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Beta Limit of Resistive Plasma in Torsatron/Heliotron

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Abstract

Stability against the interchange mode in the Torsatron/Heliotron device is investigated, taking into account the effects of the resistivity, current diffusivity, ion viscosity and thermal diffusivity. Critical beta for the low-mode-number stability is found at the finite beta value. For the range of plasma parameters of the present experiments, the resistive mode and current diffusive mode set comparable critical beta values, which are consistent with experimental observations. In future high temperature plasmas, the current diffusive mode determines the stability limit for the global mode.

keywords: beta limit, interchange instability, magnetic hill, torsatron/heliotron, resistivity, current diffusivity, anomalous transport

1. Introduction

The stability of the interchange mode[1] has been thought to determine the beta limit in stellarators with the finite shear and magnetic hill such as Torsatron/Heliotron devices[2-4] (beta value: ratio of the plasma pressure to the magnetic pressure). Recent progress in computational methods has made it possible to find a configurations which is stable against the pure MHD mode at moderately high beta ($\sim$5%)[5-7]. On the other hand, it has also been well known that the interchange mode can be destabilized by the finite value of resistivity, η, for the parameters in which the pure MHD mode is stable[8]. As a result of this, the instability persists up to the zero-beta value (the growth rate γ scales such that $\gamma \propto \beta^{2/3}\eta^{1/3}$). Finite value of the growth rate continues from the MHD unstable region to the zero beta limit[9-11]. This nature casts uncertainty in defining the beta limit of Torsatron/ Heliotron devices.

One way to establish the beta limit is to search the criterion of beta value for the occurrence of the crash phenomena. The sawtooth oscillation is known to appear when plasma pressure increases[12], and Wakatani et al. has made a model based on the nonlinear development of the resistive interchange mode[13]. It was also found in experiments, however, that the helical deformation appears above a threshold beta value, without developing into the sawtooth crash[14,15]. This observation suggests that there is a stability boundary, which does not necessarily mean the magnetic collapse, at finite beta value in real plasmas.

Although the importance of the resistivity in exciting the mode has been recognized, a little efforts have been paid to study the effects of other transport coefficients such as thermal diffusivity, χ, viscosity, ν, and especially, the current diffusivity, λ. The effects of χ and ν were studied in Ref.[9], but the impact on the beta limit was not discussed in detail. In addition to it, the important role of the current diffusivity has recently been recognized[16,17]. In this article we study the stability of the interchange mode in a cylindrical model by analytical method, taking into account the effects of the resistivity, thermal conductivity, ion viscosity and current diffusivity. The effects of χ and ν on the interchange mode is strong enough so that the critical beta for the low-mode-number stability is found at the finite beta value. For the range of plasma parameters of the present experiments, the resistive mode and current diffusive mode set comparable critical beta values, which are consistent with experimental observations. The beta limit is determined by the transport coefficients. It increases with the increment of χ or ν, and reduces by the enhancement of η and λ. In future high temperature plasmas, the beta limit is determined by the current diffusive interchange mode.

2. Model

We use a model equation based on the reduced set of equations for stellarators[18]. The cylindrical geometry is employed in order to look for the analytical insight. We also consider the case of zero equilibrium current. The model equations consist of the Ohm's law, equation of motion and the energy balance equation, as[11)]

$$\partial A/\partial t = -\nabla_{/\!/} u + \bar{\eta} j - \bar{\lambda}\nabla_{\perp}^{2} j, \tag{1}$$

$$\partial \nabla_{\perp}^{2} u/\partial t = -\nabla_{/\!/} j + (\beta_0/2)\varepsilon^{-2}\Omega'\rho^{-1}\partial p/\partial\theta + \bar{\nu}\nabla_{\perp}^{4} u, \tag{2}$$

$$\partial p/\partial t = -P_{eq}'\rho^{-1}\partial u/\partial\theta + \bar{\chi}\nabla_{\perp}^{2} p, \tag{3}$$

where j is the normalized current $j=\tilde{J}/[B_0/\mu_0 a]$, u is the normalized stream function $u=\tilde{\Phi}/[a^2R/\tau_{Ap}]$, p is the normalized pressure $p=\tilde{p}/[\beta_0\ B_0{}^2/2\mu_0]$ (the symbol tilde indicates the mode of our interest), and the operator $\nabla_{/\!/}$ is defined as

$$\nabla_{/\!/} = (\partial/\partial\zeta - \iota\,\partial/\partial\theta).$$

Length is normalized to the minor radius of the plasma column, a. Conventional notations[11)] are used for the rotational transform, ι, the inverse aspect ratio ε, and the effective curvature of the vacuum magnetic field line Ω. β_0 is the beta value at the axis, R is the major radius, B_0 is the strength of equilibrium magnetic field, and p_{eq} is the shape form of the equilibrium pressure profile, $p_{eq}(0)=1$ and $p_{eq}(1)=0$. Ω' and p_{eq}' are defined as $d\Omega/d\rho$ and $dp_{eq}/d\rho$, respectively. We study the electrostatic limit in the following. The transport coefficient χ, η, ν and λ are normalized as

$$\bar{\chi} = \chi\tau_{Ap}/a^2, \qquad (4\text{-}1)$$

$$\bar{\eta} = \eta\tau_{Ap}/\mu_0 a^2, \qquad (4\text{-}2)$$

$$\bar{\nu} = \nu\tau_{Ap}/a^2, \qquad (4\text{-}3)$$

$$\bar{\lambda} = \lambda\tau_{Ap}/\mu_0 a^4, \qquad (4\text{-}4)$$

and the time is normalized to the poloidal Alfven time τ_{Ap}

$$\tau_{Ap} = R\sqrt{\mu_0 m_i n_i}/B_0, \qquad (5)$$

where n_i is the ion density and m_i is the ion mass.

We solve the eigen value equation by the Fourier transform. We change the variable from $x=\rho-\rho_l$ to k (ρ_l being the mode rational surface) as

$$u = \exp[\gamma t + im\theta - in\zeta] \int u(k)\exp(ikx)dk \qquad (6)$$

where m and n are poloidal and toroidal mode numbers, respectively. Eliminating j and p, we have the eigenvalue equation in the k space as

$$\frac{1}{L^2}\frac{\partial}{\partial k}\frac{1}{\bar{\eta}+\bar{\lambda}k_\perp^2}\frac{\partial}{\partial k}u + \frac{D}{\gamma+\bar{\chi}k_\perp^2}u - (\gamma k_\perp^2 + \bar{\nu}k_\perp^4)u = 0, \qquad (7)$$

where $k_\perp^2 = k_\theta^2 + k^2$, $k_\theta = m/\rho_1$,

$$1/L = k_\theta s, \tag{8-1}$$

$$D = D_0 k_\theta^2 \tag{8-2}$$

$$D_0 = -\beta_0 \ \Omega' p_{eq}'/2\varepsilon^2, \tag{8-3}$$

s denotes the shear parameter $s=\rho_1 \chi'(\rho_1)$, and D_0 denotes the driving by the pressure gradient.

We obtain the eigen value by the WKB method. Introducing new coefficients Λ and σ as

$$\Lambda = \bar{\eta} + \bar{\lambda} k_\theta^2 \tag{9-1}$$

$$\sigma = \bar{\lambda}/\Lambda \tag{9-2}$$

and new variable y as

$$dy/dk = 1 + \sigma k^2, \tag{10}$$

the WKB solution for the most unstable mode satisfy the eigenvalue equation

$$\frac{\pi}{4} = \int_0^{y_c} dy \ Q, \tag{11}$$

and

$$Q = L\sqrt{\frac{\Lambda}{(1+\sigma k^2)}\left\{\frac{D}{(\gamma+\chi k_\perp^2)} - (\gamma k_\perp^2+\nu k_\perp^4)\right\}}, \tag{12}$$

$$Q(y_c) = 0.$$

3. Growth Rate of the Mode

The approximate form of the growth rate of the mode is studied in the large growth rate limit, i.e.,

$$\gamma > \bar{\chi}k_\perp^2,\ \bar{\nu}k_\perp^2. \tag{13}$$

In this limit, Eq.(11) reduces to

$$\frac{\pi}{4} = L\sqrt{\Lambda}\int_0^{y_c} dy\sqrt{\frac{1}{1+\sigma k^2}\left\{\frac{D}{\gamma} - \gamma k_\perp^2\right\}}. \tag{14}$$

The turning point y_c is given by the relation

$$k_\perp = \sqrt{D}/\gamma. \tag{15}$$

Two limiting case is obtained analytically. When the current diffusivity is negligibly small, i.e.,

$$\sigma \ll \tau^2/D, \tag{16}$$

the dispersion relation (14) gives the growth rate of the resistive instability. Neglecting σ in the integrand of Eq.(14), we have the dispersion relation

$$L\sqrt{\Lambda/\tau}[D/\tau - \tau k_\theta{}^2] = 1. \tag{17}$$

Explicit form of τ is given in limiting cases as

$$\tau = D^{2/3}L^{2/3}\Lambda^{1/3} \qquad \text{for } D_0 k_\theta{}^4 \ll s^4 \tag{18}$$

and

$$\tau = \sqrt{D}/k_\theta \qquad \text{for } D_0 k_\theta{}^4 \gg s^4. \tag{19}$$

This recovers the well known results for the resistive interchange mode[3]. The large k_θ limit, (19), is independent of the resistivity and is known as the fast interchange mode[1,9].

Stabilizing effects of the thermal conductivity and viscosity are also estimated by linear perturbation. Taking the first order correction with respect to χ and ν, the dispersion relation for the low k_θ mode can be written as

$$L\sqrt{\Lambda}D\tau^{-3/2}\{1-D\bar{\chi}/2\tau^3-3D\bar{\nu}/8\tau^3\} = 1. \tag{20}$$

The first order correction to τ is given as

$$\gamma = \gamma_0\{1-(\overline{\chi}/3+\overline{\nu}/4)L^{-2}\Lambda^{-1}D^{-1}\}, \tag{21}$$

where γ_0 is given by Eq.(18). Equation (21) implies that both χ and ν play a role to reduce the growth rate of the resistive interchange mode. The similar argument is possible for the fast interchange mode as

$$\gamma = \gamma_0 - \frac{(\overline{\chi}k_\theta{}^2+\overline{\nu}k_\theta{}^2)}{2} \tag{22}$$

where γ_0 is given by Eq.(19).

The current diffusivity can also drive the instability. Taking the other limit,

$$\sigma \gg \gamma^2/D, \tag{23}$$

the dispersion relation is approximated as

$$L\sqrt{\overline{\lambda}}D^{3/2}\gamma^{-5/2} = 3\pi/4. \tag{24}$$

This relation yields the growth rate of the current-diffusivity driven mode as

$$\gamma = (4/3\pi)^{2/5}s^{-2/5}\overline{\lambda}^{1/5}k_\theta{}^{4/5}D_0{}^{3/5}. \tag{25}$$

As is the case of the resistive mode, when mode number is high and the relation $D>\tau^2 k_\theta^2$ is violated, then the growth rate reduces to Eq.(19). The growth rate of the mode which is driven by the current diffusivity is proportional to $\beta^{3/5} k_\theta^{4/5}$ for the small β and/or k_θ limit. Higher-m modes have larger growth rate. The power index to $\bar{\lambda}$ in Eq.(25) is 1/5 and the growth rate weakly depends on $\bar{\lambda}$.

The condition for the transition from the resistive mode to the current-diffusivity driven mode is approximately given as

$$\bar{\lambda}^3 > D_0 k_\theta^{-2} s^{-4} \bar{\eta}^5. \tag{26}$$

For the parameters of D_0, s, $k_\theta \propto O(1)$, the current-diffusivity is the dominant source of the instability for $\bar{\lambda} > \bar{\eta}^{5/3}$.

We can also derive the corrections due to the small thermal conductivity and viscosity. Keeping the first order corrections with respect to $\bar{\chi}$ and $\bar{\nu}$, Eq.(24) is modified as

$$\tau^{5/2} = \frac{4}{3\pi} L\sqrt{\bar{\lambda}} D^{3/2} \{ 1 - \frac{D\bar{\chi}}{\tau^3} - \frac{4\bar{\nu}D}{5\tau^3} \}. \tag{27}$$

The growth rate is written as

$$\tau = \tau_0 - \frac{2D}{5\tau_0^2} (\bar{\chi} + \frac{4}{5}\bar{\nu}), \tag{28}$$

where τ_0 is the growth rate in the absence of the thermal conductivity and viscosity, and is given as Eq.(25).

4. Stability Criterion

The stability criterion is obtained from the dispersion relation Eq.(11) by setting $\tau=0$. We have

$$\frac{\pi}{4} = L\sqrt{\Lambda}\int_0^{y_c} dy \sqrt{\frac{1}{1+\sigma k^2}\left\{\frac{D}{\bar{\chi}k_\perp^2} - \nu k_\perp^4\right\}} \tag{29}$$

for the stability boundary.

The beta limit of the resistive mode is obtained by taking the zero-σ limit of Eq.(29) as

$$\frac{\pi}{4} = L\sqrt{\Lambda}\int_0^{k_c} dk \sqrt{\frac{D}{\bar{\chi}}\left\{\frac{1}{k^2+k_\theta^2} - \frac{\bar{\chi}\bar{\nu}}{D}k_\perp^4\right\}}. \tag{30}$$

The leading term of the integral is the logalythmic term as

$$\frac{\pi}{4} = \frac{L}{6}\sqrt{\frac{\Lambda D}{\bar{\chi}}}\,\ln\left(\frac{64D}{\bar{\chi}\bar{\nu}k_\theta^6}\right). \tag{31}$$

This result is equivalent to the one in Ref.[9]. By successive substitution, the approximate formula of the beta limit is given by the relation as

$$D_0 > D_c = \left(\frac{3\pi s}{2}\right)^2 \frac{\bar{\chi}}{\Lambda} \left(\ln\{\left(\frac{3\pi s}{2}\right)^2 \frac{1}{\bar{\nu}\Lambda k_\theta^4}\} \; 2\ln|\ln\{\frac{\bar{\chi}\bar{\nu}k_\theta^6}{64}\}|\right)^{-2} \tag{32}$$

It is noted that the right hand side is an weakly increasing function of k_θ. Higher m mode is more easily stabilized by the viscosity and thermal conductivity. As the beta value increases, the m=1 mode first appears if the other parameters are unchanged.

In the vicinity of this stability criterion ($\bar{\nu}k_c^2 > |\gamma|$), the dispersion relation, Eq.(31), is approximately given as

$$\frac{\pi}{4} = \frac{L}{6}\sqrt{\frac{\Lambda D}{\bar{\chi}}} \ln\left(\frac{64D\bar{\chi}^2}{\bar{\nu}(\bar{\chi}k_\theta^2+\gamma)^3}\right) . \tag{33}$$

The growth rate is given as

$$\gamma = \left(\frac{64D\bar{\chi}^2}{\nu}\right)^{1/3} \exp\{-\frac{\pi}{2L\sqrt{\Lambda}}\sqrt{\frac{\bar{\chi}}{D}}\} - \bar{\chi}k_\theta^2. \tag{34}$$

The condition γ=0 reduces to Eq.(31). The Taylor expansion around the critical beta value β_c, the growth rate behaves as $\gamma \propto (\beta - \beta_c)$.

In the other limit, i.e., for the current diffusive mode, $\sigma k_c^2 \gg 1$, the dispersion relation with γ=0 gives

$$\frac{\pi}{4} \simeq L\sqrt{\bar{\chi}} \int_0^{k_c} dk \sqrt{\frac{D}{\bar{\chi}} - \bar{\nu}k^6} . \tag{35}$$

This relation gives the critical value, D_c, for D_0 as

$$D_c = C\, s^{3/2} k_\theta^{-1/2} \bar{\lambda}^{-3/4} \bar{\chi} \bar{\nu}^{1/4}, \tag{36}$$

where C is a numerical constant

$$C = \left\{ \frac{4}{\pi} \int_0^1 dk \sqrt{1-k^6} \right\}^{-3/2}$$

and is close to 1.26. The result indicates that the critical beta value for the stability increases with thermal diffusivity and ion viscosity, while it decreases as the current diffusivity is enhanced.

In order to find a beta limit, it is needed to compare Eqs.(32) and (36). First, we apply Eq.(32) to a straight Heliotron E configuration[11] with ε=0.1, ι=0.51+1.69ρ^2, and $d\Omega/d\rho$= 0.1$\rho^2 d(\iota\rho^4)/d\rho$. We here use the numbers $\bar{\chi}\sim 10^{-5}$, $\bar{\nu}=2\times10^{-7}$ and $\bar{\eta}=10^{-6}$, which may be consistent with recent experiments. In this case, D_c in Eq.(32) for the m=1/n=1 mode is evaluated to be about 0.5. When we assume the pressure profile of $p_{eq}(\rho)= \beta_0(1-\rho^2)^2$, this value corresponds to $\beta_{0c}\sim 2\%$, which is in the same order as the one observed experimentally. It is also noted that the order of magnitude estimate of Eq.(32) gives

$$D_c \propto \bar{\chi}/\bar{\eta} \text{ or } \tau_r/\tau_E \tag{37}$$

where τ_r and τ_E are the resistive diffusion time and the energy confinement time. The right hand side strongly increases with the electron temperature. In future high temperature plasmas, the global resistive interchange mode is more stable.

The relation between the quantities $\bar{\chi}$, $\bar{\lambda}$ and $\bar{\nu}$ depends on the nature of the anomalous transport. For instance, for the case where the anomalous transport is driven by the magnetic braiding (which is often expected in the helical devices), we have[17)]

$$\bar{\nu}/\bar{\chi} = \sqrt{m_e/m_i} \tag{38}$$

and

$$\bar{\lambda}/\bar{\chi} = (\delta_s/\rho_i a)^2 \tag{39}$$

where m_e is the electron mass, and δ_s is the collisionless skin depth. The ratio $\bar{\lambda}/\bar{\chi}$ is about 10^4 for the present experimental devices. Using these estimations, we have

$$D_c \simeq 2\times10^2 s^{3/2}(\tau_{Ap}/\tau_E k_\theta)^{1/2} \tag{40}$$

Taking the example of $\tau_E/\tau_{Ap} \sim 3\times10^4$, which is a typical value for present experiments, we have the order of magnitude estimate for the stability beta limit

$$D_c \simeq s^{3/2}/\sqrt{k_\theta}. \tag{41}$$

Although the numerical coefficient of the order of unity may not be correct in the right hand side of Eq.(41), the beta limit for the global mode derived from Eq.(41) is within the range of experimental observations. Therefore, the beta limit determined by the current diffusive mode is comparable with the one by the resistive interchange mode for the present experimental plasma parameters.

5. Summary and Discussions

In this article, we investigated analytically the stability of the interchange mode in Torsatron/Heliotron devices in the presence of the thermal diffusivity, ion viscosity, current diffusivity and resistivity. The stability boundary was found at finite beta value. This result gives the basis for determining the beta limit in helical devices with magnetic hill and shear.

The stability boundary is investigated for the resistive mode and current-diffusivity driven mode. We found that, for the parameters of the present experiments, the resistive mode and current diffusive mode give similar critical beta value for the stability. The obtained critical beta value seems within the range of experimental observations. For future high temperature plasmas, the current diffusive mode will play more important role in determining the beta limit. The precise determination employing the computational code will be reported in a separate article.

The result indicates the important influence of the transport coefficient on the beta limit. The beta limit should be studied by taking into account the proper model of the plasma transport coefficients. It has been pointed out recently that the enhanced transport coefficient can cause the catastrophic phenomena in the dynamics of the global modes[17]. The application of such analysis on the finite beta plasma in helical devices deserves further study.

Higher m mode is more easily stabilized for the resistive interchange mode. Such a high m-modes are often studied in connection with the anomalous transport in helical devices[9,20,21]. The present analysis indicates that the current diffusive mode, the threshold beta value of which decreases as the mode number increases, would be more plausible candidate for the anomalous transport. This study is left for future research.

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