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PROBLEMS WITH THE CONCEPT OF PLASMA EQUILIBRIUM IN TOKAMAKS*

B. A. Carreras

Oak Ridge National Laboratory

Oak Ridge, Tennessee 37831, U.S.A.

ABSTRACT

The equilibrium condition for a magnetically confined plasma is normally formulated in terms of macroscopic equations. In these equations, the plasma pressure is assumed to be a function of the magnetic flux with continuous derivatives. However, in three-dimensional systems this is not necessarily the case. Here, we look at the case of an intrinsically three-dimensional realistic tokamak, and we discuss the possible interconnection between the equilibrium and anomalous transport.

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B. A. CARRERAS

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1. INTRODUCTION

The equilibrium condition for a magnetically confined plasma is normally formulated in terms of the single-fluid magnetohydrodynamic equations; cf. Freiberg [1987]. In equilibrium, these equations reflect the force balance between the magnetic and kinetic forces

$$\vec{J} \times \vec{B} = \nabla p \tag{1}$$

The magnetic field, $\vec{B}$, and current density, $\vec{J}$, obey Maxwell's equations. In stationary conditions, they are

$$\nabla \times \vec{B} = \mu_0 \vec{J} \tag{2}$$

$$\nabla \cdot \vec{B} = 0 \qquad (3)$$

To solve the equilibrium problem, it is necessary to add the adequate boundary conditions. This problem is fully three-dimensional (3-D) and nonlinear and, in the most general case, has no solution. When the magnetic configuration is invariant under a continuous group of transformations, the dimensionality of the problem can be reduced. In this case, there are well-developed methods to solve this system of equations. For instance, in the case of invariance under a uniparametric group of transformations, the problem becomes two-dimensional (2-D). In two dimensions, the system of equations (1) to (3) can be reduced to a second-order partial differential equation, the Grad-Shafranov equation; cf. Grad & Rubin [1958]; Shafranov [1966]. The form of this equation depends on the symmetry of the configuration. Examples are axisymmetry (ideal tokamak), helical symmetry (straight stellarator), axial symmetry (bumpy cylinder). Under some assumptions on the explicit form of the pressure profile, the existence and uniqueness of solutions for the Grad-Shafranov equation can be proved; cf. Mossino [1992]; Diaz [1992].

Throughout this meeting, all the presentations have followed the macroscopic equilibrium formulation, which assumes that the gradient of the plasma pressure is a function of the magnetic flux and that it has continuous derivatives. The question discussed in this paper is whether this assumption is reasonable. This question was brought up in a seminal paper by H. Grad [1967]. The problem comes about because of the intrinsic 3-D character of the toroidal magnetic configurations. The approximation of a 3-D configuration by a closely similar 2-D configuration plus a small perturbation was seriously questioned by Grad. Several types of approaches based on the averaging method (cf. Bogolyubov & Mitropolskii [1961]; Greene & Johnson [1961]; Garcia [1992]) have been successful in describing some of the large-scale properties of plasmas, such as the shapes of the magnetic surfaces and the magnetic axis shift with β. These approaches have been tested numerically under some restrictive mathematical assumptions; cf. Carreras et al. [1983], Johnson [1986], Hender et al. [1985]. However, the problem remains for the detailed description of 3-D equilibria. The non-existence of such equilibria, at least locally, could be one of the causes of the anomalous transport observed in toroidal confinement systems.

To theoretically analyze the equilibrium problem, it is necessary first to understand the magnetic field line structure. The best way of understanding this structure is to look at the magnetic field line flow as a Hamiltonian flow. Then general theorems of Hamiltonian systems such as Liouville's theorem, Noether's theorem, and the Kolmogorov-Arnold-Mosses (KAM) theorem (cf. Arnold & Avez [1968]) can be applied to the magnetic field structure. The great progress made in understanding these dynamical systems since Grad's basic paper allows a more detailed analysis of the equilibrium problem. However, the potential link between the lack of equilibrium in a 3-D configuration and the anomalous transport losses remains an outstanding question.

The rest of this paper is organized as follows. In Sec. 2, we discuss how the flow of magnetic field lines can be represented as a Hamiltonian system. The simple example of a periodic cylinder is presented in Sec. 3. The status of the problem for the tokamak configuration is reviewed in Sec. 4.

2. THE MAGNETIC FIELD LINE AS A HAMILTONIAN FLOW

For a given magnetic field, $\vec{B}$, the magnetic field line equations in Cartesian coordinates are

$$\frac{dx}{B_x} = \frac{dy}{B_y} = \frac{dz}{B_z} = \frac{d\ell}{|\vec{B}|} \tag{4}$$

Here, $d\ell$ is the element of length along the magnetic field line. These magnetic field line equations can be interpreted as the equations of motion of a particle, and the field lines as the trajectories of particles. The flow of magnetic field lines in physical space can be represented as a Hamiltonian system. This is not obvious from Eq. (4). The general proof was given by Cary & Littlejohn [1983], who used a noncanonical Hamiltonian theory of dynamical systems to show the correspondence between a Hamiltonian system and the magnetic field line flow. Here, we limit the discussion to some particular examples and use a convenient coordinate system to express Eq. (4) as a Hamiltonian system.

Any divergence-free vector field $\vec{B}$ in a toroidal configuration can be represented (cf. Cary & Littlejohn [1983]) as

$$\vec{B} = \nabla\Phi \times \nabla\Theta + \nabla\phi \times \nabla\Psi_{H} \tag{5}$$

where $\Phi \equiv B_0\rho^2/2$, with ρ a generalized radial coordinate, Θ is a generalized poloidal angle, and ϕ is the toroidal angle (Fig. 1). B_0 is a constant and is taken to be a characteristic value of the toroidal magnetic field. With this representation for the magnetic field, the magnetic field line equations are

$$\frac{\partial \Psi_{//}}{\partial \Theta} = -\frac{d\Phi}{d\phi} \tag{6}$$

$$\frac{\partial \Psi_{//}}{\partial \Phi} = \frac{d\Theta}{d\phi} \tag{7}$$

Now it is possible to interpret these equations as the equations of motion of a particle by identifying Θ as the particle position, q, Φ as its momentum, p, ϕ as the time, t, and $\Psi_{//}$ as the Hamiltonian, H. With this identification, Eqs. (6) and (7) are Hamilton's equations. Therefore, the magnetic field line equations in real space are equivalent to the particle motion equations in phase space.

To follow up the comparison between magnetic field lines and Hamiltonian systems, it is necessary to introduce the concept of rotational transform. If we follow a magnetic field line from a ϕ = constant plane once around the torus to the original ϕ = constant plane, the point of intersection of the field line with this plane has rotated an angle ι in the poloidal direction. This angle ι normalized to 2π is called the rotational transform, $ɫ$. In tokamak lore, we often use the safety factor $q = 1 / ɫ$ instead of $ɫ$. Hereafter we use $ɫ$ to avoid confusing the safety factor q with the particle position q. When $\Psi_{//}$ is a function only of ρ and with the magnetic field representation of Eq. (5), the rotational transform is given by $ɫ = d\Theta/d\phi$.

A useful way of visualizing the properties of a general dynamical system is the Poincaré representation. This representation is basically a 2-D cut of the trajectories in phase space. The trajectories are represented by the points of intersection with this 2-D surface. The same representation is used in studying magnetic field line systems. The plot resulting from the multiple intersections of a magnetic field line moving around the torus with a given ϕ = constant plane is the Poincaré plot. The Hamiltonian system (magnetic field system) is integrable when the system has a complete set of invariants; in this case all orbits (magnetic field lines) lie on invariant tori. For an integrable system, the Poincaré plot is a set of closed nested lines. Because of the dimensionality of the magnetic field line flow, the relevant invariant is the energy. That is, when the magnetic field line equations are integrable, a coordinate system such that $\Psi_{//}$ is a function only of ρ exists. The surfaces defined by $\Psi_{//}$ = constant are the invariant tori. These surfaces are called magnetic surfaces. When ι is irrational, the magnetic field line ergodically covers a whole magnetic surface. When ι is rational, $\iota = n/m$, the field line closes on itself after going m times toroidally and n times poloidally. In the latter case the orbit is called periodic, and the corresponding magnetic surface is called a resonant surface.

3. A SIMPLE EXAMPLE: A PERIODIC CYLINDER

Let us consider a simple example, the case of a periodic cylinder. We use the standard geometrical cylindrical coordinates (r, θ, z). In this geometry, the magnetic field is

$$\vec{B} = B_\theta(r)\vec{\theta} + B_z\vec{z}\,, \tag{8}$$

where $\vec{\theta}$ and $\vec{z}$ are the unit vectors in the poloidal direction and in the direction along the axis of the cylinder, respectively. Because of the symmetry, the coordinates θ and $\varphi \equiv 2\pi z/L$ are ignorable and the physical quantities are only functions of the radius r. Here, L is the length of the cylinder. We assume that B_z is constant. Using this magnetic field, the solution of the magnetic field line equations is:

$$\left.\begin{aligned} r &= r_0 \\ \theta &= \frac{B_\theta(r)}{rB_z}\frac{L}{2\pi}\varphi + \theta_0 \end{aligned}\right\} \tag{9}$$

where r_0 and θ_0 are integration constants. The rotational transform is $\iota = LB_\theta/(2\pi rB_z)$. It is clear from this model that the magnetic surfaces are concentric cylinders (Fig. 2a). For a rational value of ι the field line closes on itself, and for an irrational value the field line covers the whole cylindrical surface. Comparing with the magnetic coordinates that we have introduced in Eq. (5), $\rho = r$, $\Theta = \theta$, and $\phi = \varphi$. Here we have taken $B_z = B_0$. The Hamiltonian is

$$\Psi_H(r) = \Psi_H(0) + \frac{L}{2\pi}\int_0^r B_\theta(r) dr = \Psi_H(0) + B_0 \int_0^r ɨ(r) r \; dr \tag{10}$$

The dynamical system analogous to the magnetic field line flow of Eq. (9) is

$$\left. \begin{aligned} p &= p_0 \\ q &= \omega t + q_0 \end{aligned} \right\} \tag{11}$$

These are the equations of an unperturbed oscillator with characteristic frequency $\omega = ɨ$.

Let us consider a singular surface, $ɨ(r_S) = n/m$, where m and n are integers. The periodic orbit is characterized by a frequency $ɨ(r_S) = n/m$. A nearby magnetic surface has a characteristic frequency

$$ɨ \approx \frac{n}{m} + \hat{S}\frac{x}{r_s}, \tag{12}$$

where $\hat{S}$ is the magnetic shear parameter, $\hat{S} = r_s dɨ / dr|_{r=r_s}$, and $x = r - r_S$.

Let us assume that the cylindrical symmetry of the magnetic field is broken by introducing a magnetic field component proportional to $\cos(m\theta + n\varphi)$. An easy way to introduce the symmetry-breaking term is by adding a perturbation to the Hamiltonian, $\Psi_H(r,\theta,\varphi) = \Psi_{H0}(r) + \tilde{\Psi}(r)\cos(m\theta + n\varphi)$. The coordinates θ and φ are no longer ignorable. The Hamiltonian is time (ϕ) dependent; therefore, the energy is not conserved. However, it is possible to change coordinates and show that the system is integrable. This is because the system still has a symmetry. It is not 3-D, it has helical symmetry. By introducing a new angle variable,

$$\eta = \theta + \frac{n}{m}\varphi \qquad (13)$$

the magnetic field, Eq. (5), can be written as

$$\vec{B} = \nabla\Phi \times \nabla\eta + \nabla\phi \times \left[\nabla\left(\Psi_{ll} - \frac{n}{m}\Phi\right)\right] \qquad (14)$$

The change of coordinate implies that we can introduce a new effective Hamiltonian,

$$\Psi_{ll}^{*}(r,\eta) = \Psi_{ll0}(r) - \frac{n}{m}\Phi(r) + \tilde{\Psi}(r)\cos(m\eta) \qquad (15)$$

which is time independent. From Eq. (15) it is clear that the magnetic surfaces are no longer circular cylinders. Their topology changes near $\iota = n/m$ (Fig. 2b) with the formation of a magnetic island. This Hamiltonian, Eq. (15), can be expanded near the $\iota = n/m$ surface. Using Eq. (12), we obtain

$$\Psi_{ll}^{*}(r,\eta) \approx \Psi_{ll0}^{*}(r_s) + B_0\hat{S}\frac{x^2}{2} + \tilde{\Psi}(r_s)\cos(m\eta) \qquad (16)$$

This is the Hamiltonian of a pendulum. The magnetic surfaces inside the island correspond to the oscillating solutions of the pendulum. The island x-point corresponds to the upside-down unstable vertical position of the pendulum. Outside the island, the pendulum is rotating freely. From Eq. (16), the width, W, of the magnetic island can be estimated, and one obtains $W^2 \approx 4\left(\tilde{\Psi} / B_0\right) / \hat{S}$.

When the magnetic configuration has no continuous symmetry, the magnetic perturbation added to the cylindrically symmetric magnetic field is a general function of the three coordinates. No longer there is a simple transformation of the Hamiltonian to make it independent of φ (time independent), and magnetic flux (energy) is not an integral of the system of equations (4). The magnetic field line equation is not integrable, and the KAM theorem is the only guide for the existence of magnetic surfaces. The structure of the magnetic surfaces can be rather complicated (Fig. 3).

In the most general case and for small deviation from a symmetric configuration, it is possible to analyze the magnetic field structure in the neighborhood of each rational surface using a near-resonant-surface expansion and including only the resonant Fourier components. In this way, we can associate a magnetic island width with each surface. When nearby magnetic islands overlap, the magnetic field becomes stochastic (Chirikov criterion); cf. Chirikov [1979].

This analysis of the magnetic configurations leads to practical prescriptions for constructing magnetic confinement devices. In the case of stellarators, the magnetic field is solely generated by external conductors. These conductors must be carefully designed to ensure the existence of magnetic surfaces. Furthermore, the site of the region with unbroken magnetic surfaces should be maximized to obtain maximum efficiency in the use of the magnetic fields. This can be done using the method of stochasticity reduction developed by Cary & Hanson [1986]. The method aims at the elimination of magnetic field line stochasticity by reducing the magnetic island size in each flux surface. The parameters used to minimize the magnetic island width are those defining the coil system. The method has been applied in designing low-aspect-ratio stellarators; cf. Carreras et al. [1988].

4. TOKAMAK EQUILIBRIUM

In a tokamak, the plasma is the secondary of a transformer. The transformer induces a current in the plasma in the toroidal direction, and a poloidal field component is generated by this plasma current. To stabilize the plasma, a toroidal field is added. In an ideal tokamak, this toroidal field is created by an infinitely long conductor at the center of the torus (Fig. 4). This magnetic configuration is perfectly axisymmetric. A real tokamak cannot be built this way. In practice, the toroidal magnetic field is created by a finite set of toroidal coils. The finite set of coils gives bumpiness in the toroidal direction (toroidal ripple), and the axisymmetry is broken. Therefore, any real tokamak intrinsically constitutes a 3-D magnetic configuration.

Even if it were possible for a tokamak to maintain exact axisymmetry of the external fields, the magnetic configuration would not necessarily have this symmetry. As suggested by B. B. Kadomtsev [1991], there is the possibility of a spontaneous breakdown of symmetry. This effect is similar to other physics phenomena that are characterized by a Hamiltonian with an exact symmetry and a set of ground states with broken symmetry. In each magnetic flux surface, chains of microscopic magnetic islands (with island width of the order of the plasma skin depth, c/ω_{pe}) could be always present and the axisymmetry would be intrinsically broken. The motivation for advancing the hypothesis of spontaneous breakdown of axisymmetry is that the presence of these islands could explain the enhancement of electron heat and particle losses over the classical collision losses. The theoretical foundations of such a hypothesis are not clear. It is a serious problem to explain why the magnetic current filaments associated with these magnetic islands do not decay in a resistive time. Kadomtsev has advanced an explanation based on ion energy pumping. This requires the islands to be very small, much smaller than the ion Larmor radius. Therefore, this mechanism would not explain the long correlation length of the fluctuations in the tokamak plasma core. To counter this argument, the presence of larger magnetic islands in the plasma has also been considered. For this model, the bootstrap current gradient has been advanced as source of energy; cf. Hegna & Callen [1991].

Regardless of the relative merits of these transport models, they all introduce a breakdown of axisymmetry. From Eq. (1), $\vec{B}\cdot\nabla p = 0$. That is, the pressure is constant along a magnetic field line. Therefore, the presence of the magnetic islands means that the pressure is constant across the magnetic islands. In the case of island chains present at every rational value of ι in the plasma, the pressure profile should have flat spots at each rational value of ι (Fig. 5). Therefore, the pressure is not a differentiable function. In Fig. 5, the pressure is plotted as a function of the radius for a parabolic profile. The profile is flat at each rational value of ι, $\iota = n/m$ with n from 1 to 9, and ι varying between 1 at the origin and 1/2 at the edge. The usual parabolic profile has been plotted for reference (dashed line). If we consider higher n values, the number of steps increases. In the limit $n \rightarrow \infty$, the pressure is a fractal. Numerical techniques have been developed for the solution of a 3-D equilibrium in the presence of few islands; cf. Salas [1992]. A pressure profile with a few flat spots can be treated with these techniques. The question is how to treat the equilibrium problem when the pressure profile has a flat spot at every rational value of ι.

The assumption of spontaneous breakdown of symmetry in a real tokamak leads back to the full set of equations for the equilibrium problem in a 3-D geometry and with pressure profiles which have no derivatives. The situation is like that foreseen by Grad, with the problem of transport linked to the existence of equilibrium.

From an experimental point of view, it is not yet possible to test the assumption of the spontaneous breakdown of symmetry in tokamak plasmas. For a stellarator, the magnetic configuration is created by external windings. Therefore, for this configuration, it is possible to experimentally study the structure of the vacuum magnetic flux surfaces. The Poincaré mapping can be done by using an electron beam; cf. Anderson et al. [1992], Jaenicke et al. [1992], Ascasibar [1992]. With this technique, the presence of relatively large magnetic islands (about 1 cm or larger) can be detected. However, there is not enough resolution in the experiments to verify the micromagnetic island structure. The situation is even more complicated for tokamaks, in which the magnetic configuration is created by the plasma current and cannot be studied in the absence of plasma. So far, there are no direct measurements of magnetic perturbations in the core of a tokamak plasma. The experimental resolution in the measurement of temperature profiles is increasing but is not yet good enough to resolve the required length scales. However, magnetic islands in the plasma may be the reason for the choppy electron temperature profiles seen in the Joint European Torus (JET) (cf. Nave et al. [1992]) and TFTR (B. Grek et al. [1992]) in high-q operation.

In conclusion, the similarity in the transport properties of tokamaks and stellarators could be a consequence of the similarity of equilibrium or lack of it. By considering both types of configurations as intrinsically 3-D, the existence of local equilibrium is brought into question. Pressure profiles should not be considered as differentiable functions of ρ, but as fractals with flat spots at each rational value of the rotational transform. The mathematical formulation of the equilibrium problem from this perspective needs to be undertaken.

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Benjamin A. Carreras

Oak Ridge National Laboratory

P. O. Box 2009

Oak Ridge, Tennessee 37831-8070

Figure captions

Fig. 1. Toroidal coordinate system used in this paper.

Fig. 2. Poincaré plot for a periodic-cylinder magnetic field configuration: (a) cylindrically symmetric case and (b) helically symmetric case.

Fig. 3. Magnetic field line plot for a fully 3-D configuration showing regions of stochastic field lines, magnetic islands, and unbroken flux surfaces.

Fig. 4. Ideal tokamak. The plasma is the secondary of the transformer, which induces a plasma current, J_ϕ, in the toroidal direction, and a poloidal field component, B_θ, generated by this plasma current. There is an added toroidal field, B_ϕ, created by an infinitely long conductor at the center of the torus.

Fig. 5. Pressure as a function of radius. The profile is flat at each rational value of ι, $\iota = n/m$ for $n = 1$ to 9, and ι varying between 1 at the origin and 1/2 at the edge. The usual parabolic profile has been plotted for reference (dashed line).

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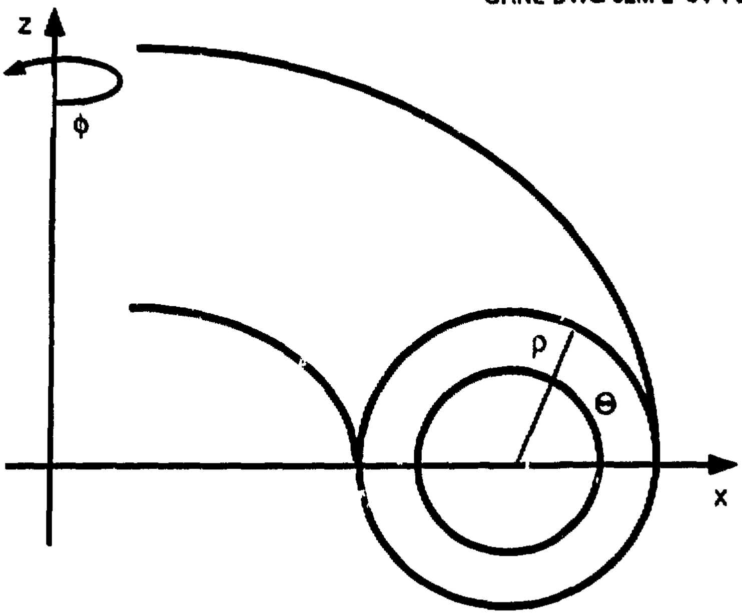

Fig. 1

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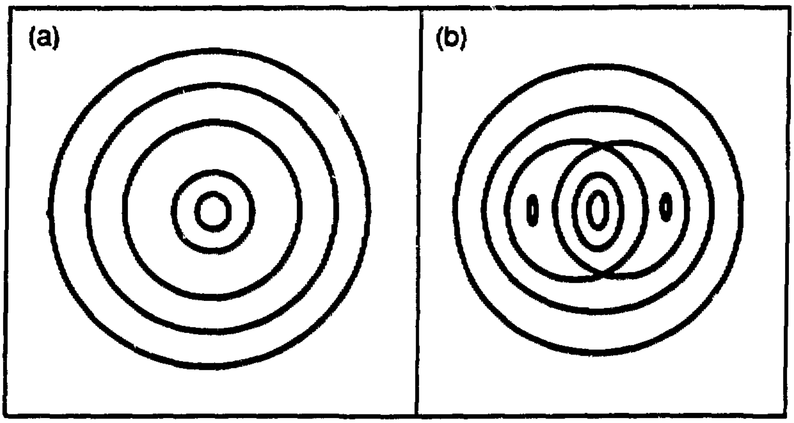

Fig. 2

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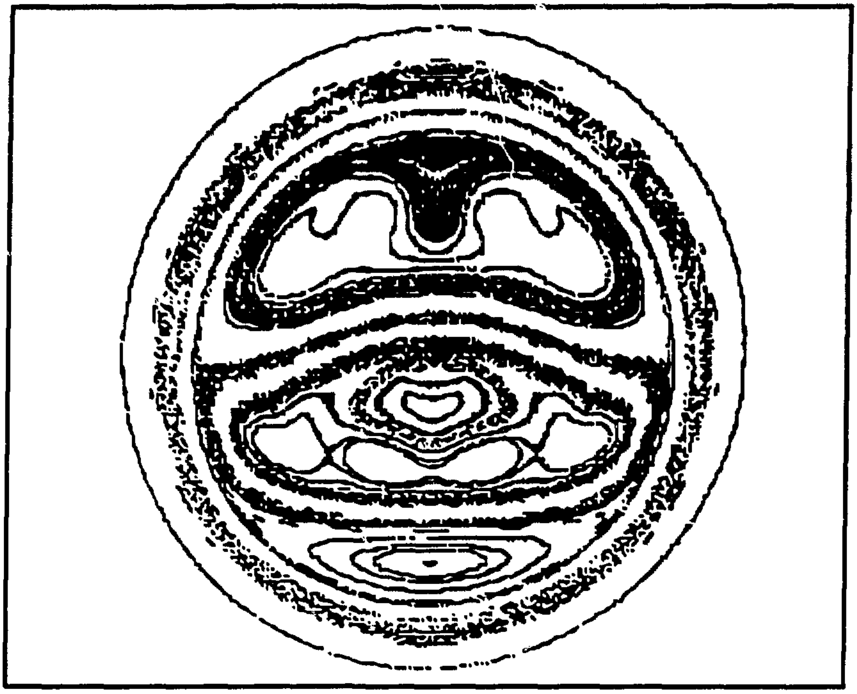

Fig. 3

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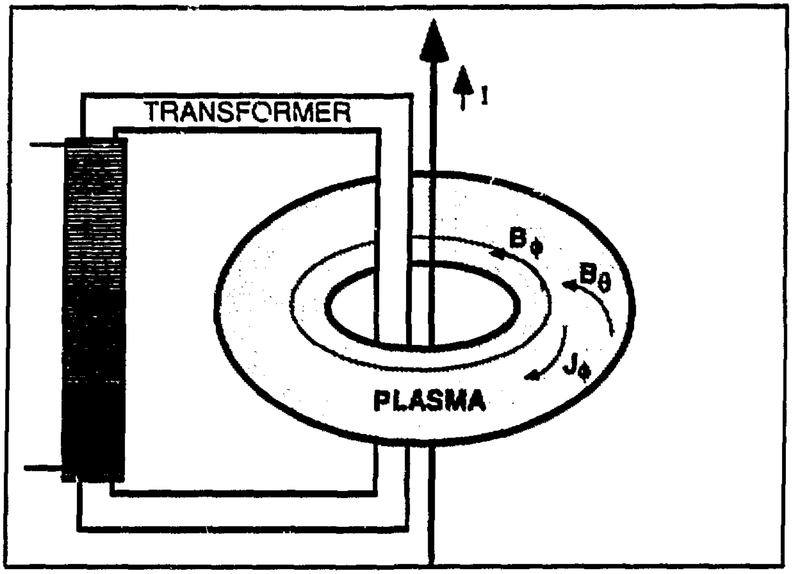

Fig. 4

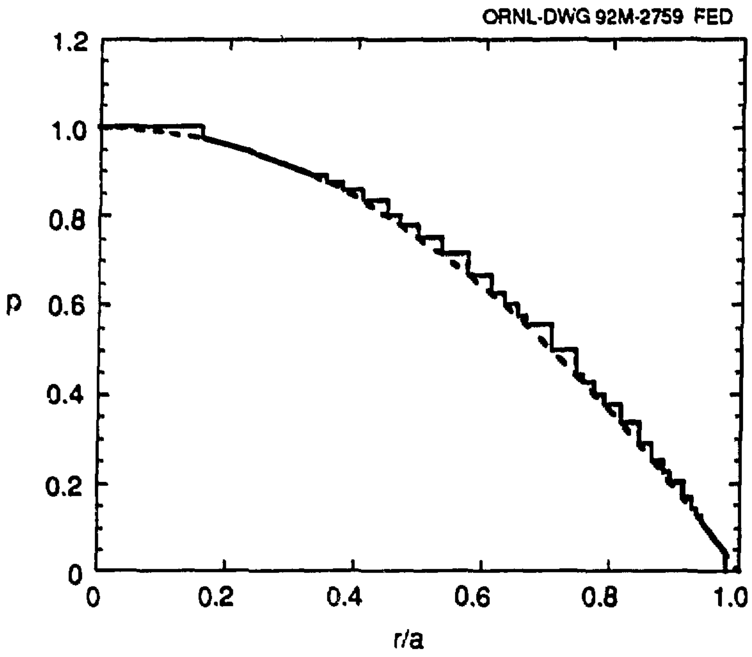

Fig. 5