Theory of Anomalous Transport in H-Mode Plasmas

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A unified model theory of the anomalous transport of high temperature plasma, L mode and H mode, is developed. Self-sustained turbulence theory for the pressure driven mode is extended, incorporating the effect of inhomogeneous radial electric field $E_r(r)$ to cover two electrically bifurcated states. Strong reductions of heat and current diffusivities, viscosity, the fluctuation amplitude, and correlation length are predicted in nonlinearly sustained turbulence. A consistent radial width of $E_r(r)$ is obtained. A new stability window is also found.

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After the discovery in the ASDEX tokamak [1] of the H mode, which is associated with the sudden reduction of the thermal conductivity χ, the energy flux per particle divided by the temperature gradient, at some critical temperature, the H mode has been observed in almost all tokamaks [2]. The change is most prominent near the edge, and is referred to as a transport barrier. The H-mode characteristic is a generic nature of high temperature toroidal plasmas.

The understanding of transition physics has been advanced. A theory based on the bifurcation of the radial electric field E_r was proposed, involving the interaction of the ion orbit loss and E_r [3]. The sudden change of E_r in a narrow edge region has been confirmed by experiments [4], and an external bias could cause H mode [5]. A rapid change of ion loss at the transition was observed [6]. Elaboration of the model of the transition has followed [7–10]. In H-mode plasmas it is now considered that E_r probably takes a large negative or positive value, the magnitude of which is given by $|X| \sim 1$ ($X \equiv e\rho_p E_r/T$, ρ_p: poloidal ion gyroradius, T: temperature). The effects of E_r on the fluctuations and anomalous transport have been theoretically analyzed and the steep radial gradient of E_r was found to improve linear stability of various modes and to reduce the fluctuation levels [11–16]. Nevertheless the understanding of transport in the H mode is not complete, since it ultimately requires that of the L-mode plasma as well.

Recently, a new development of the theory for the self-sustained turbulence and associated anomalous transport in toroidal hot plasmas was explored [17,18]. It was found that nonlinearly driven viscosities and thermal conductivity play important roles in controlling the fluctuation. The balance between them determines the self-sustained turbulence and associated transport. The theoretical model of the L-mode transport was obtained based on the solution of a nonlinear instability.

In this article, we apply this method to the tokamak plasma, incorporating the effect of the inhomogeneous E_r to make a bridge between L- and H-phase plasmas. Based on this approach, it is now possible to construct a unified transport model of the L-mode phase and the H-mode phase. The anomalous transport coefficients are obtained in a unified and explicit form in terms of the profile parameters. Strong reductions of the thermal conductivity, the electron and ion viscosities, μ_e and μ, and the turbulent level in the H-phase plasma are explained. The change in the fluctuations is also analyzed. Furthermore, an additional stability window due to E_r' is discovered in the higher pressure regime. The anomalous ion viscosity determines Δ, the typical scale length of E_r [19]. Self-consistent solutions of Δ and μ are obtained.

We study a circular tokamak with the toroidal coordinates (r,θ,ζ). The reduced set of equations [20] is employed. The basic equations consist of the equation of motion,

$$n_i m_i \{d(\nabla_\perp^2\phi)/dt - \mu\nabla_\perp^4\phi\} = B^2\nabla_\parallel \cdot \mathbf{J} + B[\nabla p \times \nabla(2r\cos\theta/R)]\cdot\boldsymbol{\zeta},$$

the generalized Ohm's law, $\mathbf{E}+\mathbf{v}\times\mathbf{B}=\mathbf{J}/\sigma - \nabla_\perp^2\lambda\mathbf{J}$, and the energy balance equation, $dp/dt = \chi\nabla_\perp^2 p$. Notations are as follows: m_i is the ion mass, n_i is the ion density, ϕ is the static potential, $\mathbf{B}$ is the main magnetic field, p is the plasma pressure, $\mathbf{J}$ is the current, σ is the classical conductivity, and λ is the current diffusivity ($\lambda \propto \mu_e$). The $\mathbf{E}\times\mathbf{B}$ nonlinear interactions are renormalized in the form of χ, μ, and λ. The detailed derivation was reported in Ref. [21]. The derivative d/dt is $\partial/\partial t + [\phi, \,]/B$ where $[\,,\,]$ denotes the Poisson bracket. The Doppler shift of frequency is offset by the homogeneous $\mathbf{E}\times\mathbf{B}$ rotation. Only the contribution of E_r' to d/dt is retained.

To reduce the partial differential equation to an ordinary differential one, the ballooning transformation [22] is employed as

$$p(r,\theta,\zeta) = \sum_m \exp(-im\theta + in\zeta) \times \int p(\eta)\exp\{im\eta - inq\eta\}\,d\eta$$

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(q is the safety factor) for the microscopic modes. The problem is essentially two dimensional, but the ballooning formulation may be applied by assuming formally the flow shear frequency to be of the order of $1/n$ [23,24]. Eliminating $\tilde{\phi}$ and $\tilde{J}$ from the basic equation, we have the eigenmode equation for p:

$$\frac{d}{d\eta}\frac{F}{\hat{\gamma}+\Xi F+\Lambda F^2}\frac{d}{d\eta}\left(\hat{\gamma}+\omega_{E1}\frac{d}{d\eta}+KF\right)p + \alpha[\kappa+\cos\eta+(s\eta-\alpha\sin\eta)\sin\eta]p-\left(\hat{\gamma}+\omega_{E1}\frac{d}{d\eta}+MF\right)F\left(\hat{\gamma}+\omega_{E1}\frac{d}{d\eta}+KF\right)p=0. \quad (1)$$

We use the normalizations $r/a\rightarrow\hat{r}$, $t/\tau_{Ap}\rightarrow\hat{t}$, $\chi\tau_{Ap}/a^2\rightarrow\hat{\chi}$, $\mu\tau_{Ap}/a^2\rightarrow\hat{\mu}$, $\tau_{Ap}/\mu_0\sigma a^2\rightarrow 1/\hat{\sigma}$, $\lambda\tau_{Ap}/\mu_0 a^4\rightarrow\hat{\lambda}$, $\tau_{Ap}\equiv a\sqrt{\mu_0 m_i n_i}/B_p$, $\gamma\tau_{Ap}\rightarrow\hat{\gamma}$, and notation $\Xi=n^2q^2/\hat{\sigma}$, $\Lambda=\hat{\lambda}n^4q^4$, $K=\hat{\chi}n^2q^2$, $M=\hat{\mu}n^2q^2$, γ is the growth rate, $s=r(dq/dr)/q$, $F=1+(s\eta-\alpha\sin\eta)^2$, $\kappa\equiv-(r/R)(1-1/q^2)$ (average well), $B_p=Br/qR$, $\alpha=q^2\beta'R/a$, $\varepsilon=r/R$, a and R for the major and minor radii, $\beta=2\mu_0 p/B^2$, and $\beta'\equiv d_\beta/d\hat{r}$. The parameter ω_{E1} denotes the shear effect of the radial electric field,

$$\omega_{E1}=\tau_{Ap}(dE_r/d\hat{r})(srB)^{-1}. \quad (2)$$

If we neglect ω_{E1}, Eq. (1) reduces to the transport-driven ballooning mode equation for the L-mode plasma [18]. The ideal magnetohydrodynamic mode equation [22] is recovered by further taking $1/\hat{\sigma}=\hat{\lambda}=\hat{\chi}=\hat{\mu}=0$. In order to apply the ballooning formalism, the subsidiary ordering $\omega_{E1}\sim O(\varepsilon)$ is adopted [24]. It is assumed that $1/\hat{\sigma}=0$, for the plasma of our interest.

Equation (1) predicts that the current-diffusive ballooning mode has a large growth rate. The stability boundary of the nonlinear mode determines the relation between the anomalous transport coefficients $\{\hat{\chi},\hat{\lambda},\hat{\mu}\}$ and the plasma inhomogeneity $\{\beta',E_r'\}$. We study here the case when the ballooning mode is caused by the normal curvature, not by the geodesic curvature, i.e., $\frac{1}{2}+\alpha>s$. For the strongly localized mode, where $s^2\eta^2<1$ and $\eta^2<1$ are satisfied, this eigenvalue equation is approximated by a type of Weber equation, where we neglect the $dp/d\eta$ term as in Refs. [18,21]. The effects of ω_{E1} are incorporated by a perturbation method. These analytic simplifications with $\gamma=0$ reduce Eq. (1) to

$$d^2p/d\eta^2+(\alpha\Lambda/K)\{1-(\tfrac{1}{2}+\alpha-s)\eta^2\}p - M\Lambda[1+3(s-\alpha)^2\eta^2]p+\mathcal{L}p=0, \quad (3)$$

where the operator $\mathcal{L}$ is defined as

$$\mathcal{L}p\equiv\omega_{E1}\left(\frac{K}{F}\frac{d^3p}{d\eta^3}-\Lambda(1+M/K)F^2\frac{dp}{d\eta}\right). \quad (4)$$

The stability boundary is examined by the following perturbation method. Let $\{u_j\}$ be the jth eigenfunction of the original Weber type equation ($\omega_{E1}=0$, $\mathcal{L}p=0$). The most unstable fundamental eigenmode u_0 and the first harmonics u_1 are expressed as $u_0=\xi^{1/4}\pi^{-1/2}\exp\{-\xi\eta^2/2\}$ and $u_1=\sqrt{2/\pi}\xi^{3/4}\eta\exp\{-\xi\eta^2/2\}$, respectively, where $\xi^2=H_0(\frac{1}{2}+\alpha-s)N^2\{1+C(s)N^4\}$, $H_0=\alpha^{3/2}\hat{\lambda}\hat{\chi}^{-3/2}\times\hat{\mu}^{-1/2}$, N is the normalized mode number $N^4=n^4q^4\times\hat{\chi}\hat{\mu}/a$, and $C(s)=3(s-\alpha)^2(\frac{1}{2}+\alpha-s)^{-1}$. The most unstable mode satisfies $N^2=1/f(s)$ where $f(s)$ was given as $(1+2\alpha-2s)\sqrt{2+C(s)}$ [18]. The odd and even parity modes are mixed by the operator $\mathcal{L}$ when $\omega_{E1}\neq 0$. Writing the eigenfunction as

$$p(\eta)=u_0+\rho u_1\cdots, \quad (5)$$

the first order correction to the eigenvalue is given approximately as

$$\frac{H^2}{\xi^2}\simeq 1-\frac{\langle 0|\mathcal{L}|1\rangle\langle 1|\mathcal{L}|0\rangle}{2\xi^2}, \quad (6)$$

where $H\equiv H_0N^2(1-N^4)$. Substituting the eigenfunctions and N, the integrals $\langle 0|\mathcal{L}|1\rangle$ and $\langle 1|\mathcal{L}|0\rangle$ are obtained. Equation (6) is rewritten as

$$H_0=[1+G_1(\alpha,s)\omega_{E1}^2](1+2\alpha-2s)f(s), \quad (7)$$

where

$$G_1=\left[\frac{9}{8}\left(1+\frac{4}{3\alpha f}\right)^2+\frac{25(s-\alpha)^2}{4(1+2\alpha-2s)}\left(\frac{2+C}{1+C}\right)^{1/2}\times\left(1+\frac{4}{3\alpha f}\right)\left(1+\frac{4}{5\alpha f}\right)\right]\frac{\alpha}{f^2}. \quad (8)$$

From the definition of H_0 and Eq. (7), the anomalous transport coefficient with the E_r' contribution is obtained as

$$\hat{\chi}=\frac{f(s)^{-1}\alpha^{3/2}(\hat{\lambda}/\hat{\chi})(\hat{\chi}/\hat{\mu})^{1/2}}{1+G_1\omega_{E1}^2}. \quad (9)$$

The numerator of Eq. (9) is the L-mode transport coefficient χ_L [18]. The effect of E_r' on the thermal conductivity is quantified in Eq. (9), which unifies the L- and H-mode plasma transport. The coefficients are explicitly expressed in terms of the equilibrium quantities, since the self-sustained turbulence is determined. The suppression of the transport is significant when $\omega_{E1}\sim 1/\sqrt{G_1}$. A noticeable reduction of χ occurs when ω_{E1} approaches unity.

Figure 1 shows the results for the transport coefficient. Equation (1) is also solved numerically. The dependence of $\hat{\chi}/\hat{\chi}_L$ on ω_{E1} is illustrated in Fig. 1(a) with other parameters fixed. The dotted line indicates the Lorentzian fit of $\hat{\chi}/\hat{\chi}_L$, giving $\hat{\chi}/\hat{\chi}_L\sim 1/(1+g\omega_{E1}^2)$ with $g\simeq 24$. The result confirms the theoretical result of Eq. (9), with a

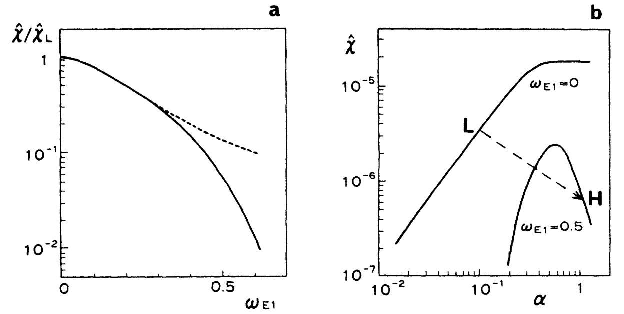

FIG. 1. Dependence of the normalized thermal conductivity $\hat{\chi}$ on ω_{E1} (a) and on α (b). Parameters are $\hat{\mu}/\hat{\chi}=1$, $\hat{\lambda}/\hat{\chi}=10^{-4}$, $\varepsilon=\frac{1}{3}$, $q=3$, and $s=0.5$. In (a) α is fixed as 0.3 and the dotted curve shows the fitting $\hat{\chi}/\hat{\chi}_L=1/(1+g\omega_{E1}^2)$ with $g\simeq 24$. In (b) the arrow indicates typical change in the L- to H-mode transition.

numerical correction of G_1. For the parameters of the calculation, G_1 is predicted to be 12; a factor 2 enhancement in the coefficient G_1 is partly due to the finite value of the magnetic well ($\kappa=-0.29$ in Fig. 1) which is not kept in Eq. (9). When ω_{E1} becomes larger, the suppression becomes more prominent. Figure 1(b) demonstrates the α dependence of χ for $\omega_{E1}=0$ (L mode) and for the case with finite ω_{E1}. The suppression of $\hat{\chi}$ relative to $\hat{\chi}_L$ is stronger in either the small α or large α limits, as is shown in Eq. (8). The reduction of χ at large α in the case of $\omega_{E1}\neq 0$ (which is related to second stability) is important for the establishment of the H mode. When the L mode plasma becomes the H mode, the parameter α would be enhanced from $O(0.1)$ to above $O(1)$. Figure 1(b) shows that the enhancement in α by a factor of 10 allows a reduction of $\hat{\chi}$ by a factor of 10 when the reduction from ω_{E1} is effective. A numerical solution of Eq. (1) confirms that similar results hold for a wide range of the parameter s.

The theory also predicts the change in the fluctuation characteristics. The relation $nq=(\alpha/\hat{\chi}\hat{\mu})^{1/4}N$ shows that the poloidal mode number becomes larger as χ becomes smaller. Combining Eq. (9) and the estimate of $k_\theta(=m/r)$ for the L-mode turbulence [18], we have

$$k_\theta^2=\delta^{-2}\alpha^{-1}(1+G_1\omega_{E1}^2)\,, \tag{10}$$

where δ is the collisionless skin depth. The radial correlation length l_r is also reduced. Its value is given by $l_r^2=\langle k_r^2\rangle^{-1}$ and $\langle k_r^2\rangle=\int n^2q^2s^2\eta^2p(\eta)^2d\eta[\int p(\eta)^2d\eta]^{-1}$, $\langle k_r^2\rangle$ being the average of the squared radial wave number of the fluctuations. Substituting Eq. (5), we have $\langle k_r^2\rangle=[n^2q^2s^2/2\xi](1+2\rho^2)$. Noting the relation $\rho=\langle 1|\mathcal{L}|0\rangle/2\xi$, l_r is given as

$$l_r^2=\frac{2f(s)s^{-2}\xi\hat{\chi}\alpha^{-1/2}}{1+G_r(\alpha,s)\omega_{E1}^2}\,, \tag{11}$$

$$G_r(\alpha,s)=\left(\frac{3}{4\sqrt{2}}+\frac{1}{\alpha f(s)}+\frac{23s^2}{8\sqrt{2}}\right)\frac{\alpha}{f(s)}\,. \tag{12}$$

The fluctuation level is also reduced. The renormalized diffusivity satisfies $\chi\sim(r\tilde{E}/B)^2/\chi n^2q^2$ [21], which gives the estimate $\tilde{\phi}/B\sim\chi$. Using Eq. (9), we have

$$e\tilde{\phi}/T\sim(\chi_L eB/T)(1+G_1\omega_{E1}^2)^{-1}\,. \tag{13}$$

The coefficients G_1 and G_r have a similar magnitude.

The thickness of the transport barrier, Δ, and μ are to be determined self-consistently. In this paper the radial derivative E_r' is simply estimated as E_r/Δ or $XT/e\rho_p\Delta$. It was shown that Δ reduces to the thickness of the source region (ρ_p) in the limit of $\mu\to 0$, and $\Delta\sim\sqrt{\mu/\nu_i}$ in the large μ limit [19,25] (ν_i is the ion collision frequency). We may use the interpolation formula

$$\Delta\sim\sqrt{\rho_p^2+\mu/\nu_i} \tag{14}$$

apart from a numerical factor of order unity. Substituting these relations into ω_{E1} and using the relation $\hat{\mu}\sim\hat{\chi}$ [18,21], we have a convoluted equation for μ in terms of X as

$$\mu=\mu_L\left[1+\frac{\beta}{2}\left(\frac{r}{\rho_p s}\right)^2\frac{\alpha}{1+\mu/\nu_i\rho_p^2}X^2\right]^{-1}\,, \tag{15}$$

where $\mu_L=\chi_L$. This equation shows that μ (and χ as well) is reduced by an increment of $|X|$. Order-of-magnitude reductions of μ, χ, and λ from L-mode values are expected for $|X|\sim 1$.

It is a straightforward extension to study the influence of the curvature of the E_r profile E_r''. The result is given as

$$\chi=\frac{\chi_L}{1+G_1\omega_{E1}^2+G_2\omega_{E2}^2}\,, \tag{16}$$

where $\omega_{E2}=(aE_r''/s^2B)\tau_{Ap}$, $E''=d^2E_r/dr^2$, and $G_2\sim\frac{2}{5}\times(\delta/a)^2$. Details will be given in a forthcoming paper.

In summary, we have developed a unified theory of anomalous transport in the low confinement mode and high confinement mode in toroidal high temperature plasmas. The nonlinear self-sustained turbulence theory for pressure driven modes has been extended to make a bridge between the two bifurcated states. In this theory, the $\mathbf{E}\times\mathbf{B}$ nonlinearity is renormalized in the form of a diffusion operator. To determine the diffusion coefficient, the mean field approximation for the dressed test wave is employed.

The explicit form of χ is obtained [Eq. (16)], and an order-of-magnitude reduction of χ due to the E_r' effect, regardless of the sign of E_r', is predicted. In our framework, not only χ but also $\mu_e(\lambda)$ and μ are reduced by a similar magnitude. The theory predicts that there should be established different transport barriers, namely, for the particle, heat (χ), current (λ), as well as momentum, $E_r(\mu)$ in the H mode. The width of the strong E_r region in the H mode is affected by the reduction of μ due to E_r'. The self-consistent width is estimated [Eq. (15)].

The fluctuation level is also found to be reduced [Eqs. (10), (11), (13)] due to the $|E_r'|$ effect. The mode number of the least stable mode becomes higher for a fixed pressure gradient. Furthermore, another stability window due to E_r' is found, which resembles the second stability window due to low s.

We have adopted the ballooning representation. This would limit the applicable range to $\omega_{E1}/\varepsilon \sim O(1)$. It was shown [24] that the large-flow-shear limit for the classical ballooning formalism converges properly to the small-flow-shear limit of the general theory. The extension to the case of $\omega_{E1}/\varepsilon \gg 1$ is really necessary, since strong reductions of χ, μ, and λ are expected from Eq. (1) when ω_{E1} approaches unity. The recent progress on two-dimensional theory [24,26] would enable this extension. The other simplification is that the Prandtl numbers, $\hat{\mu}/\hat{\chi}$ and $(a/\delta)^2\hat{\lambda}/\hat{\chi}$, are taken to be unity. They can depend on the change of k_θ, which is caused by ω_{E1}. However, this dependence is found to be weak [27].

Connor has succeeded in reproducing our results on the L mode by use of the scale invariance technique [28], supporting the physics basis of our model. The present result (9) is obtained except for the numerical factor. Nonlinear simulation would give this coefficient and allows us to examine the validity of the ansatz in detail. The investigation of effects such as the diamagnetic drift for kinetic corrections, parallel flow or perpendicular compressibility are necessary. These research topics are open for future study.

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