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Modification of the Tearing Mode Growth Rate by the Presence of a Stochastic Magnetic Field

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MODIFICATION OF THE TEARING MODE GROWTH RATE BY THE
PRESENCE OF A STOCHASTIC MAGNETIC FIELD

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ABSTRACT

We have studied the effect of a stochastic magnetic field on the growth of a tearing mode. An analytic solution is given for the case in which the stochastic magnetic field is static. In that case, an unstable tearing mode is further destabilized (stabilized) if the Δ' values of the nonlinearly driven modes are positive (negative). This mechanism explains the destabilization of the ($m = 3$; $n = 2$) tearing mode observed in nonlinear three-dimensional calculations.

Three-dimensional numerical calculations^{1,2} have shown that the nonlinear interaction of tearing modes with different helicities can lead to a very rapid growth of the modes involved. This destabilization is triggered by the overlap of magnetic islands, i.e., by the stochasticization of magnetic field lines. The mechanism by which the destabilization takes place is investigated in this report.

We study here the effect on the linear growth of a test wave tearing mode in the resistive magnetohydrodynamic (MHD) approximation of a static stochastic magnetic field. Treating the stochastic field as static is justified by the fact that the growth rates of the nonlinearly driven waves we will discuss are very rapid when compared with the slow evolution of the stochastic background. For the same reason, we neglect the slow fluid motions associated with the stochastic background.

The equilibrium, adequately modeled by slab geometry, has a magnetic field given by: $\vec{B}_{eq}(x) = B_0 b(x) \hat{y} + B_0 \hat{z} + \vec{B}_{st}(x, y, z)$, where $\vec{B}_{st}$ is the static stochastic magnetic field (which is small compared with B_0 and perpendicular to the z direction). The total magnetic field can be expressed in terms of the poloidal flux function ψ , $\vec{B} = \vec{\nabla}\psi \times \hat{z} + B_0 \hat{z}$. The function ψ is given in terms of Fourier components by

$$\psi = \psi_{eq}(x) + \sum_{k'} \psi_{k'}^{st}(x) e^{i(k'_1 y + k'_2 z)} + \sum_k \tilde{\psi}_k(x, t) e^{i(k_1 y + k_2 z)} + c.c. ,$$

where $d\psi_{eq}/dx = b(x)$, $\psi_{k'}^{st}(x)$ are the components of the stochastic magnetic field poloidal flux function, and $\tilde{\psi}_k(x, t)$ are the corresponding components of the linear test wave. The velocity is given in terms of the stream function ϕ , $\vec{v}_1 = \vec{\nabla}\phi \times \hat{z}$, and it is contributed to only by the perturbation:

$$\phi = \sum_k \tilde{\phi}_k(x, t) e^{i(k_1 y + k_2 z)} + c.c.$$

In addition, there are the driven modes, $\tilde{\psi}_{k-k'}^{(2)}$, and $\tilde{\phi}_{k-k'}^{(2)}$. We work to first order in $\tilde{\psi}_k$ and $\tilde{\phi}_k$ and to second order in the stochastic fields.

The functions Φ and Ψ are the solutions of the reduced set of resistive MHD equations:^{1,3}

$$\frac{\partial \Psi}{\partial t} + \frac{\partial \Phi}{\partial y} \frac{\partial \Psi}{\partial x} - \frac{\partial \Phi}{\partial x} \frac{\partial \Psi}{\partial y} = S^{-1} \eta J - E_z - \frac{\partial \Phi}{\partial z}, \quad (1)$$

$$\frac{\partial \nabla_{\perp}^2 \Phi}{\partial t} + \frac{\partial \Phi}{\partial y} \frac{\partial \nabla_{\perp}^2 \Phi}{\partial x} - \frac{\partial \Phi}{\partial x} \frac{\partial \nabla_{\perp}^2 \Phi}{\partial y} = - \frac{\partial \Psi}{\partial x} \frac{\partial J}{\partial y} + \frac{\partial \Psi}{\partial y} \frac{\partial J}{\partial x} - \frac{\partial J}{\partial z}. \quad (2)$$

Here, $J = \nabla_{\perp}^2 \psi$ is the current density and η is the resistivity, normalized to a typical value η_0 . Both equations are in dimensionless form with the poloidal flux function normalized to $B_0 a$ and the stream function normalized to a^2/τ_h ; a , the width of the slab, and $\tau_h = a\sqrt{\rho_m}/B_0$ are the length and time units, respectively. The parameter $S = \tau_r/\tau_h$ is the magnetic Reynolds number, where $\tau_r = a^2/\eta_0$ is the skin time.

From Eqs. (1) and (2), we can calculate the time evolution of the perturbation. Keeping terms up to third order, we find the equations for $\tilde{\psi}_k$ and $\tilde{\phi}_k$ are:

$$\begin{aligned} \frac{\partial \tilde{\psi}_k}{\partial t} = & -ik_1 \frac{x - x_k}{L_s} \tilde{\phi}_k - i \frac{\partial}{\partial x} \left[\sum_{k'} (k_1 - k_1') \psi_{k'}^{st} \tilde{\phi}_{k-k'}^{(2)} \right] \\ & + ik_1 \sum_{k'} \psi_{k'}^{st} \frac{\partial \tilde{\phi}_{k-k'}^{(2)}}{\partial x} + S^{-1} \eta \tilde{J}_k \end{aligned} \quad (3)$$

and

$$\begin{aligned} \frac{\partial}{\partial t} \left(\frac{\partial^2 \phi_k}{\partial x^2} - k_1^2 \phi_k \right) = & -ik_1 \frac{x - x_k}{L_s} \tilde{J}_k + ik_1 \frac{dJ_{eq}}{dx} \tilde{\psi}_k \\ & + i \frac{\partial}{\partial x} \sum_{k'} k_1' \left[\psi_{k'}^{st} \tilde{J}_{k-k'}^{(2)} - J_{k'}^{st} \tilde{\psi}_{k-k'}^{(2)} \right] \\ & + ik_1 \sum_{k'} \left[\frac{\partial J_{k'}^{st}}{\partial x} \tilde{\psi}_{k-k'}^{(2)} - \frac{\partial \psi_{k'}^{st}}{\partial x} \tilde{J}_{k-k'}^{(2)} \right]. \end{aligned} \quad (4)$$

Now, x_k is defined by $\vec{k} \cdot \vec{B}_{eq}(x_k) = 0$. In these equations we have retained only the nonlinear terms that generate a coherent contribution to the mode k .⁴ The J_k is the k component of the toroidal current density. We have assumed that the shear is constant near the singular surface and L_s is the shear length $(q'/q^2)^{-1}$.

The evolution of the driven modes is similarly given by

$$\frac{\partial \psi_{k-k'}^{(2)}}{\partial t} = -i \frac{(k_1 - k'_1)(x - x_{k-k'})}{L_s} \tilde{\phi}_{k-k'}^{(2)} - i \left(k'_1 \psi_{-k'}^{st} \frac{\partial \tilde{\phi}_k}{\partial x} + k_1 \frac{\partial \psi_{-k'}^{st}}{\partial x} \tilde{\phi}_k \right), \quad (5)$$

$$\begin{aligned} \frac{\partial}{\partial t} \left[\frac{\partial \tilde{\psi}_{k-k'}^{(2)}}{\partial x^2} - (k_1 - k'_1)^2 \tilde{\phi}_{k-k'}^{(2)} \right] \\ = -i \frac{k_1 - k'_1}{L_s} (x - x_{k-k'}) \tilde{J}_{k-k'}^{(2)} + i(k_1 - k'_1) \frac{dJ_{eq}}{dx} \tilde{\psi}_{k-k'}^{(2)} \\ + i \left(k'_1 J_{-k'}^{st} \frac{\partial \psi_k}{\partial x} - k'_1 \psi_{-k'}^{st} \frac{\partial \tilde{J}_k}{\partial x} - k_1 \frac{\partial \psi_{-k'}^{st}}{\partial x} \tilde{J}_k + k_1 \frac{\partial J_{-k'}^{st}}{\partial x} \tilde{\psi}_k \right), \quad (6) \end{aligned}$$

where we keep only terms which give the coherent contribution to the mode k in Eqs. (3) and (4). For simplicity, we have neglected the resistivity here. Because the stochastic field is static, it is implied that the time dependence can be factorized in the form $e^{\gamma t}$. Then, by assuming $\gamma \tau_h \ll 1$ and $L_s/a \sim 1$, Eqs. (5) and (6) can be solved, using the fact that the homogeneous part of the equation describes the well-known outer region in tearing mode theory. We find:

$$\tilde{\phi}_{k-k'}^{(2)} = \frac{iL_s}{(k_1 - k'_1)(x - x_{k-k'})} \left[S_1(x_{k-k'}) g_{k-k'}(x) - S_1(x) \right],$$

$$\tilde{\psi}_{k-k'}^{(2)} = \frac{S_1(x_{k-k'})}{\gamma} g_{k-k'}(x),$$

where

$$\left[\frac{d^2}{dx^2} - (k_1 - k'_1)^2 - \frac{L_s}{x - x_{k-k'}} \frac{dJ_{eq}}{dx} \right] g_{k-k'}(x) = \Delta'_{k-k'} \delta(x - x_{k-k'}),$$

$$g_{k-k'}(x_{k-k'}) = 1,$$

and

$$S_1 \equiv -i \left(k'_1 \psi_{-k}^{st}, \frac{\partial \tilde{\phi}_k}{\partial x} + k_1 \frac{\partial \psi_{-k}^{st}}{\partial x} \tilde{\phi}_k \right).$$

This solution is substituted into Eqs. (3) and (4), and after neglecting nonresonant contributions, we obtain the renormalized equations for the mode k :

$$\gamma \tilde{\psi}_k = -i \frac{k_1(x - x_k)}{L_s} \tilde{\phi}_k + iL_s \sum_{k'} \frac{\Delta'_{k-k'}}{k_1 - k'_1} \delta(x - x_{k-k'}) \left| k'_1 \psi_{k'}^{st} \right|^2 \frac{d\tilde{\phi}_k}{dx} + nS^{-1} \tilde{J}_k, \quad (7)$$

$$\gamma \left(\frac{d^2 \tilde{\phi}_k}{dx^2} - k_1^2 \tilde{\phi}_k \right) = -i \frac{k_1(x - x_k)}{L_s} \tilde{J}_k + ik_1 \frac{dJ_{eq}}{dx} \tilde{\psi}_k + \frac{1}{\gamma} \frac{d}{dx} \left\{ \left[\sum_{k'} \Delta'_{k-k'} \delta(x - x_{k-k'}) \left| k'_1 \psi_{k'}^{st} \right|^2 \right] \frac{d\tilde{\phi}_k}{dx} \right\}. \quad (8)$$

We note that summations over k' may be done, and for reasonable stochastic fields, they will give weakly varying functions of x , which we may consider to be constant. We then have:

$$\gamma \tilde{\psi}_k = -i \frac{k_1(x - x_k)}{L_s} \tilde{\phi}_k - iC_1 \frac{d\tilde{\phi}_k}{dx} + \eta S^{-1} \frac{d^2 \tilde{\psi}_k}{dx^2}, \quad (9)$$

$$(\gamma^2 - C_2) \frac{d^2 \tilde{\phi}_k}{dx^2} = -ik_1 \frac{x - x_k}{L_s} \gamma \frac{d^2 \tilde{\psi}_k}{dx^2}. \quad (10)$$

These may be solved in terms of parabolic cylinder functions; matching this solution to the known solution outside the layer,⁵ we obtain the dispersion relation:

$$\Delta'_k = \pi \left(\frac{L_s}{k_1} \right)^{1/2} (\gamma^2 - C_2)^{1/4} \gamma^{3/4} \frac{S^{3/4}}{2^\alpha} \frac{\Gamma\left(\frac{3}{4} - \frac{\alpha}{2}\right)}{\Gamma\left(2 + \frac{1}{4} - \frac{\alpha}{2}\right)} \\ \times \frac{(1 - 4\alpha^2)^{3/4}}{(1 + 2\alpha)^{3/2}} F_2^1 \left(\frac{\alpha}{2} - \frac{1}{4}, \quad \frac{3}{2}; \quad 2 - \frac{1}{4} + \frac{\alpha}{2}; \quad \frac{2\alpha - 1}{2\alpha + 1} \right), \quad (11)$$

where

$$C_1 = L_s \sum_{k'} \frac{\Delta'_{k-k'} \delta(x - x_{k-k'}) |k'_1 \psi_{k'}^{st}|^2}{k_1 - k'_1},$$

$$C_2 = \sum_{k'} \Delta'_{k-k'} \delta(x - x_{k-k'}) |k'_1 \psi_{k'}^{st}|^2,$$

and

$$4\alpha^2 = \frac{C_1^2 \gamma S}{4(\gamma^2 - C_2) + C_1^2 \gamma S}.$$

For $C_1 = 0$, a particular solution is $\gamma^2 \equiv C_2$ and $\tilde{\psi}_k = 0$, as may be seen by inspection of Eqs. (9) and (10). For $k_1 C_1 > 0$, Eq. (11) also yields $\gamma^2 \equiv C_2$, if $C_2 \gg (\Delta'_k/3)^{8/5} (L_s/k_1)^{4/5} S^{-6/5}$ and $C_2^2 \gg (\Delta'_k)^2 k_1 |C_1|^{3/2} \pi^3 L_s$. The growth rate is thus given by

$$\gamma^2 = \sum_{k'} \left| \frac{B_{rk'}^{st}}{B_0} \right|^2 \frac{\Delta'_{k-k'}}{L_s / |k_1 - k'_1|} \tau_h^{-2} \approx \left| \frac{\vec{B}_{st}}{\vec{B}} \right|^2 \frac{\Delta'}{W} \tau_h^{-2}, \quad (12)$$

where $W \sim L_s / |k_1 - k'_1|^{-1}$ is the radial mode extent. Note that the growth is purely on the MHD time scale; i.e., it is faster than usual tearing modes. The mode has the structure of a highly localized vortex at $x = x_k$. One thus expects a rapid conversion of magnetic to kinetic energy by this nonlinear mechanism.

Three-dimensional calculations of the evolution of tearing modes of different helicity^{1,2} have shown that modes are in fact nonlinearly destabilized. In the particular case of tokamak disruptions, equilibria that are unstable to the $(m = 2; n = 1)$ and $(m = 3; n = 2)$ modes show nonlinear evolution, which leads to overlap of the magnetic islands. At this point, the $(m = 3; n = 2)$ mode is strongly destabilized. The overlap of islands produces field line stochastization, and the main component of this stochastic field is the $(m = 2; n = 1)$ mode. The evolution of the 2/1 magnetic island is in the Rutherford⁶ regime (linear with time) and is practically unaffected by the interaction with the $(m = 3; n = 2)$ mode. Therefore, we can assume that the stochastic magnetic field is static and is essentially given by the $(m = 2; n = 1)$ mode. From this point, the effect of this stochastic field on the growth of the $(m = 3; n = 2)$ mode can be deduced from the previous calculation. The driven mode is essentially the $(m = 5; n = 3)$ mode, whose Δ' can be determined from the numerical calculations and is found to be positive when the $(m = 3; n = 2)$ mode is nonlinearly destabilized. As a result, the overlap of the 2/1 and 3/2 magnetic islands has a double effect. On the one hand, it generates the stochastic magnetic field; on the other hand, due to the quasilinear deformation of the

toroidal current profile near the $q = 5/3$ singular surface, the Δ' of the $(m = 5; n = 3)$ mode becomes positive. From the previously derived expression for γ , Eq. (12), the growth rate for the $(m = 3; n = 2)$ mode can be calculated, and it agrees well with the results of the numerical computation (see Fig. 1). The comparison is done for the kinetic energy growth rate of the $(m = 3; n = 2)$ mode, which starts increasing before the magnetic energy does. This is to be expected from the solutions given for Eqs. (7) and (8), which have already been mentioned. Further evidence of the role of the $(m = 5; n = 3)$ mode in the destabilization of the $(m = 3; n = 2)$ mode can be obtained by redoing the numerical calculation without including the $5/3$ helicity in it. The nonlinear growth rate of the $(m = 3; n = 2)$ mode is larger than in the previous case up to the time at which $\Delta'_{5/3}$ becomes positive (see the point designated by the arrow in Fig. 1); at that time it becomes smaller and no longer has a sharp peak. The numerical results clearly support the existence of the destabilizing mechanism suggested by this work.

For a more complete understanding of the effect of a stochastic magnetic field on the growth of a tearing mode, it will be necessary to consider the nonstatic case. The assumption that the stochastic field is no longer static seriously complicates the calculations, and it is not a trivial extension of the work presented here. The nonstatic case is important in a later phase of the numerical calculations. In this case, due to the presence of fluid turbulence, terms that can be interpreted as anomalous resistivity and fluid viscosity become important. There is some evidence of these effects from the numerical results,⁷ and we are extending our calculations to this more general case. Needless to say, comparison with actual experiments, rather than with numerical computations, may also require consideration of the effects of stochastic fields at the microscopic level (e.g., electron viscosity, modified thermal transport, etc.).

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Fig. 1. The $(m = 3; n = 2)$ kinetic energy growth rate obtained in the numerical calculations of Ref. 2 (—) is compared with the value obtained from Eq. (12) (---). Also given is the rate obtained in the numerical calculation when the 5/3 helicity was not included (— — —).

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