Axisymmetric Stability of Vertically Asymmetric
Tokamaks at Large Beta Poloidal

K. Yamazaki,* H. Fishman, and M. Okabayashi
Plasma Physics Laboratory, Princeton University,
Princeton, New Jersey 08544

A. M. M. Todd,
Grumman Aerospace Corporation
105 College Road East,
Princeton, New Jersey 08540

DISCLAIMER

* On leave from: Institute of Plasma Physics, Nagoya University, Nagoya 464, Japan.

ABSTRACT

The stability of high- β vertically asymmetric tokamak equilibria to rigid displacements is investigated analytically. It is found that vertical stability at large beta poloidal is mainly determined by a coupling between the shape of the plasma surface and the Shafranov shift of the magnetic axis. To the lowest order, symmetric components of the plasma surface shape are found to be the critical destabilizing elements. Asymmetric components have little effect. The inclusion of higher order terms in the high β tokamak expansion leads to further destabilization. Qualitative agreement between these analytic results and numerical stability calculations using the PEST code is demonstrated.

1. INTRODUCTION

The beneficial effects of ellipticity and triangularity on the non-axisymmetric stability of tokamaks have long been recognized [1]. Unfortunately, vertical elongation induces an axisymmetric vertical instability which requires active or passive feedback for stabilization. Consequently, a firm understanding of shaping trade-offs with regard to axisymmetric and non-axisymmetric stability is critical for tokamak reactor designs.

The role of ellipticity in tokamak stability was first studied analytically by Rutherford [2]. The effect of rectangular shaping on axisymmetric stability was discussed by Rosen [3]. However, these two works did not take toroidal effects into account. We will show that with increasing beta, the toroidal effects begin to dominate the shaping effects. Haas [4] and Wakatani [5] have performed analytical calculations including toroidal effects for elliptical plasmas only.

Recently, a single-divertor configuration with up-down asymmetry has been considered for the ETF and INTOR designs [6] in order to simplify certain aspects of the engineering design. The vertically asymmetric configuration has been investigated experimentally in PDX [7] and Doublet III [8]. There have been some numerical calculations of axisymmetric stability in these asymmetric configurations, but all at low beta [9, 10].

In this paper, we extend Haas' elliptical analysis [4] to arbitrarily-shaped high- β equilibria, and discuss the dependence of poloidally symmetric and asymmetric shaping on vertical stability. Furthermore, we discuss the effect on stability of a coupling between toroidicity and shaping.

Section 2 describes the double-expansion analytic equilibrium where a shaping parameter and the conventional high- β toroidal expansion parameter are used. Section 3 presents analytic stability calculations for a rigid vertical displacement utilizing the ideal MHD energy principle. Finally, in Section 4, these analytic results are compared with numerical results generated by the PEST code [11].

2. ANALYTICAL EQUILIBRIUM MODEL

In order to treat arbitrarily shaped tokamaks, we use a non-orthogonal non-circular quasi-toroidal coordinate system (ρ, ϕ, θ) , as well as a circular quasi-toroidal system (r, ϕ, θ) and a cylindrical system (R, ϕ, z) . Here

$$\rho = r / (1 + \delta(\theta)) , \quad (1)$$

and the shaping factor $\delta(\theta)$ is defined by

$$\begin{aligned} \delta(\theta) &= \sum_{\ell=1}^{\infty} (\delta_{\ell}^C \cos \ell\theta + \delta_{\ell}^S \sin \ell\theta) , \\ &= \sum_{\ell=1}^{\infty} \delta_{\ell} \cos (\ell\theta + \alpha_{\ell}) . \end{aligned} \quad (2)$$

Figure 1 shows the relationships between these coordinate systems. The Fourier coefficients δ_{ℓ} and phases α_{ℓ} are determined by the shape of the plasma boundary, $(\rho = a)$. We assume that departure $\delta \sim \epsilon = a/R_0$, the inverse aspect ratio, and expand in ϵ .

Axisymmetric toroidal equilibria obey the Grad-Shafranov equation,

$$L\Psi \equiv \left(R \frac{\partial}{\partial R} \frac{1}{R} \frac{\partial}{\partial R} + \frac{\partial^2}{\partial z^2} \right) \Psi ,$$

$$= - \mu_0 R j_\phi , \quad (3)$$

$$= - \frac{1}{2} \frac{d}{d\Psi} I^2 - \mu_0 R^2 \frac{d}{d\Psi} P .$$

Here $p = p(\Psi)$ is the plasma pressure, $I = I(\Psi)$ is the toroidal field function, and the equilibrium magnetic field is defined in terms of the poloidal magnetic flux Ψ as

$$\vec{B} = \frac{I}{R} \vec{e}_\phi + \frac{\nabla \Psi \times \vec{e}_\phi}{R} .$$

We restrict our analysis to equilibria characterized by the following pressure and toroidal field functions:

$$P = P_0 \left(1 - \Psi/\Psi_s \right) ,$$

$$I^2 = I_0^2 + \left(I_s^2 - I_0^2 \right) \Psi/\Psi_s = \left(R B_\phi \right)^2 , \quad (4)$$

with Ψ normalized to zero at the plasma center ($\rho = 0$), and having the value Ψ_s at plasma surface ($\rho = a$).

We now adopt the conventional high- β ordering for expansion in ε ,

$$\beta_p \equiv \frac{2\mu_o^2 R_o P_o}{2\mu_o R_o^2 P_o - (I_s^2 - I_o^2)} \sim O(\epsilon^{-1}), \quad (5)$$

and utilize the restriction $\delta \sim \epsilon$ to also expand in δ .

Thus Eq. (3) can be expressed up to second order as:

$$a^2(1 + \delta)^2 L\Psi = (\Delta^1 + \Delta^\epsilon + \Delta^\delta + \Delta^{\epsilon\delta} + \Delta^{\delta\delta} + \dots) \cdot (\Psi^1 + \Psi^\epsilon \Psi^\delta + \Psi^{\epsilon\epsilon} + \Psi^{\epsilon\delta} + \Psi^{\delta\delta} + \dots) \quad (6)$$

$$= -4\Psi_o^\epsilon \left[1 + 2\epsilon\beta_p x \cos\theta + \epsilon^2\beta_p^2 x^2 \cos^2\theta + 2\delta(1 + 3\epsilon\beta_p x \cos\theta) + \dots \right], \quad (7)$$

with the superscripts ϵ and δ denoting the relevant ordering. Here,

$$\begin{aligned} \Delta^1 &= \frac{1}{x} \frac{\partial}{\partial x} x \frac{\partial}{\partial x} + \frac{1}{x^2} \frac{\partial^2}{\partial \theta^2}, \\ \Delta^\epsilon &= -\epsilon \left(\cos\theta \frac{\partial}{\partial x} - \frac{\sin\theta}{x} \frac{\partial}{\partial \theta} \right), \\ \Delta^\delta &= -\left(\frac{2\delta}{x} \frac{\partial}{\partial \theta} + \frac{\delta}{x} \right) \frac{\partial}{\partial x}, \end{aligned}$$

and

$$x \equiv \rho/a.$$

Order by order, we find

$$\Psi^1 = \Psi^0, \quad (8)$$

$$\Psi^E = \Psi_0^E (1 - x^2) (1 + \epsilon \beta_P x \cos \theta), \quad (9)$$

$$\Psi^\delta = 0, \quad (10)$$

$$\begin{aligned} \Psi^{EE} = \Psi_0^E \epsilon \left\{ -\frac{\epsilon \beta_P}{4} x^2 (1 - x^2) + \frac{1}{4} x (1 - x^2) \cos \theta \right. \\ \left. + \frac{\epsilon \beta_P}{4} x^2 (1 - x^2) \cos 2\theta \right\}, \end{aligned} \quad (11)$$

$$\begin{aligned} \Psi^{E\delta} = \Psi_0^E \sum_l \delta_l \{ \cos(l\theta + \alpha_l) \cdot [2(x^l - x^2) \\ + \epsilon \beta_P (x - 3x^3 + x^{l-1} + x^{l+1}) \cos \theta] \\ + \sin(l\theta + \alpha_l) \cdot [\epsilon \beta_P (x^{l-1} - x^{l+1}) \sin \theta] \}, \end{aligned} \quad (12)$$

and

$$\Psi^{\delta\delta} = 0. \quad (13)$$

These results reduce to those of Ref. [12] in the circular limit $\delta \rightarrow 0$.

3. VERTICAL STABILITY ANALYSIS USING THE ENERGY PRINCIPLE

In this section, we describe the vertical stability of the analytic equilibria described in the previous section. Assuming a rigid Z-displacement

$$\vec{\xi} = \xi_Z \vec{e}_Z,$$

the energy integral for the plasma region simplifies to

$$\delta W_p = \frac{\xi_Z^2}{2\mu_0} \int dV \left\{ \left(\frac{\partial B_Z}{\partial Z} \right)^2 + \frac{\partial B_Z}{\partial R} \frac{\partial B_R}{\partial Z} \right\}, \quad (14)$$

with the volume element

$$dV = 2\pi R_0 a^2 \{1 + \epsilon x(1 + \delta(\theta)) \cos \theta\} (1 + \delta(\theta))^2 x dx d\theta.$$

Since our equilibrium model has no skin current, the remaining contribution to the potential energy is that of the vacuum region surrounding the plasma,

$$\delta W_v = \frac{1}{2\mu_0} \int \frac{[\nabla \tilde{v}_v]^2}{R^2} dV. \quad (15)$$

Here $\tilde{v}_v$ is the perturbed poloidal vacuum flux function corresponding to a perturbed magnetic field,

$$\vec{b} = \frac{\nabla \tilde{v}_v \times \vec{e}_\phi}{R},$$

which satisfies

$$L \tilde{V}_v = 0. \quad (16)$$

Inside the plasma, we write the perturbed flux

$$\tilde{V}_p = -\epsilon_z \frac{\partial}{\partial z} \Psi,$$

and use the connection condition

$$\tilde{V}_p(\theta) = \tilde{V}_v(\theta) \quad \text{at } \rho = a,$$

to determine $\tilde{V}_v$. The problem is closed by requiring $\tilde{V}_v = 0$, on the surface of a perfectly conducting shell at $\rho = \rho_w$, which exactly matches the equilibrium plasma shape. Expanding $\delta W = \delta W_p + \delta W_v$ in terms of ϵ and δ as in Section 2, vertical stability is determined by

$$\delta W > 0. \quad (17)$$

To the lowest order, we obtain the cylindrical limit and the circular toroidal limit,

$$\delta W_p^{\epsilon\epsilon} = -\frac{4\pi^2 \epsilon_z^2}{\mu_0 a^2 R_0^2} (\Psi_0^\epsilon)^2 \left(1 + \frac{1}{2} (\epsilon \beta_p)^2\right), \quad (18)$$

$$\delta W_v^{\epsilon\epsilon} = \frac{4\pi^2 \epsilon_z^2 (\Psi_0^\epsilon)^2}{\mu_0 R_0^2 a^2} \left(\frac{x_w^2 + 1}{x_w^2 - 1} + \frac{1}{2} (\epsilon \beta_p)^2 \frac{x_w^4 + 1}{x_w^4 - 1} \right), \quad (19)$$

with

$$\begin{aligned} \tilde{v}_v^\varepsilon = \frac{\varepsilon_Z}{a} \psi_0^\varepsilon \left[\frac{2x(1 - x_w^2/x^2)}{\sin\theta} \right. \\ \left. + \frac{\varepsilon\beta_p x^2(1 - x_w^4/x^4)}{1 - x_w^4} \sin 2\theta \right], \end{aligned} \quad (20)$$

and

$$\tilde{v}_v^\delta = 0. \quad (21)$$

The resultant stability criterion is:

$$\left(\frac{\Gamma}{\Gamma_{PA}}\right)^2 = -4 \frac{x_w^2 + 1}{x_w^2 - 1} + 4 - 2(\varepsilon\beta_p)^2 \left(\frac{x_w^4 + 1}{x_w^4 - 1} - 1\right) < 0. \quad (22)$$

Where the poloidal Alfvén growth rate Γ_{PA} is defined by

$$\Gamma_{PA}^2 = (v_{PA}/a)^2 = (\psi_0^\varepsilon)^2 / \mu_0 M R_0^2 a^4. \quad (23)$$

v_{PA} is the poloidal Alfvén velocity, and M is the mass density. The growth rate Γ was estimated from

$$\Gamma^2 = -2\delta W/\xi^2 \int M dV.$$

In Eq. (22), the first and third terms correspond to shell and toroidal stabilization respectively, while the second term is the instability driving force. In the absence of a conducting wall ($\rho_w = \infty$), the plasma is marginally

stable, so that a wall always provides stability to the lowest order. As $x_w \rightarrow \varepsilon^{-1/2}$ we must include the next-order terms in δW to check the stability. In this order, the plasma terms are

$$\delta W_p^{\varepsilon\varepsilon\delta} = - \frac{4\pi^2 \xi_z^2}{\mu_0 R_0 a^2} (\Psi_0^\varepsilon)^2 \left\{ 2 (\varepsilon\beta_p)^2 \delta_2 \cos\alpha_2 - \frac{15}{8} \varepsilon\beta_p \delta_3 \cos\alpha_3 \right\}, \quad (24)$$

and

$$\delta W_p^{\varepsilon\varepsilon\varepsilon} = - \frac{4\pi^2 \xi_z^2}{\mu_0 R_0 a^2} (\varepsilon\beta_p)^2 \frac{3}{8} \varepsilon. \quad (25)$$

The vacuum fluxes

$$\begin{aligned} \tilde{V}_v^{\varepsilon\delta} = & \frac{\xi_z \Psi_0^\varepsilon}{a} \sum_l \delta_l \left\{ \cos(l\theta + \alpha_l) \left[\left(-\frac{2}{x} - \frac{2(l-1)}{x^{l-1}} + \frac{2}{x^{l+1}} \right) \sin\theta \right. \right. \\ & \left. \left. + \varepsilon\beta_p \left(-\frac{2}{x^2} - \frac{l-3}{x^{l-2}} + \frac{2}{x^{l+2}} \right) \sin 2\theta \right] \right. \\ & \left. + \sin(l\theta + \alpha_l) \left[\varepsilon\beta_p \frac{l+1}{x^l} + \left(\frac{2}{x} \frac{(l-1)}{l-1} + \frac{2}{x^{l+1}} \right) \cos\theta \right. \right. \\ & \left. \left. + \varepsilon\beta_p \left(\frac{l-3}{x^{l-2}} + \frac{2}{x^{l+2}} \right) \cos 2\theta \right] \right\}, \quad (26) \end{aligned}$$

and

$$\begin{aligned} \tilde{V}_V^{\epsilon\epsilon} = \frac{\epsilon_z^2 \Psi_0^2}{a} \epsilon \left\{ -\frac{\epsilon\beta_p}{2} \frac{1}{x} \sin\theta + \left(\frac{1}{2} - \frac{5}{4} \frac{1}{x^2} \right) \sin 2\theta \right. \\ \left. + \epsilon\beta_p \left(\frac{1}{4} \frac{1}{x} - \frac{1}{2} \frac{1}{x^3} \right) \sin 3\theta \right\}, \end{aligned} \quad (27)$$

lead to the following vacuum terms,

$$\begin{aligned} \delta W_V^{\epsilon\epsilon\delta} = \frac{4\pi^2 \epsilon_z^2}{\mu_0 R_0^2 a} (\Psi_0^2)^2 \left\{ (2 + 3 (\epsilon\beta_p)^2) \delta_2 \cos\alpha_2 + \frac{37}{6} \epsilon\beta_p \delta_3 \cos\alpha_3 \right. \\ \left. + (\epsilon\beta_p)^2 \delta_4 \cos\alpha_4 \right\}, \end{aligned} \quad (28)$$

and

$$\delta W_V^{\epsilon\epsilon\epsilon} = \frac{4\pi^2 \epsilon_z^2}{\mu_0 R_0^2 a} (\Psi_0^2)^2 \frac{\epsilon}{8} (\epsilon\beta_p)^2. \quad (29)$$

Thus to the third order, the stability criterion becomes

$$\begin{aligned} (\Gamma/\Gamma_{PA})^2 = -\frac{8}{x_w^2} - 4 \left[2 + (\epsilon\beta_p)^2 \right] \delta_2 \cos\alpha_2 - \frac{113}{6} \epsilon\beta_p \delta_3 \cos\alpha_3 \\ - 4 (\epsilon\beta_p)^2 \delta_4 \cos\alpha_4 + \epsilon \frac{\epsilon\beta_p}{2} (1 + 3 \epsilon\beta_p) < 0. \end{aligned} \quad (30)$$

The first term in Eq. (3) describes the stabilizing effect of the conducting wall for the case where $x_w \sim 0$ ($\bar{\epsilon} \sim 1/2$). The second term includes

the effect of elliptical shaping of the plasma cross section in cylindrical and toroidal plasma geometries. The third and fourth terms describe the effect of triangular and rectangular shapings, respectively. The higher-order toroidicity destabilization is given by the final term. It should be noted that Fourier harmonics with $l \geq 5$ have no effect on stability to this order. The lower harmonics, which give rise to ellipticity ($l=2$), triangularity ($l=3$) and rectangularity ($l=4$), can be coupled by toroidicity and determine the vertical stability of non-circular plasmas. It should be noted that asymmetric components ($\delta_l^S \equiv \delta_l \sin \delta_l$) have no effect on vertical stability. The effect of toroidal coupling is summarized in Fig. 2. Standard 'D' shaping ($\delta_3^S \equiv \delta_3 \cos \alpha_3 > 0$) and tilted square shaping ($\delta_4^C > 0$) enhance stability, relative to inverse 'D' shaping ($\delta_3^C < 0$) and standard square shaping ($\delta_4^C < 0$) respectively.

In a similar manner we can obtain the lowest-ordered growth rate for a horizontal shift as

$$(\Gamma/\Gamma_{PA})_H^2 = (\Gamma/\Gamma_{PA})_V^2 - \frac{q^2}{2}, \quad (31)$$

where the safety factor

$$q = \left| \frac{a^2 I_O}{R_O \Psi_O} \right|,$$

and $(\Gamma/\Gamma_{PA})_V$ is the normalized growth rate for the vertical instability defined by Eq. (22). Equation (31) shows that slightly non-circular equilibria are always stable to a rigid horizontal displacement at high β .

As has been previously reported [1], Eq. (30) shows that elliptical shaping ($\delta_2 \cos \alpha_2 < 0$) to vertical instability even in the limiting case of low- β cylindrical plasmas ($\epsilon\beta_p \rightarrow 0$). The Shafranov shift of the magnetic axis in the high- β toroidal case leads to a coupling with the elliptical shaping which destabilizes the vertical mode. Triangular ($l=3$) or rectangular ($l=4$) shaping has no effect in the limiting case of low beta ($\epsilon\beta_p \rightarrow 0$). On the other hand, Rosen [3] has shown that cylindrical rectangular shaping has an effect on the vertical stability with a rigid displacement in the δ^3 order. This discrepancy is due to the difference in definition of the plasma boundary and the expansion ordering. Unlike in Rosen's case, the energy integrals $\delta W^{\delta\delta}$ and $\delta W^{\epsilon\delta\delta}$ turned out to be zero in our calculation. Toroidal $l=3$ and $l=4$ components results in similar behavior to the $l=2$ component. In higher order, toroidicity is destabilizing due to the increasing shift of the magnetic axis and the resultant vertical elongation of the magnetic surfaces near the axis [13].

Within our ordering, asymmetric shaping has no effect when the symmetric shaping component (δ_λ^c) is held constant. However, if the total shaping coefficient $\sqrt{(\delta_\lambda^c)^2 + (\delta_\lambda^s)^2}$ is held constant, the configuration tends toward marginal stability as δ_λ^s is increased.

4. COMPARISON WITH NUMERICAL CALCULATIONS

We now compare our analytic results with numerical calculations using the PEST code [11,13]. In order to accomplish this, the PPL equilibrium code [14]

was adopted to generate the plasma shape of Eqs. (1) and (2). The basic parameters of our equilibria are quite similar to those of the INTOR design, except for the shape of the plasma cross section.

In Fig. 3, poloidal plane projections of axisymmetric mode displacement vectors are shown for a medium β_p case ($\beta_p = 1$) and a high β_p case ($\beta_p = 1/\epsilon$). It should be noted that, particularly at high- β , the numerical axisymmetric mode is not incompressible. In addition a pure vertical displacement is not the most unstable perturbation. Thus, our rigid-mode analysis will always return optimistic results. On the other hand, the theoretical model has an almost uniform current profile which leads to pessimistic results. As discussed in the previous section, the outward shift of the magnetic axis due to high β leads to a deformation of the magnetic surface inside the plasma and makes the total system unstable. In the low β_p case, the displacement of the unstable mode is uniform and almost vertical. However, in the high β case the unstable mode is localized near the magnetic axis.

Comparisons between the rigid-mode analysis and the FEST code results are shown for several shapes in Fig. 4. The theoretical results are nearly two times greater than the computational results, which may be mainly due to the differences between analytic and computational current profiles. The discrepancy between the analytical and numerical studies is quite large in the inverse-D shaping case because the unstable mode is extremely localized near the elliptic magnetic axis, because the displacement of the most unstable axisymmetric mode is not purely vertical, and because the mode is compressive (Fig. 5). In that case, by decreasing β_p slightly, we can get more reasonable agreement between the theoretical and numerical analyses.

5. SUMMARY

We have analytically studied the vertical stability of slightly non-circular tokamak equilibria with vertical asymmetry and compared the results with numerical calculations using the PEST code. It is concluded that:

- (1) Toroidal high- β equilibria with vertically asymmetric cross sections were obtained using the double expansion in shape and toroidicity. The lowest order energy integral shows that the horizontal rigid-mode is always stable while the vertical rigid mode is marginally stable in the absence of a stabilizing conductive wall.
- (2) To the next order, growth rate of the vertical instability depends on a coupling between toroidicity and shaping.
- (3) Inverse 'D' ($\delta_3^C < 0$) and standard square ($\delta_4^C < 0$) shapes are unstable to rigid vertical displacements.
- (5) Asymmetric shaping has no dramatic effect on vertical instability.

ACKNOWLEDGMENT

This work was supported by the U.S. Department of Energy Contract No. DE-AC02-76-CHO-3073.

REFERENCES

- [1] WESSON, J. A., Nucl. Fusion 18 (1978) 87.
- [2] RUTHERFORD, P. H., Princeton PPL Report MATT-976 (1973).
- [3] ROSEN, M.D., Phys. Fluids 18 (1975) 482.
- [4] HAAS, F.A., Nucl. Fusion 15 (1975) 407.
- [5] WAKATANI, M., J. Phys. Soc. Jpn. 38 (1975) 1555.
- [6] INTOR, International Tokamak Reactor, Zero Phase IAEA, Vienna, 1980.
- [7] MEADE, D., et al., in Plasma Physics and Controlled Nuclear Fusion Research (Proc. 8th Int. Conf. Brussels, 1980) CN-38/X-1.
- [8] NAGAMI, M., et al., in Plasma Physics and Controlled Nuclear Fusion Research (Proc. 8th Int. Conf. Brussels, 1980) CN-38/PD.
- [9] DALHED, H.E., GRIMM, R.C., JOHNSON, J.L., Princeton Report PPPL-1766 (1981).
- [10] TODD, A.M.M., MILLER, A.E., GRIMM, R.C., OKABAYASHI, M., DALHED, H.E., Princeton Report PPPL-1787 (1981).
- [11] GRIMM, R.C., GREENE, J.M., JOHNSON, J.L., in Methods in Computational Physics 16 (Killeen, J., Ed.) Academic Press, New York (1976) 253.
- [12] YAMAZAKI, K., UCHIDA, T., Jpn J. Appl. Phys. 18 (1979) 981.
- [13] TODD, A.M.M., MANICKAM, J., OKABAYASHI, M., et al., Nucl. Fusion 19 (1979) 743.
- [14] JOHNSON, J.L., DALHED, H.E., GREENE, J.M., et al., J. Comput. Phys. 32 (1979) 212.

FIGURE CAPTIONS

Fig. 1 Coordinate systems.

(R, ϕ, Z) : cylindrical coordinates,

(r, ϕ, θ) : circular quasi-toroidal coordinates,

(ρ, ϕ, θ) : non-circular quasi-toroidal coordinates.

Fig. 2 Effects of shaping components of the plasma cross section on the vertical instability, only due to the coupling between shaping and toroidicity.

Fig. 3 Comparison of the axisymmetric mode displacement between a medium β case and a high β case.

(a) $\beta_p = 1, \Gamma/\Gamma_{pA} = 0.52$

(b) $\epsilon\beta_p = 1, \Gamma/\Gamma_{pA} = 1.30$

Fig. 4 Normalized growth rate, Γ/Γ_{pA} , of axisymmetric instability for vertically elongated plasmas ($\delta_2^C = -0.2$) with different shapes.

Solid line: rigid mode analysis (for any δ_3^S),

Broken line: PEST code results,

for symmetric shaping ($\delta_3^S = 0$),

and asymmetric shaping ($\delta_3^S = 0.1$) .

Points (a), (b), and (c) denote the shaping cases of Fig. 5 and a point (b*) represents the case (b) of Fig. 2.

Fig. 5 Poloidal projections of the displacement for

- (a) inverse 'D' shaping ($\delta_2^C = -0.2$, $\delta_3^C = -0.1$, $\delta_3^S = 0$) ,
- (b) asymmetric shaping ($\delta_2^C = -0.2$, $\delta_3^C = 0$, $\delta_3^S = 0.1$) , and
- (c) standard 'D' shaping ($\delta_2^C = -0.2$, $\delta_3^C = 0.1$, $\delta_3^S = 0$) .

81X0820

Fig. 1

Fig. 2

Fig. 3

Fig. 4

Fig. 5