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A MODEL FOR THE L-H TRANSITION IN TOKAMAKS*

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ABSTRACT

Fluctuation-driven transport fluxes in the plateau regime are calculated with the methodology of neoclassical transport theory. Particle and heat fluxes are the most sensitive to fluctuations; the modification to plasma resistivity is the least sensitive. The fluctuation-driven bootstrap current and Ware pinch flux are moderately sensitive and depend on the radial mode structure. One of the thermodynamic forces depends on the radial electric field E_r. Changing E_r can change the fluctuation spectrum and thus the transport fluxes. The effects of E_r on the fluctuation spectrum are caused by the radial shear of the angular velocity, which is proportional to E_r/r. Studies of the dynamic evolution and the saturation of MHD turbulence under the influence of E_r show that the saturation amplitudes are lower and the confinement is thus better for a more negative value of E_r. A proposed model for the L-H transition is based on the improved confinement with more negative E_r. A scaling for the power threshold P_{th} is $P_{th} \propto N^3 q/(I_p^2 M_i)$, with the N plasma density, q the safety factor, I_p the plasma current, and M_i the ion mass.

1. INTRODUCTION

Traditionally, neoclassical transport theory and turbulence transport models have been developed separately. We have begun to integrate these two approaches by applying the methodology of neoclassical theory to calculate the fluctuation-driven transport matrix [1,2]. The sensitivities of the various transport fluxes to fluctuations are different. The most sensitive ones are particle and heat fluxes, which are proportional to m^2 with m the poloidal mode number. The least sensitive one is the modification to plasma resistivity, which is proportional to the parallel wave vector $|m-nq|$, with n the toroidal mode number. The bootstrap current and Ware pinch flux are proportional to m and depend on the radial mode structure. One of the thermodynamic forces depends on the radial electric field E_r. The effects of E_r cannot be transformed away by a Doppler shift in mode frequency, because a radial shear of $\vec{E} \times \vec{B}$ angular velocity $\omega_f \propto E_r/r$ exists even with constant E_r. Transport fluxes obtained depend on the fluctuation spectrum. We use the spectrum resulting from the dynamic evolution of MHD turbulence to study confinement in tokamaks—in particular, the effects of E_r on fluctuation amplitudes [3]. For rippling modes, the more negative the value of E_r, the lower the fluctuation amplitudes and, thus, the better the confinement. Experimental results from several tokamaks also indicate improved confinement with a more negative value of E_r [4–6]. On the bases of these results, we suggest that the L-H transition is triggered by a sudden change of E_r to a more negative value [7]. This change is caused by enhanced orbit losses at the edge of tokamaks when $\nu_{*i} = \nu_i Rq/v_{ti}\epsilon^{3/2}$ decreases. A scaling for the power threshold $P_{th} \propto N^3 q/(I_p^2 M_i)$ is found. We make qualitative comparisons of the model and the experimental observations.

2. NEOCLASSICAL TRANSPORT THEORY OF FLUCTUATIONS

It is shown in neoclassical theory that the bootstrap current, Ware pinch flux, and modification to plasma resistivity are induced by mirror forces, which reflect particles from magnetic and electrostatic potential wells, acting on plasma parallel flows [8]. We include such effects in the calculations of the fluctuation-driven transport matrix. The resultant electron Onsager symmetric transport matrix for electrostatic fluctuations in the plateau regime is

$$\begin{pmatrix} \Gamma_e \\ q_e/T_e \\ J_{\|}/T_e \end{pmatrix} = \begin{pmatrix} D_e & -\frac{3}{2}D_e & W \\ -\frac{3}{2}D_e & \frac{13}{4}D_e & -0.87W \\ W & -0.87W & \sigma_{e\phi}/T_e \end{pmatrix} \begin{pmatrix} X_1^e \\ X_2^e \\ E_{\|}^{(A)} \end{pmatrix} \tag{1}$$

where

$$X_1^e = -\left\langle \frac{\omega}{m} \right\rangle_e \frac{eBr}{cT_e} - \frac{e}{T_e}\Phi' + \frac{P_e'}{P_e},\ X_2^e = \frac{T_e'}{T_e},$$

$$D_e = -\frac{\sqrt{\pi}}{4}\rho_{pe}^2 \frac{v_{te}}{Rq} N \sum_{m,n,\omega} \frac{m^2}{|m-nq|} \left|\frac{e\Phi_{mn\omega}}{T_e}\right|^2,\ W = -\frac{\sqrt{\pi}}{2}\frac{c}{B_p}\frac{v_{te}}{Rq}\frac{N}{\nu_{ei}} \sum_{m,n,\omega} \frac{m(m-nq)}{|m-nq|} \left|\frac{e\Phi_{mn\omega}}{T_e}\right|^2$$

$$\sigma_{e\phi} = \sigma_s \left(1 - 0.35\sqrt{\pi}\,\frac{v_{te}/Rq}{\nu_{ei}} \sum_{m,n,\omega} |m-nq| \left|\frac{e\Phi_{mn\omega}}{T_e}\right|^2\right)$$

In Eq. (1), $E_{\|}^{(A)}$ is the inductive parallel electric field, ρ_{pe} is the electron poloidal gyroradius, B_p is the poloidal magnetic field strength, Ω_{pi} is the ion poloidal gyrofrequency, ν_{ei} is the electron-ion collision frequency, σ_s is the Spitzer conductivity, $\Phi_{mn\omega}$ is the complex amplitude for the (m,n) mode with frequency ω, and angular brackets $\langle\rangle_e$ denote the spectrum average $\langle A\rangle_e = \sum_{m,n,\omega} AS^e_{mn\omega}/\sum_{mn\omega} S^e_{mn\omega}$, where $S^e_{mn\omega} = m^2|\Phi_{mn\omega}|^2 \exp[-(\omega^E_{mn}/\omega^e_{tmn})^2]$, with $\omega^E_{mn} = \omega - mcE_r/Br$, $\omega^e_{tmn} = v_{te}|m-nq|/Rq$, and the prime denoting d/dr. To obtain the scalings of the corresponding tokamak neoclassical fluxes, we replace the spectrum summations in D, W, and $\sigma_{e\phi}$ by ϵ^2, with ϵ the inverse aspect ratio. For D two order of magnitude larger than the neoclassical value, we need $\sum_{mn\omega} m^2|e\Phi_{mn\omega}/T_e|^2/|m-nq| \gtrsim 100\epsilon^2$. However, since $m \gg |m-nq|$ for localized modes, the modification to σ_s due to fluctuations remains comparable to or less than that due to neoclassical theory. If the mode is either symmetric or antisymmetric relative to the mode rational surface where $m = nq$, W vanishes. If the mode is shifted away from the mode rational surface, $W \neq 0$ and net fluctuation-driven bootstrap currents and Ware pinch fluxes exist. The qualitative behavior of the matrix in Eq. (1) is consistent with experimental observations.

The ion transport matrix for electrostatic fluctuations in the plateau regime is

$$\begin{pmatrix} \Gamma_i \\ q_i/T_i \end{pmatrix} = \begin{pmatrix} D_i & -\frac{3}{2}D_i \\ -\frac{3}{2}D_i & \frac{13}{4}D_i \end{pmatrix} \begin{pmatrix} X_1^i \\ X_2^i \end{pmatrix} \tag{2}$$

where $$X_1^i = \left\langle \frac{\omega}{m} \right\rangle_i \frac{eBr}{cT_i} + \frac{P_i'}{P_i} + \frac{e\Phi'}{T_i} X_2^i = \frac{1}{T_i}\frac{dT_i}{dr}$$

$$D_i = -\frac{\sqrt{\pi}}{4}\rho_{pi}^2 \frac{v_{ti}}{Rq} N \sum_{m,n,\omega} \frac{m^2}{|m-nq|} \left|\frac{e\Phi_{mn\omega}}{T_i}\right|^2 \exp\left[-(\omega^E_{mn}/\omega^i_{tmn})^2\right]$$

The variables in Eq. (2) are the same as those in Eq. (1) except that the electron quantities are replaced by the corresponding ion quantities. Because of the exponential damping factor in D_i, fluctuation-driven ion fluxes cannot be enhanced over their neoclassical values as much as fluctuation-driven electron fluxes.

When E_r changes, then to maintain $\Gamma_e = \Gamma_i$ the frequency spectrum and thus the wave number spectrum need to change. If E_r becomes more negative, X_1^i (X_1^e) and, thus, Γ_i (Γ_e) may decrease (increase). To maintain $\Gamma_i = \Gamma_e$, the spectrum must readjust. Since D_i is sensitive only to the low-frequency, long-wavelength part of the spectrum, it is less sensitive than D_e to the spectrum change. To compensate for the increase of X_1^e, the fluctuation level should decrease, especially in the high-frequency, short-wavelength part of the spectrum, to reduce D_e so that Γ_e is reduced to the level of Γ_i. This kinematic argument is valid if X_1^e or X_1^i is affected by E_r, which is possible if there is a radial shear in the angular velocity $\omega_j \propto E/r$. The quantitative change can be obtained by studying the dynamic evolution of the turbulence.

3. EFFECTS OF E_r ON RESISTIVE FLUID TURBULENCE

As an example, we study the evolution and saturation of the resistivity-gradient-driven turbulence under the influence of E_r in cylindrical geometry. The model consists of the vorticity and resistivity evolution equations derived from resistive MHD theory,

$$\frac{\rho_m}{B_z^2}\frac{d}{dt}\nabla_\perp^2\Phi = -\frac{1}{\langle\eta\rangle}\nabla_\parallel^2\Phi + \frac{\langle J_z\rangle}{B_z}\nabla_\parallel\left(\frac{\tilde{\eta}}{\langle\eta\rangle}\right) \tag{3}$$

$$\frac{d}{dt}\tilde{\eta} - \chi_\parallel\nabla_\parallel^2\tilde{\eta} = \frac{1}{r}\frac{\partial\Phi}{\partial\theta}\langle\eta\rangle' \tag{4}$$

Here, $d/dt \equiv \partial/\partial t + \hat{z}\times\nabla_\perp\Phi\cdot\nabla_\perp$ is the total convective derivative, with electrostatic potential $\Phi = \langle\Phi\rangle + \tilde{\Phi}$ and resistivity $\eta = \langle\eta\rangle + \tilde{\eta}$. Angular brackets $\langle\rangle$ denote ensemble average, the tilde denotes the fluctuation part, and the remaining symbols are defined in [9]. We assume that E_r is constant in radius. The theory of resistivity-gradient-driven turbulence has been studied extensively, and it may be related to edge fluctuations in tokamaks. Following previous studies of resistive MHD turbulence theory, we derive a renormalized two-point resistivity evolution equation

$$\frac{\partial}{\partial t}\langle\tilde{\eta}\tilde{\eta}\rangle_- + \frac{(E_r - \langle r_+^2\rangle^{1/2}E_r')r_-}{\langle r_+^2\rangle}\frac{\partial}{\partial\theta_-}\langle\tilde{\eta}\tilde{\eta}\rangle_- - \frac{\partial}{\partial r_-}D_-\frac{\partial}{\partial r_-}\langle\tilde{\eta}\tilde{\eta}\rangle_-$$

$$-\frac{2\chi_\parallel}{L_s^2}\left[\Delta_r^2 + \frac{r_-^2}{4}\right]\frac{\partial^2}{\partial\theta_-^2}\langle\tilde{\eta}\tilde{\eta}\rangle_- = \langle S\rangle \tag{5}$$

where D_- is the scale-dependent relative diffusion term, $\langle S\rangle$ is the source function for driving fluctuations, the subscript $-$ denotes the relative coordinates, r_+ is the average radius, and Δ_r is the radial mode width.

The effects of E_r are introduced in Eq. (5). Even though E_r is constant in f, $E_r/r \propto \omega_f$ of the fluid element in the θ direction is not. Two fluid elements located at different radii thus rotate at different values of ω_f. The angular frequency ω_{cM} of the center of mass for these two fluid elements gives rise to the conventional Doppler shift in the frequency spectrum. The difference of ω_f relative to ω_{cM}, which is described by $E_r r_-/\langle r_+^2\rangle$ in Eq. (5), changes the correlation time τ_{cl} between fluid elements obtained for the $E_r = 0$ case when coupled to turbulent radial diffusion. The change in τ_{cl} leads to changes in the fluctuation spectrum. These effects of E_r are not specific to the rippling modes but exist for other turbulence modes as well. We see that in general the spectrum can be affected by the radial shear in ω_f, namely, $d(E_r/r)/dr$.

To determine the stationary fluctuation spectrum, we invert the evolution operator in Eq. (5) to obtain the spectrum balance equation

$$\langle \tilde{\eta}\tilde{\eta} \rangle_{-} \simeq \tau_{cl} \langle S \rangle \tag{6}$$

where $\tau_{cl} = \tau_c/(1+\sqrt{1+R_c^{-1}})(1+\tau_c^2/2\tau_f^2)\,\ell n[(1+R_c^{-1})/(DN+\sigma\tau_c k_{0r}k_{0\theta}r_-\theta_-/2\tau_f)]$, $DN = a_r k_{0r}^2 r_-^2 + a_\theta k_{0\theta}^2\theta_-^2 + a_z k_{0z}^2 z_-^2 + a$, the a are form factors for decorrelation length [9], k_0 are averaged wavenumbers, the decorrelation time $\tau_c = (Dk_{0r}^2)^{-1}$, the parallel dissipation time $\tau_d = (\chi_{\|} k_{0\theta}^2/L_s^2 k_{0r}^2)^{-1}$, the relative $\vec{E}\times\vec{B}$ time $\tau_f = [2(|E_r| - \langle r_+^2\rangle^{1/2} E_r'/\sigma)k_{0\theta}/\langle r_+^2\rangle k_{0r}]^{-1}$, and σ is the sign of E_r. These characteristic time scales represent different dynamic mechanisms in the system; $\tau_d > \tau_f > \tau_c$ for tokamak edge parameters. The Fourier transform of Eq. (6) determines the wavenumber spectrum of the saturated turbulence under the influence of E_r. The wavenumber spectrum for the negative sign of E_r (i.e., $\sigma = -1$) is narrower than that for the positive sign; consequently the saturation level is lower. The numerical calculations are carried out by using a nonlinear multiple-helicity initial value code and the evolutions of the fluctuation level with different values of E_r are shown in Fig. 1(a). We see that the saturation amplitude is lower for $\sigma = -1$. The wavenumber spectra for the two cases are shown in Fig. 1(b). The case with $\sigma = -1$ exhibits a narrower wavenumber spectrum than the other case. These results are consistent with the two-point theory and qualitative results of the kinematic argument. The E_r dependence of the calculated frequency spectra is similar to that of the wavenumber spectra.

4. A MODEL FOR THE L-H TRANSITION

Both qualitatively and quantitatively, a more negative value of E_r can reduce the fluctuation level and thereby improve confinement. Improved particle confinement with more negative E_r has also been observed in the biased limiter experiment and the Impurity Study Experiment (ISX-B). The improvement of energy confinement time τ_E with more negative E_r is not observed in these experiments, probably because of the enhanced radiation from accumulated impurities at the plasma core when E_r is more negative. On the basis of these experimental observations and theoretical studies, we suggest that the L-H transition in tokamaks is triggered by a sudden change of E_r to a more negative value, which in turn reduces fluctuation level and improves the confinement.

Before the L-H transition, E_r in the edge region (but still inside the separatrix) of a tokamak is probably determined by the nonambipolar ion particle flux Γ_{i+cx} induced by the plasma momentum loss associated with charge exchange and ionization processes,

$$\Gamma_{i+cx} = \frac{c}{eB_p} N M_i \langle \sigma v \rangle_{i+cx} N_n U_t\,, \tag{7}$$

where $\langle \sigma v \rangle_{i+cx}$ is the $\langle \sigma v \rangle$ for ionization and charge exchange processes, N_n is the neutral density, and $U_t = -(cT_i/eB_p)(e\Phi'/T_i + P_i'/P_i + \alpha T_i'/T_i)$ with α a constant. At the ambipolar state, $U_t = 0$ determines the value of E_r for the L-mode. Because of ion heating, the ion collisionality ν_{*i} decreases and ion orbit loss increases. The increasing ion orbit loss makes E_r more negative. To model the ion orbit loss, we assume that the loss cone is determined by the resonance between parallel speed $v_{\|}$ and poloidal $\vec{E}\times\vec{B}$ drift, which gives rise to the resonance condition [10] $v_{\|}/v = v_{ti}/v\,\rho_{pi}\,eE_r/T_i$, where v is the particle speed. Since $|v_{\|}/v| \leq 1$, the resonance condition can be satisfied if $v/v_{ti} \geq \rho_{pi} eE_r/T_i$. Furthermore, the ion orbit loss is most important in the banana regime, where particles with speed v satisfy $\nu_{*i}(v) \equiv \nu_i(v)Rq/(\epsilon^{3/2}v) \ll 1$. Here $\nu_i(v)$ is the ion collision frequency for particles with speed v. For $\nu_{*i}(v) \ll 1$, $v/v_{ti} \gg [\nu_{*i} \equiv \nu_{*i}(v = v_{ti})]^{1/4}$. With these two constraints on v/v_{ti}, we estimate the nonambipolar ion orbit loss rate to be

$$\left(\frac{\partial N}{\partial t}\right)_{\text{orbit}} = -N\nu_i \frac{G}{(\nu_{*i} + x^4)^{1/2}} \exp[-(\nu_{*i} + x^4)^{1/2}] \quad (8)$$

where $x = \rho_{pi} e E_r / T_i$ and G is a geometric factor that depends on the details of the loss cone boundary in the phase space. To reach the ambipolar state, we must have

$$\frac{1}{r}\frac{\partial}{\partial r}(r\Gamma_{i-cx}) = -N\,\nu_i \frac{G}{(\nu_{*i} + x^4)^{1/2}} \exp\left[-(\nu_{*i} + x^4)^{1/2}\right] \quad (9)$$

which is an equation for E_r. We approximate Eq. (9) with an algebraic equation by replacing the $\partial/\partial r$ operator with $1/\Delta r$, where Δr is the typical radial scale length in the edge region of a tokamak. The simplified equation for E_r is then

$$-(x + \lambda_i) = F_i \frac{\nu_{*i}}{(\nu_{*i} + x^4)^{1/2}} \exp\left[-(\nu_{*i} + x^4)^{1/2}\right] \quad (10)$$

where $F_i = (\Delta r/\rho_{pi})[2\nu_i/(\nu_{*i}\langle\sigma v\rangle_{i+cx} N_n)]$, and $\lambda_i = -\rho_{pi}(P_i'/P_i + \alpha T_i'/T_i)$

The qualitative behavior of the solution to Eq. (10) can be understood by examining the intersections of the two functions $y_1(x) = -(x + \lambda_i)$ and $y_2(x) = F_i\,[\nu_{*i}/(\nu_{*i} + x^4)^{1/2}] \times \exp\left[-(\nu_{*i} + x^4)^{1/2}\right]$ with x as the variable. Schematic plots for $y_1(x)$ and $y_2(x)$ are shown in Fig. 2. For $\nu_{*i} \gg 1$, or high N_n, which implies small F_i, the approximate solution to Eq. (10) is $x \simeq -\lambda_i$, which is the value of the radial electric field for the L-mode in this model. For $\nu_{*i} < 1$ and low N_n, which implies large F_i, the normalized radial electric field x in Eq. (10) is more negative than $-\lambda_i$. This is the value of the radial electric field in the H-mode. Confinement times are higher for this value of x than for $x = -\lambda_i$ in the L-mode, as discussed previously. However, for $\nu_{*i} \to 0$, the solution to Eq. (10) again approaches $x = -\lambda_i$, implying that the confinement is reduced to that of the L-mode. Therefore, a range of ν_{*i} exists over which the H-mode can occur.

For low N_n, the onset of the L-H transition occurs approximately when $\nu_{*i} \simeq 1$, which can be achieved with ion heating. If the plasma is in the L-mode and $\nu_{*i} \gg 1$ (the Pfirsch-Schlüter regime), then the direction of the ion ∇B drift relative to the null position of the divertor is important for neoclassical ion energy confinement [11]. Less power is needed to reach $\nu_{*i} \simeq 1$ if the ion ∇B drift is directed toward the null position of the divertor than if it is directed away from this position because the ion energy confinement is better in the former case.

The onset condition $\nu_{*i} \simeq 1$ gives rise to the relation for hydrogen isotopes: $NR\epsilon^{-3/2}q = CT^2$, where C is a constant. To obtain the power threshold P_{th} we assume $T_e = T_i = T$. From a typical energy confinement scaling for the L-mode [12], we have $NTV/P_{\rm th} \simeq \tau_E \propto I_p P_{th}^{-1/2} M_i^{1/2} a^{-0.4} R^{1.75}$, where V is the plasma volume, a is the minor radius, and P_{th} is the total heating power at the onset of the L-H transition. Substituting τ_E into the onset condition, we obtain a scaling for the power threshold,

$$P_{th} \propto \frac{N^3}{I_P^2}\frac{q}{M_i}\epsilon^{3.3} R^{4.3}\,. \quad (11)$$

In obtaining Eq. (11), we used a global scaling for τ_E to find a relationship between T and P_{th}. Because we are discussing the value of ν_{*i} in the edge region, we should use a local edge energy confinement time scaling, but no such scaling is available. Each tokamak generally has a unique τ_E scaling, so the power threshold P_{th} is likely to vary from device to device. For fixed N and I_p, P_{th} is proportional to the toroidal magnetic field strength B_t. If N is linearly proportional to I_p, then P_{th} does not depend on I_P. We also note that P_{th} is inversely proportional to M_i.

Several qualitative features of the neutral model for the L-H transition are consistent with the experimental results. Both high N_n and $\nu_{*i} \gg 1$ have detrimental effects on the L-H transition. If the plasma density is too low for a fixed value of T_i which implies that ν_{*i} is very much smaller than unity, the improvement in τ_p and τ_{Ee} is unlikely to occur. Because E_r is more negative in the H-mode, the impurities accumulate in the plasma core in the H-mode. One possible problem for this model is the "suddenness" of the transition. For the simplest version of the model, given in Eq. (10), the suddenness is provided by a rapid increase in T_i, which causes the corresponding rapid decrease in ν_{*i}. It is not clear whether this suddenness is adequate to explain the experimental observations. However, we note that N_n is itself coupled to the plasma transport, both in the edge region of the tokamak and in the divertor; this may provide an additional suddenness mechanism associated with the bifurcation phenomenon observed in the one-dimensional divertor model [13].

We have discussed E_r in the edge region of a tokamak. To improve τ_E, E_r must become more negative over a wide range of plasma radius. The coupling of E_r in the interior of a tokamak to E_r in the edge region, through an anomalous plasma viscosity induced by fluctuations [14], is governed by

$$\frac{\partial}{\partial t} U_t = \frac{1}{r} \frac{\partial}{\partial r} r\mu \frac{\partial U_t}{\partial r}, \tag{12}$$

where μ is a diffusion coefficient for U_t or E_r. In the fluid regime, the electrostatic fluctuation-induced momentum diffusion coefficient is

$$\mu = \frac{2}{3\sqrt{\pi}} \rho_{pi}^2 \frac{v_{ti}^2}{R^2 q^2} \sum_{m,n,\omega} m^2 \left| \frac{e\Phi_{mn\omega}}{T_i} \right|^2 \int_0^\infty dZ \frac{Z^4 \nu_i e^{-z^2}}{(\omega_{mn}^E)^2 + \nu_i^2},$$

where $Z = v/v_{ti}$. Note that μ is of the same order as the fluctuation-driven heat conductivity in the fluid regime. One way to model the L-H transition is to impose a set of boundary and initial conditions for E_r on Eq. (12). The edge boundary condition for E_r is determined by Eq. (10). When the value of E_r in the edge region changes because of changes in ν_{*i} and F_i, the solution to Eq. (12) changes and leads to improved confinement, according to Eqs (1) and (2).

The sign of E_r in the H-mode in our model is opposite to that in the Itoh-Itoh model [10]. Questions such as whether changing E_r is indeed the mechanism that triggers the L-H transition, and, if it is, which sign of E_r is releavant to the H-mode must be resolved by experimental measurement of E_r.

5. CONCLUDING REMARKS

We apply the methodology of neoclassical theory to calculate the transport matrix induced by fluctuations. We find different sensitivities to fluctuations for the various transport fluxes. The most sensitive are particle and heat fluxes, and the least sensitive is the modification to plasma conductivity. One of the thermodynamic forces depends on E_r. Changing E_r can change the fluctuation spectrum and thus the transport fluxes. The fluctuation level is lower, and confinement is thus better, for a more negative value of E_r, consistent with experimental observations. We propose a model for L-H transition based on the improvement of confinement with a more negative value of E_r. The change of E_r to a more negative value is triggered by the enhanced ion orbit loss when ν_{*i} decreases. A power threshold scaling $P_{th} \propto N^3 q/(I_p^2 M_i)$ is obtained.

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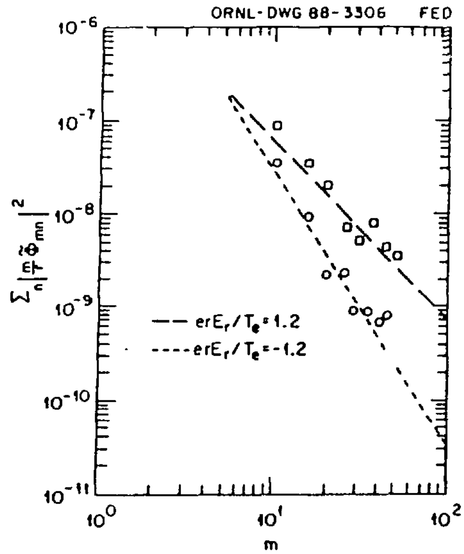
FIGURE CAPTIONS

FIG. 1.(a) Evolution of fluctuation level for three different values of erE_r/T_e. (b) Wavenumber spectrum for $erE_r/T_e = 1.2$ and -1.2.

FIG. 2. Schematic plots of the functions $y_1(x) = -(x + \lambda_i)$ and $y_2(x) = F_i[\nu_{*i}/(\nu_{*i} + x^4)^{1/2}]$ $\exp[-(\nu_{*i} + x^4)]^{1/2}$.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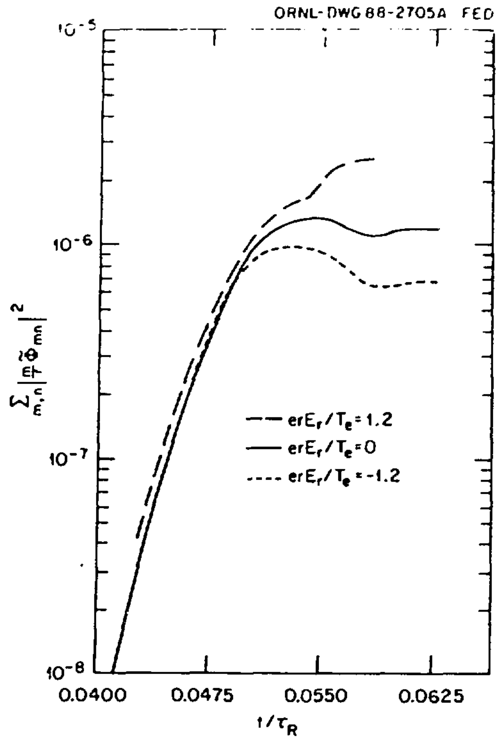
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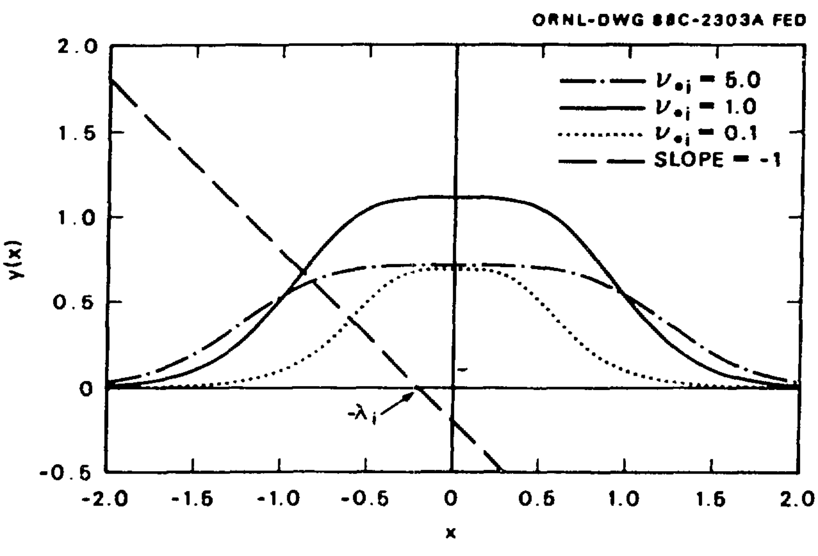
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