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MASTER

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I. High β MHD stability studies with PEST

J. Manickam

Abstract The PEST code has been modified to use higher order cubic finite elements. This has resulted in a more rapid convergence of the eigenvalue with increasing numbers of functions, extending the applicability of the code to high- β , high- q regimes which were previously difficult to compute.

PEST is a variational ideal MHD stability code^{1,2} which determines the linear stability of an equilibrium to infinitesimal perturbations. The stability analysis is best accomplished in a straight field line coordinate system. This is readily obtained when the equilibrium field is expressed in the form,

$$B = \nabla\zeta \times \nabla\Psi + q(\Psi)\nabla\Psi \times \nabla\theta,$$

where q is the safety factor. Here Ψ is the radial coordinate and θ and ζ are generalized poloidal and toroidal angles respectively. The Jacobian $\mathcal{J}$ is defined by

$$\mathcal{J}^{-1} \equiv \nabla\Psi \times \nabla\theta \cdot \nabla\zeta.$$

The perturbations are represented by a finite subset of a complete set of functions, $\xi = \sum a_m \Phi_m$. The variational problem can then be reduced to the determination of the eigenvalues and eigenvectors of the generalized matrix eigenvalue problem

$$\sum_m [\omega^2 \langle \Phi_m^* | K | \Phi_m \rangle - \langle \Phi_m^* | \delta W | \Phi_m \rangle] a_m = 0.$$

We find it convenient to use a mixed finite-element and fourier harmonics representation for Φ_n

$$\Phi(\mathbf{r}) = \sum_l \sum_m \sum_n \Phi_{l,m}(\Psi) e^{i(l\theta - n\phi)},$$

where $\Phi(\Psi)$ represents a finite element in the Ψ direction and θ and ϕ represent poloidal and toroidal angles respectively. In an axisymmetric geometry the different toroidal, n , modes decouple and we can drop this subscript. This leaves a double sum over l , the poloidal mode number, and m the radial basis functions. An accurate representation of the eigenvalues and eigenvectors generally requires a very large number of expansion functions. Since this number is limited by computer memory and time, it is often necessary to extrapolate the results from the use of successive sets of functions.

This task is simplified if the extrapolation formulas are known. The expected convergence behavior^{2,3} is:

$$\omega^2 = \omega_\infty^2 + C_L e^{-\alpha \Delta L}$$

Figure 1: Linear finite elements used in PEST. This example shows 10 elements identified by superposed labels.

in the poloidal direction, and

$$\omega^2 = \omega_\infty^2 + C_M/M^2$$

in the radial direction. Here ω_∞^2 is the converged result and ω^2 the value obtained with the use of ΔL fourier harmonics or M finite elements.

In the radial direction the choice of basis functions is finite elements. The simplest elements consistent with the requirements of the function space are linear elements, see Fig. 1. In a typical stability analysis run approximately one hundred finite elements are used. The choice of ΔL is determined by n and the range of variation in the q profile. Typically we need $l_{\min} \leq l \leq l_{\max}$ with $l_{\min} \sim -5$ and $l_{\max} \sim nq_{\max} + 10$. These requirements are dependent on the choice of the Jacobian which defines the straight field-line coordinate system. The two most commonly used Jacobians are,

	Coordinate
$\mathcal{J} \propto x^2$	; PEST
$\mathcal{J} \propto \frac{x}{ \nabla\psi }$	; Eq. Arc length

The choice of the Jacobian is particularly useful in limiting the number of poloidal harmonics required to represent the instability. This is readily inferred from an examination of the distribution of constant Θ curves for PEST and equal arc-length coordinates.

Figure 2: Lines of constant Ψ and θ in A) PEST and B) equal arc length coordinates.

Clearly θ_{PEST} has a sparser distribution on the low field side compared to $\theta_{\text{eq.arcs}}$ and would require more fourier harmonics to represent the displacement vector. Consequently, the equal arc length coordinate system has been the standard choice for most of the previous studies. These previous studies were mainly limited to relatively low- β and low- q equilibria and the convergence properties followed the established behavior.

Recently the interest in second stability which requires $q_{\text{axis}} \sim 2$ and $\epsilon\beta_{\text{pol}} > 1$ has required extending these studies into these new regimes.

In the high $\epsilon\beta_{\text{pol}} > 1$, high q regimes we encounter two new difficulties.

- the high q implies that the range ΔL must be increased substantially to ensure an accurate representation of the high l components.
- high β_{pol} implies a separatrix very close to the plasma edge and a rapidly changing radial contribution.

The first issue makes it difficult to use the PEST coordinates which have a poor spatial representation on the low-field side, though they carefully tailor the Θ grid in the vicinity of the separatrix. The second makes it difficult to use equal arc length coordinates. As a consequence we find that the PEST co-ordinates treat the Ψ dependence adequately but have poor Θ convergence. In contrast the equal arc length co-ordinates have the opposite virtue and disadvantage. This is illustrated in Fig. 3, where we show the convergence properties of a low- q high- β equilibrium, $R/a = 3$, $\kappa = 1.75$, $\delta = 0.5$, $\beta_{\text{pol}} = 1.37$, $g = 4$.

We note that in equal arc length coordinates the eigenvalue saturates when approximately 15 fourier harmonics are included, Fig. 3a. In contrast we need more than 35 harmonics in PEST coordinates. On the other hand the PEST coordinates converge more rapidly in the radial direction as seen in Fig. 3b. It is also clear that the convergence in equal arc length coordinates is not in the asymptotic regime where the $1/M^2$ dependence is satisfied.

In order to remedy this we have introduced higher order finite elements in the Ψ direction to improve the radial convergence for the equal-arc length co-ordinates. The new representation is based on Hermite cubic polynomials. Both even and odd functions are included and each element spans five grid points. A schematic of these functions is shown in Fig. 4. With this representation we find significantly improved convergence in equal arc length coordinates, comparable to that obtained in PEST coordinates. The results using cubic elements is shown in Fig. 5, which should be compared with Fig. 3b. The convergence in PEST coordinates with cubic rather than linear elements is relatively unchanged, unlike the results in equal arc length coordinates.

The improved convergence with cubic elements is most dramatically observed when examining the stability of high- q high- β equilibria. We show the results for an ARIES-2 equilibrium, $R/a = 4$, $\kappa = 1.8$, $\delta = 0.6$, $\beta_{pol} = 4.65$, $g = 6.4$, in Fig. 6. It is clear that with linear elements the code has great difficulty converging to the correct answer as the region of asymptotic convergence is not reached even when we use 1200 finite elements! In contrast with just 200 cubic elements we are within a few per cent of the converged solution.

In conclusion we report that the PEST stability code has been significantly improved by the introduction of higher order finite elements, cubic as opposed to linear, and is now capable of correctly determining the stability properties of high- β , high- q equilibria.

Figure 3: Convergence of ω^2 with increasing numbers of A) poloidal fourier harmonics, and B) radial finite elements. In both figures the results in equal arc length coordinates is labelled "1" and in PEST as "2".

Figure 4: Hermite cubic functions used as finite elements. Each element is labelled and both even, Φ_E , and odd, Φ_O , functions are shown.

Figure 5: Convergence in M , the number of radial cubic finite elements in 1) equal arc length and 2) PEST coordinates.

Figure 6: Convergence in M in equal arc length coordinates for a high- β high- q ARIES-2 equilibrium. The values of M are given at the top of the figure.

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II. High β : Is it the Promised Land?

Steve Cowley

Abstract The transport properties of ultra high β tokamaks are discussed. Qualitative considerations suggest that transport may improve with β .

The advantages of high β reactors ($\beta \sim \mathcal{O}(1)$) are not quite clear. However, if the energy flux to the wall or divertor is a real limiting factor (and it probably is), it is clear that high β is desirable only if the energy confinement increases significantly with β . Given our lack of understanding of transport at low β , predictions of transport at high β are at best speculative. In this note, we sketch the properties of ultra high β tokamaks that suggest that transport may indeed improve at high β . We will make the standard assumption that transport is caused by instabilities and further assume that stabilizing instabilities will reduce transport.

First, let us discuss the equilibrium of high β large aspect ratio tokamaks. (The mathematical solution of this problem is given in Ref. 1.) Consider the forces on the plasma: the pressure force, $-\nabla p$; the toroidal field force, $-\nabla B_T^2/2 + \mathbf{B}_T \cdot \nabla \mathbf{B}_T$; and the poloidal field force, $-\nabla B_p^2 + \mathbf{B}_p \cdot \nabla \mathbf{B}_p$. Nominally, $\mathbf{B}_p \sim \mathcal{O}(\epsilon \mathbf{B}_T/q)$ (where $\epsilon = a/R_0$, the inverse aspect ratio) and $p \sim \mathcal{O}(B_T^2)$. Thus the poloidal field forces are small, $\mathcal{O}(\epsilon^2)$, compared to the pressure and toroidal field forces. Since $B_T = F(\psi)/R$ and $p = p(\psi)$, the balance between toroidal field forces and pressure ($\nabla p = -(1/R^2)F\nabla F$) yields $\psi = \psi(R)$. This is not magnetic confinement since the flux surfaces are not closed nested toroids. Closure of the flux surfaces is achieved by squashing the poloidal flux into a boundary layer (see figure) thereby enhancing the poloidal field. Typically, flux surfaces take up the shape shown in the figure. Most of the volume is in the *core region* where $\psi = \psi(R)$ and the poloidal field is negligible. The flux surfaces are closed in the *boundary layer region*. Note that $\nabla p + 1/R^2 F\nabla F \sim \epsilon \nabla p$ in the boundary layer. Thus a simple estimate of the poloidal field in the boundary layer, B_{pbl} , is obtained by balancing $-\nabla B_p^2/2$ with the force $\epsilon \nabla p$ to yield $B_{pbl} \sim \sqrt{\epsilon p}$. Since the poloidal flux in the core is equal to the poloidal flux in the boundary layer, the width of the boundary layer, δ , can be approximated by $\delta/a \sim B_{pcore}/B_{pbl} \sim \sqrt{\epsilon/(q^2\beta)}$. Our concern here is with $\beta \gg \epsilon/q^2$ so that $\delta \ll a$. The magnetic field strength varies **across** the surface by a factor of order β and **in** the surface by a factor of order ϵ . (Note that $|\mathbf{B}|$ is constant along the surface in the core.)

Now let us consider five properties of these equilibria that are beneficial for stability and transport.

(i) Reduced Volume in the Boundary Layer.

Since the magnetic field does not vary on the flux surface in the core, there are no

trapped particles in the core. Trapped particles are a fraction $\epsilon^{1/2}$ of the particles in the boundary layer and the boundary layer now occupies a fraction $\sqrt{\epsilon/(q^2\beta)}$ of the total volume. Thus, the trapped particles are a fraction $\epsilon^{1/2}\sqrt{\epsilon/(q^2\beta)}$ of all the particles.² A direct consequence of this is that high β neoclassical transport coefficients are reduced from their low β values by the factor $\sqrt{\epsilon/(q^2\beta)}$. The reduction in the trapped particle fraction may also be beneficial for anomalous transport since it reduces the drive for trapped particle modes.

(ii) Strong Average Good Curvature.

The average curvature is dominated by the core (because the core has more volume) and the curvature in the core is “good.” It is relatively straightforward to show that this ensures interchange or “Mercier” stability.³ At these high values of β , the curvature effects on stability are important even for kink modes. Tearing modes are strongly stabilized by average good curvature.⁴ One expects instabilities to be localized in the boundary layer so as to avoid the good curvature region. This possible reduction in the volume where turbulence occurs is obviously good. However, the gradients in the boundary layer are large and therefore turbulence in the boundary layer may be more effective in causing transport.

(iii) Short Connection Length in the Boundary Layer.

The connection length, ℓ_c , (the length of field line in the boundary layer) is $a/\sqrt{\epsilon\beta}$ at high β ($\ell_c \sim qR$ at low β). Thus, the $k_{\parallel}$ of a mode localized to the boundary layer is less than $\sqrt{\epsilon\beta}/a$. Part of the stabilization of ballooning modes is due to this shortening of the connection length and part is due to the strong shear (see the next point). The reduction of ℓ_c also makes ion Landau damping more effective at stabilizing drift wave-type instabilities.

(iv) Strong Local Shear in the Boundary Layer.

The local shear of the field lines in the boundary layer is large.³ A simple measure of local shear (which is related to the usual definition³) is obtained by considering an infinitesimal vector $\delta\mathbf{r}$ separating two field lines. We choose $\delta\mathbf{r}$ to point radially (i.e. parallel to $\nabla\psi$) at the outermost point of the flux surface and $\delta\mathbf{r} \cdot \mathbf{B} \equiv 0$. As one transports $\delta\mathbf{r}$ along $\mathbf{B}$ the angle between $\delta\mathbf{r}$ and $\nabla\psi$, λ , changes. If we go a distance, ℓ_c , along the field line we find $\tan \lambda \sim \mathcal{O}((\beta q^2/\epsilon)^{1/2})$, i.e. $\delta\mathbf{r}$ is roughly perpendicular to $\nabla\psi$. Obviously, we would have reached an identical conclusion if we had chosen $\delta\mathbf{r}$ to point radially at some other poloidal location in the boundary layer. Short wavelength ($k_{\perp}a \gg 1$) instabilities with $k_{\parallel} \ll k_{\perp}$ have perpendicular displacements parallel to a $\delta\mathbf{r}$. Thus, the strong shear tends to make the displacements perpendicular to $\nabla\psi$.³ However, such displacements do not extract energy from the temperature or density gradients. Clearly, the strong shear helps stabilize all modes that are driven by temperature or density gradients.

(v) Minimum B.

When $\beta \gg \epsilon$, the variation of $|\mathbf{B}|$ across the surface far exceeds the variation in the surface. $|\mathbf{B}|$ takes its minimum value at the magnetic axis. This has an important effect on trapped particle drifts. The ∇B drift clearly dominates the guiding center drifts. When $\beta \gg \epsilon/q^{2/3}$ the guiding center drift and the diamagnetic drift are in opposite directions (so called drift reversal). Thus, for modes with $\omega \sim \omega_*$ there is no resonant interaction between the mode and the trapped particles. The precession frequency, ω_D , is given by the formula $\omega_D = (\partial J / \partial \psi)_\mu / (\partial J / \partial E)_\mu$, where $J = \oint v_{\parallel} dl$ is the "Second Adiabatic Invariant." One may easily verify that, for $\beta \gg \epsilon/q^{2/3}$, outward displacements of trapped particles in which J and μ are conserved, raise the particle energy. Clearly, for a normal density profile, i.e. a centrally peaked profile, trapped particles stabilize positive energy modes for which μ and J are conserved. It is possible that negative energy modes are destabilized by the trapped particles.⁵ (This is the subject of current research.)

Is high β the promised land? These five properties do make high β tokamaks look attractive. They are, however, largely qualitative indications; not predictions. One should also note that transport, while being a critical issue, is not the only issue of importance to high β tokamaks. Of course, the question will ultimately be answered by experiments. Much recent effort has been focused on obtaining "high $\epsilon\beta_p$," i.e. $\beta > \epsilon/q^2$, often with $q \gg 1$ and $\beta \leq \epsilon/q^{2/3}$. Unfortunately, such discharges do not benefit from minimum B (point (v)) and therefore may not be exploiting all the benefits of high β .

Figure: Typical Equilibrium Flux Surfaces for a Circular Boundary.

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III. Theory Publication List 1/92 - 3/92

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Biglari, H. and M. Ono, P. Diamond, and G.G. Craddock, *Flow Shear Suppression of Turbulence Using Externally Driven Ion Bernstein and Alfvén Waves*, PPPL Report 2811.

Hahm, T.S., *Flow Shear Induced Compton Scattering of Electron Drift Instability*, Phys. Fluids B and PPPL Report 2830.

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Valeo, E.J. and S.C. Cowley, *A Device to Produce Lasing Action at Wavelength of 400Å and Below*, Phys. Rev. A and PPPL Report 2835.

Jardin, S.C., M. Bell, and N. Pomphrey, *TSC Simulation of Ohmic Discharges in TFTR*, Nucl. Fusion and PPPL Report 2839.

III. Theory Visitor List 1/92 - 3/92

Bhattacharjee, Amitava, University of Columbia, New York, February.

Dewar, Robert, Australian National University, Australia, January.

Dolginov, Arkady, Bartol Research Institute, University of Delaware, on leave from Ioffe Physical-Technical Institute, Russia, January.

Hegna, Chris, University of Wisconsin, Madison, February.

Heikkila, Walter, Physics Department of University of Texas, Dallas, February.

Rayevski, Dimitri, New York University, Courant Institute, January, February.

Roy, Steve, Caltech, Ca, February.

Schneider, Ralf, IPP, Garching, Germany, March - April.

Shaing, Kerchung, ORNL, Oak Ridge, TN, February.

Tanaka, Motohiko, National Institute for Fusion Science, Nagoya, Japan, March.

Tessarotto, Massimo, University of Trieste, Italy, January.

Wootton, Alan, Fusion Research Center, University of Texas, Austin, February.

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