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Resistive Ballooning Mode Equation

G. Batemen
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ABSTRACT

A second-order ordinary differential equation on each flux surface is derived for the high mode number limit of resistive MHD ballooning modes in tokamaks with arbitrary cross section, aspect ratio, and shear. The equation is structurally similar to that used to study ideal MHD ballooning modes computationally. The model used in this paper indicates that all tokamak plasmas are unstable, with growth rate proportional to resistivity when the pressure gradient is less than the critical value needed for ideal MHD stability.

A new kind of large-scale instability, called a ballooning mode, is predicted to appear in tokamaks when the ratio of plasma pressure to magnetic pressure (β) is raised to several percent.¹⁻⁶ Theoretical predictions indicate that the maximum stable β can be increased by proper choice of the plasma shape, proximity of the wall, and profiles of the pressure and current. Estimates of finite Larmor radius and kinetic effects⁷⁻¹⁰ have not appreciably changed the predicted onset of this instability.

Ballooning modes are analogous to the Rayleigh-Taylor instability of a plane slab plasma in a gravitational field and a horizontal magnetic field intercepting fixed endplates.¹¹ The mode is unstable when the restoring force due to field line bending is overwhelmed by a sufficiently strong gravitational force. Resistivity, however, allows a slow instability to persist in an arbitrarily small gravitational field, since the plasma can flow through the magnetic field with less field line bending. Instead of a clear-cut onset of instability at some critical value of the driving force, there is always a resistive instability with growth rate increasing gradually as the driving force is increased and with a transition to the ideal growth rate at a sufficiently large driving force. We show in this Letter that toroidal ballooning modes behave in much the same way in the presence of resistivity.

The analytic theory of toroidal ballooning modes is significantly simplified in the limit of large toroidal mode number ($e^{-in\zeta}$, $n \rightarrow \infty$). To lowest order in this limit, the ideal MHD theory reduces to finding the eigenvalue and eigenfunction associated with a second-order ordinary differential equation for the displacement along the cross-sectional

contour of each flux surface.^{4,5} To next order in this large n expansion, it is found that the instability is concentrated around the flux surface with the largest eigenvalue and that its growth rate is well approximated by and bounded from above by the largest eigenvalue. We show that the resistive MHD problem simplifies to a similar ordinary differential equation in the large toroidal mode number limit.

Consider the linearized resistive MHD equations for the perturbed velocity $\underline{v}$, magnetic field $\underline{B}^1$, and pressure p^1 ,

$$\rho \partial \underline{v} / \partial t = -\nabla p^1 + [(\underline{\nabla} \times \underline{B}) \times \underline{B}]^1, \quad (1)$$

$$\partial \underline{B}^1 / \partial t = \underline{\nabla} \times (\underline{v} \times \underline{B} - \eta \underline{\nabla} \times \underline{B})^1, \quad (2)$$

$$\partial p^1 / \partial t = -\underline{v} \cdot \underline{\nabla} p, \quad (3)$$

which describe small motions about a static, scalar pressure, axisymmetric equilibrium with prescribed pressure $p(\psi)$, toroidal field $RB_\zeta = F(\psi)$, mass density ρ , and resistivity η . Here $\psi(R, y)$ is the poloidal magnetic field stream function which is uniform on each flux surface, (R, y, ζ) are cylindrical coordinates around the centerline of the torus, and the equilibrium satisfies the equations

$$\underline{\nabla} p = (\underline{\nabla} \times \underline{B}) \times \underline{B}, \quad (4)$$

$$\underline{B} = \underline{\nabla} \psi \times \underline{\nabla} \zeta + F(\psi) \underline{\nabla} \zeta. \quad (5)$$

Note that the isothermal approximation is used in Eq. (3) so that the perturbed pressure is determined by convection alone and the sound speed is neglected. This approximation properly retains the behavior of

ideal MHD ($\eta = 0$) at marginal stability. Previous work¹ shows this approximation to be reasonable for determining the growth rate of resistive ballooning modes also.

The derivation proceeds by using an orthogonal flux coordinate system (ψ, χ, ζ) with Jacobian J . Perturbed quantities are expressed in the form

$$A = \tilde{A}(\psi, \chi) \exp(\gamma t + i n \phi) , \quad (6)$$

$$\phi(\psi, \chi, \zeta) \equiv -\zeta + \int_0^\chi d\chi' J B_{\zeta'} / R , \quad (7)$$

where the phase ϕ is uniform along equilibrium magnetic field lines ($\underline{B} \cdot \nabla \phi = 0$). To leading order in $n \gg 1$, Eq. (2) reduces to

$$\gamma G \underline{B}^1 = \nabla \times (\underline{v} \times \underline{B}) , \quad (8)$$

where

$$G \equiv 1 + \eta n^2 |\nabla \phi|^2 / \gamma . \quad (9)$$

Equation (8) indicates that resistivity reduces the perturbed magnetic field for a given plasma velocity. The equation makes sense only if $\eta n^2 / \gamma \ll n$; otherwise $B^1 \sim 0$ and the plasma behaves as if it were unmagnetized. This requirement on γ will be satisfied by the resistive ballooning mode derived here.

Although Eq. (8) is very similar to the ideal MHD Faraday's law (recovered in the limit $\eta \rightarrow 0$, $\gamma \neq 0$), it is difficult to construct a useful variational form from Eqs. (1), (3), and (8) since the equations

are no longer Hermitian when $G \neq 1$. Hence, instead of a δW treatment, we proceed with a direct solution of the equations.

The ψ component of the equation of motion, Eq. (1), is

$$\rho \gamma X = R^2 B_{\chi}^2 [-in\phi' P^1 + \underline{B} \cdot \underline{\nabla} (B_{\psi}^1 / R B_{\chi}) - 2 B_{\chi} B_{\chi}^1 (J B_{\chi})' / J B_{\chi} - 2 B_{\zeta} B_{\zeta}^1 R' / R] , \quad (10)$$

where $X \equiv \underline{v} \cdot \underline{\nabla} \psi = v_{\psi} R B_{\chi}$, $P^1 \equiv p^1 + \underline{B} \cdot \underline{B}^1$ is the total perturbed pressure, and the prime denotes $(\partial/\partial\psi)_{\chi, \zeta}$. From this and other components of the equation of motion it can be seen that $P^1 \sim O(1/n)$ in the $n \gg 1$ expansion. Then, using Eq. (3), this implies

$$\gamma \underline{B} \cdot \underline{B}^1 = p' X \quad (11)$$

to leading order. When Eq. (8) is used to express $\underline{B}^1$ in terms of $\underline{v}$, one finds that $\underline{\nabla} \cdot \underline{v} \sim O(1)$ and to leading order

$$\phi' R B_{\chi} X = B_{\chi} v_{\zeta} - B_{\zeta} v_{\chi} , \quad (12)$$

$$\phi' R^2 B_{\chi}^2 B_{\psi}^1 = B_{\chi} B_{\zeta}^1 - B_{\zeta} B_{\chi}^1 . \quad (13)$$

These equations, together with the $\underline{\nabla} \psi \times \underline{B}$ component of the equation of motion, can be used to eliminate $in P^1$ and $\underline{B}^1$ from Eq. (10) and reduce it to an equation in X alone. The Jacobian J can be eliminated by using the identity

$$-R^2 (J B_{\chi}^2)' / J = R^2 p' + F F' , \quad (14)$$

which is derived by expressing the equilibrium equations, Eqs. (4) and (5), in (ψ, χ, ζ) coordinates.

The resulting second-order ordinary differential equation for the large toroidal mode limit of resistive ballooning modes in an axisymmetric toroidal plasma is

$$\mathbf{B} \cdot \nabla (|\nabla \phi|^2 / B^2 G) \mathbf{B} \cdot \nabla X + (D_R - \rho \gamma^2 |\nabla \phi|^2 / B^2) X = 0, \quad (15)$$

where the driving term D_R has the form

$$\begin{aligned} D_R &\equiv \frac{dp}{d\psi} \left[2 \frac{\partial}{\partial \psi} \left(p + \frac{1}{2} B^2 \right) - \frac{\partial \phi}{\partial \psi} \frac{RB}{B^2} \zeta \mathbf{B} \cdot \nabla B^2 \right] / B^2 \\ &= 2p' \mathbf{B} \times \nabla \phi \cdot \nabla / B^2 \end{aligned} \quad (16)$$

and κ is the magnetic field line curvature,

$$\kappa \equiv \hat{\mathbf{B}} \cdot \nabla \hat{\mathbf{B}} = \left[\mathbf{B} \times \nabla \left(p + \frac{1}{2} B^2 \right) \right] \times \mathbf{B} / B^4. \quad (17)$$

Expressions for ϕ and G are given by Eqs. (7) and (9), and

$$|\nabla \phi|^2 = (RB_X \phi')^2 + (B/RR_X)^2. \quad (18)$$

The operator $\mathbf{B} \cdot \nabla$ in Eq. (15) can be written as $B_X \partial / \partial \ell$, where ℓ is the poloidal arc length around a flux contour. All equilibrium quantities are periodic in ℓ but ϕ' grows secularly with ℓ . To achieve periodicity of the mode in a sheared magnetic field we use the Fourier transform technique⁵ or the infinite series technique.⁷ In either case, Eq. (15) should be solved on the infinite interval $-\infty < \ell < \infty$ for the square-integrable eigenfunction X and its eigenvalue γ .

While Eq. (15) is amenable to computational solution for arbitrarily prescribed equilibria or to analytic approximation in some cases, it is instructive at this point to examine some general features of the solution using dimensional analysis and variational forms.

Note that resistivity appears only in the field line bending term (the first term) in Eq. (15) and only through the factor $G = 1 + \eta n^2 |\nabla \phi|^2 / \gamma$. When the driving term D_R is sufficiently large, the resulting growth rate is so large that $\gamma \gg \eta n^2 |\nabla \phi|^2$, $G \approx 1$, and essentially the ideal MHD ballooning mode is recovered with

$$\rho \gamma^2 \approx \text{Max} \frac{\int d\ell [X^2 D_R / B_X - (\partial X / \partial \ell)^2 B_X |\nabla \phi|^2 / B^2]}{\int d\ell X^2 |\nabla \phi|^2 / B^2 B_X} \quad (19)$$

The driving term D_R is generally positive around the outer edge of the toroid and negative around the inner edge, giving X its typically ballooning character. Below a critical pressure gradient, depending upon geometry and profiles, the driving term is so small that Eq. (19) predicts ideal MHD stability.

However, as the pressure gradient is decreased to this critical value, the growth rate diminishes to the point where $\gamma \ll \eta n^2 |\nabla \phi|^2$. Then the large factor $G \approx \eta n^2 |\nabla \phi|^2 / \gamma \gg 1$ substantially reduces the field line bending term in Eq. (15), and the variational form becomes

$$\rho \gamma^2 \int d\ell X^2 |\nabla \phi|^2 / B^2 B_X + (\gamma / \eta n^2) \int d\ell (\partial X / \partial \ell)^2 B_X / B^2 - \int d\ell X^2 D_R / B_X = 0. \quad (20)$$

Note that the plasma is always unstable according to this model — even

with arbitrarily small driving term. At low β , the growth rate is proportional to $\eta n^2 D_R L^2$, where L is the poloidal connection length,

$$\gamma \approx \eta n^2 \int d\ell (X^2 D_R / B_\chi) / \int d\ell [(\partial X / \partial \ell)^2 B_\chi / B^2] . \quad (21)$$

Using $\partial X / \partial \ell \sim X / L$, $D_R \sim \beta_p / R^3 a_p$, and $q \sim CB / 2\pi R B_\chi$ [where $\beta_p \equiv 2p / B_\chi^2$ is the poloidal beta, $a_p \equiv p / |\nabla p|$ is the pressure gradient scale length, and $C \equiv \oint d\ell$ is the poloidal circumference (note that usually $L \sim C$)], then the resistive growth rate is approximated by

$$\gamma \sim (\eta / a^2) n^2 q^2 (\beta_p a / R) (L^2 / C^2) (a / a_p) , \quad (22)$$

where a is the minor radius. It can be seen that this growth rate is generally much faster than resistive magnetic diffusion, $\gamma_R \sim \eta / a^2$, especially near the ideal MHD ballooning marginal point $\beta_p \sim R / a$. This agrees with similar estimates with compression included in the model.^{1,6}

The transition between ideal and resistive ballooning behavior is somewhat complicated by the fact that with shear, $|\nabla \phi|$ is secular as well as periodic in ℓ ,

$$\phi' = 2\pi q' \ell / C + \tilde{\phi}' , \quad (23)$$

where $\tilde{\phi}'$ is a periodic function of ℓ and $q(\psi) = F \oint d\ell / 2\pi B_\chi R^2$. If $\eta = 0$ exactly, an unstable eigenfunction behaves at infinity like

$$X \sim \exp(-\sqrt{\rho} \gamma \int^\ell d\ell / B_\chi) . \quad (24)$$

However, for $\eta \neq 0$ we have at large ℓ

$$X \sim \exp(-nq' \sqrt{\rho \eta \gamma} \int^{\ell} d\ell \ell 2\pi R/C) . \quad (25)$$

An arbitrarily small amount of resistivity in a system with shear always leads to an $\exp(-\ell^2)$ dependence¹⁰ rather than the intrinsically slower $\exp(-\ell)$ dependence because resistivity allows a relatively larger $\partial X/\partial \ell$ with less field line bending as $\ell \rightarrow \infty$.

In summary, Eq. (15) is derived for the $n \rightarrow \infty$ limit of ballooning modes in an arbitrary axisymmetric toroidal plasma using a resistive MHD model with isothermal equation of state. It is shown that resistive ballooning modes are always unstable. Below the critical β for ideal ballooning modes, the growth rate is proportional to the resistivity but is generally faster than resistive diffusion.

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