the amount of diffusion necessary to stabilize mode $k_\perp$, is given by

$$D=0.07\left(\frac{T_e}{T_e+T_i}\right)^4\frac{m_e}{m_i\beta_e}\left(\frac{L_s}{L_n}\right)^2\frac{v_i\rho_i^2}{L_n}b^{1/2}. \qquad (8)$$

This is not yet the transport coefficient since it depends on the wave parameter $k_\perp$. Actually ion collisions, as well as higher-order spatial turbulent broadening effects and ion nonlinearities,[10] act to give additional damping for $b>1$, with the consequence that $D(b)$ will have a maximum near $b_{\max}=2$. The dependence of $b_{\max}$ on plasma parameters affects the scaling of D very weakly. For convenience, in practical applications, we simply put $b=2$ in Eq. (8), giving the formula quoted in the abstract, which coincides with the values obtained by numerical solutions to (4) and (5).

The D of Eq. (8) is an electron-test-particle diffusion coefficient. When inserted in the electron kinetic equation, appropriately using the ambipolar potential,[11] one can obtain transport equations in the usual way. The transport equations in Ref. 11 apply to the present case when $\beta_e > m_e/m_i$. For crude estimates, D equals the electron thermal conductivity χ_e and particle diffusion is not anomalous (there is no fast ion motion along the field lines).

Rewriting (8) to display the scaling (with lengths in centimeters and temperatures in electron volts, μ the ion mass in units of the proton mass) gives

$$D=2.65\times10^{16}\frac{1}{n\mu^{1/2}}\left(\frac{T_e}{T_e+T_i}\right)^4\frac{T_i^{3/2}}{T_e}\left(\frac{L_s}{L_n}\right)^2\frac{1}{L_n}, \quad (9)$$

which is independent of magnetic field. For $T_e = T_i$, D increases with temperature, but if the dependence on $\tau=T_e/T_i$ is considered, (9) is consistent with recent observations on PLT (Princeton Large Torus),[12] where confinement improved with increased T_i. The behavior with ion mass, $D\propto\mu^{-1/2}$ is in agreement with ISX-A (Impurity Study Experiment) measurements.[13] Further, Eq. (9) shows a tendency for improved confinement with decreased aspect ratio, [note that $L_s/L_n=(Rq/r)d\ln n/d\ln q$ is approximately R/a] a feature which, qualitatively, has been seen in the data.[14]

Note that Eq. (8) can be written

$$D=\frac{v_e}{L_s}\frac{c^2}{\omega_{pe}^2}\left[6\times10^{-3}\left(\frac{m_e}{m_i}\right)^{1/2}\left(\frac{L_s}{L_n}\right)^3\right].$$

This, to the extent that the factor in brackets is 1 (it is of the order for tokamaks), is the Ohkawa[15] formula. Of course, the physics underlying the present diffusion coefficient is much different, being due to low-frequency modes, as is the scaling with ion mass, the T_e/T_i ratio, q, and several other parameters.

Finally taking Alcator[16] profiles at $r=5$ cm, $L_s \cong 100$ cm, $L_n\cong5$ cm, and $T=1$ keV, and defining $\tau_E\equiv a_L^2/D$, we find $\tau_E=2.4\times10^{-19}\,na_L^2$, which is within 50% of the Alcator empirical scaling law,[14,17] $\tau_E=3.8\times10^{-19}na_L^2$. We might note that this transport mechanism is also consistent with the interpretation of the soft-x-ray anomaly given in Ref. 11, and that the absolute value of D from Eq. (8) agrees very well with that found empirically from the x-ray data.

Obviously, the theory as described in this paper is not accurate to 50%, and many omitted effects and questions about the finite-β_e eigenvalue problem remain to be considered. However, the main attributes of Eq. (8), namely, absolute magnitude and density scaling, have been verified by the numerical solution of Eqs. (4) and (5). The significant point is that, when saturated by shear-induced resonance broadening, the most fundamental drift instability, the universal mode, gives inverse density scaling of the anomalous thermal conductivity. One then feels that this feature will be retained when complicating effects are added, and that on this basis, the established empirical scaling, $\tau_E\propto n$, can be understood.

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Parametric Instabilities in Lower-Hybrid-Frequency Heating of a Tokamak

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Decay spectra during the rf heating were identified as the parametric decay into the cold lower-hybrid waves and the ion-cyclotron waves at the plasma surface. The decay instabilities absorbed some fraction of the rf power, hence they affected the plasma heating. However, they did not prevent the penetration of the lower hybrid wave to the plasma center. Drastic reduction of the decay spectra was obtained when the resonant parametric decay condition was destroyed.

Lower-hybrid heating is one of the most attractive techniques for the further heating of a tokamak plasma aiming at a thermonuclear fusion reactor, because of its engineering feasibility.[1] But physical understanding about the interaction of intense rf fields with a tokamak plasma does not seem to be sufficient. Since the required rf power is more than Ohmic heating power for the further heating, nonlinear phenomena will inevitably take place. Among them, parametric instability is the most important in the lower-hybrid range of frequency.[2] In fact, decay spectra were observed in the lower-hybrid heating of the ATC, Wega, and other tokamaks.[3-5] Prokolab has predicted that the heating in the ATC correlated with the presence of the decay spectra.[2,4] In the Wega, correlation between generation of high-energy ions and parametric decay was stressed.[5,6] It is recently discussed the existence of parametric instabilities must lead to blocking and subsequent accumulation of lower-hybrid waves on the low-density surface plasma. The role of parametric instabilities in the lower-hybrid heating experiments is still not clear, despite the enormous works on parametric instabilities.[2,4,6-10]

Lower-hybrid-heating experiments have been successfully carried out in the JFT-2 tokamak at Japan Atomic Energy Research Institute (JAERI), and efficient ion heating was obtained.[11] Parametric decay was also observed in the JFT-2. In this paper, properties of the decay spectra measured with electrostatic probes placed in the plasma periphery will be studied. Identification of the observed decay spectra in plasma heating will be presented.

A 200-kW, 650-MHz (or 750-MHz) radio-frequency system was used for lower-hybrid-heating experiments on the JFT-2. A lower-hybrid-wave launcher was a phased array of four waveguides.[1,12] The dimensions of the each waveguide at a plasma edge were 14 mm×290 mm. Experimental parameters are summarized in Table I. Schematic arrangement of the launcher is shown in Fig. 1. The front edge of the launcher was embedded in the scrape-off layer. The electron temperature and density in the scrape-off layer measured with a double probe were $T_e = 15-25$ eV and $n_s = (5-10)\times 10^{17}$ m^{-3}, respectively.

The wave signal was measured with a molybdenum coaxially shielded electrostatic probe, 3 mm long and 1 mm in diameter, immersed in the scrape-off layer. Typical frequency spectra during the rf pulse are shown in Fig. 2(a). Input frequency was 650 MHz. The low-frequency and sideband modes satisfy usual frequency conservation law and the upper sideband intensity is negligibly small. The frequency separation between the pump wave and the nth peak of lower sideband

