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Magnetic helicity transport and reversed field pinch

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ABSTRACT

The Reversed Field Pinch plasma is described by a model based on magnetic helicity transport. It is shown that the mean correlation between fluctuations in fluid velocity and in magnetic field, $<\tilde{\underline{u}}\times\tilde{\underline{B}}>$, provides non-dissipative transport of helicity that persists whether or not the plasma is fully relaxed, as well as a non-Ohmic energy sink that vanishes when the fully relaxed state is attained. Such energy losses account for the anomaly in loop voltage or resistance seen in experiment, which can be interpreted as a loss of helicity associated with obstruction of parallel current flow and magnetic flux emerging from the plasma due to plasma-wall contact. Loss of magnetic flux through the bounding surface of the plasma enhances the helicity transport which increases the loop voltage necessary to sustain the toroidal current. The increase which is shown to be proportional to the flux intercepted and the local electron temperature compares well with that observed in experiments.

(Submitted for publication in Nuclear Fusion)

May 1987 (Revised Dec 1987)

INTRODUCTION

In Reversed Field Pinch (RFP) experiments [1], it is observed that the plasma magnetic field configuration never quite reach the fully relaxed minimum energy state satisfying $\nabla \times \underline{B} = \mu \underline{B}$ with μ uniform across the plasma, predicted by Taylor [2,3]. The departure from the fully relaxed state is indicated by the measured profile of μ which is not uniform but decreasing towards the plasma edge [4,5,6]. Such departure has been associated with the energy throughput in the system [7].

The RFP equilibrium configuration consists of nested helices. On the axis, the field is purely toroidal while at increasing radius the field lines form spirals which become more and more tightly wound until at the reversal surface the field is purely poloidal. Between the reversal surface and the plasma edge, field lines actually spiral backwards, justifying the reversed field designation bestowed upon this type of discharge. In the force-free configuration, current density $\underline{j}$ is everywhere parallel to magnetic field $\underline{B}$ ie. $\mu_0 \underline{j} = \mu \underline{B}$ by virtue of Ampere's Law, so there must be components of the current flowing perpendicular to and even opposite to the applied emf in a toroidal RFP. Evidently Ohm's Law, $\underline{E} + \underline{u} \times \underline{B} = \eta \underline{j}$, cannot be satisfied at the reversal surface when higher order effects are ignored. Nevertheless the field configuration of RFP discharges is sustained against resistive decay by the applied toroidal electric field to which poloidal current is somehow coupled and maintained. It has been suggested that interactions of magnetic field fluctuations can result in an effective electric field parallel to the magnetic field needed to sustain the field reversal [8,9,10]. In particular, the mean field magnetohydrodynamic theory invokes correlation between fluctuations in fluid velocity and field of the form $\tilde{\underline{u}}\times\tilde{\underline{B}}$. It has been shown [11] that the global magnetic helicity of a steady state plasma bounded by a perfectly conducting wall where all fluctuations vanish is conserved with the externally applied electric field balancing the resistive dissipation. It has also been shown [12,7] that the helicity balance is more correct than the magnetic energy balance in computing the plasma resistivity when fluctuations are ignored. Bhattacharjee and Hameiri [13] showed that, independently of the μ profile, fluctuations of the $\tilde{\underline{u}}\times\tilde{\underline{B}}$ type dissipate magnetic energy but not helicity. Assuming these properties, Boozer [14] obtained a functional form for the flux of helicity which depends on the gradient of μ and vanishes when $\nabla\mu = 0$ ie in the fully relaxed configuration. This appears to imply that the fully relaxed state is free from helicity transport and that a substained RFP cannot be fully relaxed. In the sustained RFP helicity is being injected continously by the applied toroidal electric field, and unless this input is locally balanced against dissipation everywhere, helicity transport must exist if the configuration is to remain stationary. Furthermore, the fact that toroidal magnetic field is reversed in the outer region of the plasma means that the applied toroidal emf is not injecting helicity at this location but absorbing it. Consequently helicity transport is called for in a sustained RFP even when the plasma is fully relaxed. If transport depended on a non-vanishing gradient of μ, then the fully relaxed state could never be realized in a sustained RFP.

An alternative view advanced in this paper is that departure from the fully relaxed state is a consequence of non-resistive energy loss due to the $\tilde{\underline{u}}\times\tilde{\underline{B}}$ fluctuations as they act to redistribute helicity. It will be shown that if there were no such energy loss, a sustained RFP would attain the fully relaxed Taylor state with $\nabla\mu=0$, in the presence of the non-vanishing $\tilde{\underline{u}}\times\tilde{\underline{B}}$ needed to redistribute helicity.

In the model advocated here, power can be thought of as flowing into the plasma through two channels. Flow through the first balances the resistive dissipation of the plasma current giving rise to the classical heating of the electrons. The second drives $\tilde{\underline{u}}\times\tilde{\underline{B}}$ fluctuations responsible for the helicity transport. Non-resistive energy loss from the fluctuations, manifested by the departure of the field configuration from the fully relaxed state, may account for the ion heating . . Fluctuating radial velocity on the plasma surface can also give rise to such eneryg loss. It is suggested that variations in resistivity profile alone cannot inhibit the plasma from attaining the fully relaxed state.

In the HBTX experiment [15], the resistivity calculated from the global helicity balance equation is higher than the Spitzer resistivity based on the electron temperature and the known concentration of impurities by more than an order of magnitude in some cases. In this paper it is shown that such differences can be attributed to the leaking of helicity from the plasma edge region where field lines are intersecting the vacuum vessel wall and the limiters. A model to describe such leakage is presented and the theoretical results are compared to data from HBTX. The helicity transport activity needed to compensate the leakage enhances non-resistive energy loss from the fluctuations. Similarly, an increase of resistivity near the edge of the plasma which may occur in some conditions [16] leads to additional helicity transport and the associated fluctuations will result in enhanced energy throughput in the second energy flow channel; this process influences the global energy balance much more than the small local resistive dissipation of energy at the edge.

THE ROLE OF FLUCTUATIONS IN HELICITY AND ENERGY TRANSPORT

When there are fluctuations [17,18,19,4] the Ohm's law has to be modified to include their contributions [see, for example,14]. With variables separated into the mean and fluctuating components such as $\underline{E} = \underline{E}_0 + \tilde{\underline{E}}$, the Ohm's law becomes

$$\underline{E}_0 = \eta \underline{j}_0 - \langle \tilde{\underline{u}} \times \tilde{\underline{B}} \rangle \tag{1}$$

$$\tilde{\underline{E}} = \eta \tilde{\underline{j}} - \tilde{\underline{u}} \times \underline{B}_0 \tag{2}$$

where $\langle\ \ \rangle$ denotes the appropriate average over time and poloidal and toroidal space and $\underline{E}_0 = \langle \underline{E} \rangle$. For simplicity, plasma mean flow $\underline{u}_0$ and resistivity fluctuation $\tilde{\eta}$ are taken to be zero.

The global helicity balance [20,21] of a steady state RFP bounded by a perfectly conducting wall [22,12,7] is

$$V_\phi \Phi(a) = \int_{V_a} \langle \underline{E} \cdot \underline{B} \rangle \, d^3x$$

and with fluctuations,

$$V_\phi \Phi(a) = \int_{V_a} \underline{E}_0 \cdot \underline{B}_0 \, d^3x + \int_{V_a} \langle \underline{\tilde{E}} \cdot \underline{\tilde{B}} \rangle \, d^3x \tag{3}$$

where V_a is the plasma volume and V_ϕ and $\Phi(a)$ denote the toroidal loop voltage and flux. Realizing that in steady state, $\underline{E}_0 = (V_\phi/2\pi R)\underline{\phi}$ where $\underline{\phi}$ is the toroidal unit vector and R the major radius,

$$\int_{V_a} \underline{E}_0 \cdot \underline{B}_0 \, d^3x = \int_{V_a} \frac{V_\phi}{2\pi R} (\underline{B}_0 \cdot \underline{\phi}) \, d^3x = V_\phi \, \Phi(a)$$

identically, implying that $\int_{V_a} \langle \underline{\tilde{E}} \cdot \underline{\tilde{B}} \rangle d^3x \equiv 0$. From eq (2), this can be written as

$$\int_{V_a} \langle \underline{\tilde{E}} \cdot \underline{\tilde{B}} \rangle d^3x = \int_{V_a} \eta \langle \underline{\tilde{j}} \cdot \underline{\tilde{B}} \rangle d^3x + \int_{V_a} \langle \underline{\tilde{u}} \times \underline{\tilde{B}} \rangle \cdot \underline{B}_0 d^3x \equiv 0. \tag{4}$$

Experiments [18,19,4] indicate that $|\tilde{B}|/B_0 \sim 10^{-2}$ and the perpendicular scale length of $\tilde{B}$ is $\Lambda/a \sim 1/25$. From $\mu_0 \underline{j}_0 = \mu \underline{B}_0$ and $\mu_0 \underline{\tilde{j}} = \nabla \times \underline{\tilde{B}}$,

$$\frac{|\tilde{j}|}{j_0} \sim \frac{a}{\Lambda} \frac{1}{\mu a} \frac{|\tilde{B}|}{B_0}$$

where $\mu a \sim 1.5$ typically. From these estimates,

$$\frac{|\underset{\sim}{\tilde{j}}\cdot\underset{\sim}{\tilde{B}}|}{|\underset{\sim}{j}_0\cdot\underset{\sim}{B}_0|} \sim \frac{a}{\Lambda}\,\frac{1}{\mu a}\,\frac{|\underset{\sim}{\tilde{B}}|^2}{B_0^{\,2}} \sim 10^{-3}$$

implying that $\int_{V_a} \eta\langle\underset{\sim}{\tilde{j}}\cdot\underset{\sim}{\tilde{B}}\rangle d^3x = 0$ when compared to $\int_{V_a} \eta\underset{\sim}{j}_0\cdot\underset{\sim}{B}_0\, d^3x$.

Accordingly, with eq (4),

$$\int_{V_a} \langle\underset{\sim}{\tilde{u}}\times\underset{\sim}{\tilde{B}}\rangle\cdot\underset{\sim}{B}_0\, d^3x = 0 \qquad (5)$$

though the integrand is not necessarily small. The interpretation of eq (5) is that helicity is transported through the action of the fluctuations from one region of the plasma to another with approximately no dissipation and that the transport does not affect its global balance.

In order to reveal the role of fluctuations in transporting helicity, the magnetic helicity balance equation for a sub-volume bounded by a mean flux surface inside the plasma volume needs to be derived first. In addition to the average operator $\langle\quad\rangle$ which denotes averages over the mean flux surface and the appropriate time, a second average operator $\{\quad\}$ which denotes averages over the mean flux surface but not time is introduced to extract the non-spatial fluctuating component. Expressed in these averages, a variable can be separated into a mean and two fluctuating parts such as $\underset{\sim}{A} = \langle\underset{\sim}{A}\rangle + \{\underset{\sim}{\tilde{A}}\} + \underset{\sim}{\tilde{A}}'$ or $\underset{\sim}{A} = \underset{\sim}{A}_0 + \underset{\sim}{\tilde{A}}$ where $\underset{\sim}{A}_0 = \langle\underset{\sim}{A}\rangle$ and $\underset{\sim}{\tilde{A}} = \{\underset{\sim}{\tilde{A}}\} + \underset{\sim}{\tilde{A}}'$. In terms of Fourier modes with spatial mode numbers (m,n), the two fluctuating terms $\{\underset{\sim}{\tilde{A}}\}$ and $\underset{\sim}{\tilde{A}}'$ correspond to the (m,n)=(0,0) mode and the higher spatial harmonic modes respectively. Following Bevir and Gray [21], the mean relative helicity which excludes the linkage between the external poloidal flux and the internal toroidal flux for a sub-volume v bounded by a mean flux surface s is defined as

$$\langle K\rangle = \int_V \langle \underline{A}\cdot\underline{B}\rangle\, d^3x \;-\; \langle \oint_S \underline{A}\cdot d\underline{\theta}\, \oint \underline{A}\cdot d\underline{\phi}\, \rangle \tag{6}$$

where $\underline{\theta}$ and $\underline{\phi}$ are the poloidal and toroidal unit vectors on s. This expression defines the magnetic helicity content within any mean flux surface.

The time rate change of the mean helicity in v is

$$\langle \dot{K}\rangle = 2\int_V \langle \dot{\underline{A}}\cdot\underline{B}\rangle d^3x - \int_S \langle \underline{A}x\dot{\underline{A}}\rangle\cdot(d\underline{\theta}xd\underline{\phi})$$

$$- \langle \oint_S \dot{\underline{A}}\cdot d\underline{\theta}\, \oint \underline{A}\cdot d\underline{\phi} + \oint \underline{A}\cdot d\underline{\theta}\, \oint \dot{\underline{A}}\cdot d\underline{\phi}\, \rangle$$

In terms of the mean and fluctuating parts,

$$\langle \dot{K}\rangle = 2\int_V \dot{\underline{A}}_0\cdot\underline{B}_0\, d^3x + 2\int_V \langle \dot{\tilde{\underline{A}}}\cdot\tilde{\underline{B}}\rangle d^3x - \int_S \langle \tilde{\underline{A}}'x\dot{\tilde{\underline{A}}}'\rangle\cdot(d\underline{\theta}xd\underline{\phi})$$

$$- 2\Phi_0(r)\dot{\psi}_0(r) - 2\langle \tilde{\Phi}(r)\dot{\tilde{\psi}}(r)\rangle \tag{7}$$

where $\Phi_0(r)=\oint_S \underline{A}_0\cdot d\underline{\theta}$, $\psi_0(r)=\oint_S \underline{A}_0\cdot d\underline{\phi}$, $\tilde{\Phi}(r)=\oint_S \{\tilde{\underline{A}}\}\cdot d\underline{\theta}$ and

$\tilde{\psi}(r)=\oint_s \{\underline{\tilde{A}}\}\cdot d\underline{\phi}$. With $\underline{E} = -\underline{\dot{A}} - \nabla\chi$ and the assumption that the mean field lines close on themselves, $<\dot{K}>$ becomes

$$<\dot{K}> = -2\int_V \underline{E}_0\cdot\underline{B}_0 d^3x - 2\int_V <\underline{\tilde{E}}\cdot\underline{\tilde{B}}> d^3x - 2\Phi_0(r)\dot{\psi}_0(r) - 2H(r) \tag{8}$$

where

$$H(r) = <\tilde{\Phi}(r)\dot{\tilde{\psi}}(r)> + \int_S <\tilde{\chi}\underline{\tilde{B}}>\cdot(d\underline{\theta}\times d\underline{\phi}) + \tfrac{1}{2}\int_S <\underline{\tilde{A}}'\times\underline{\dot{\tilde{A}}}'>\cdot(d\underline{\theta}\times d\underline{\phi}) \tag{9}$$

represents the rate of change of helicity in v due to an outward helicity flow across the surface s as a result of the actions of fluctuations. The first term in eq (9) corresponds to the helicity injection current drive by oscillating field [21].

In steady state, $<\dot{K}>=0$ and $\dot{\psi}_0(r)=V_\phi$, a constant, so that eq (8) becomes

$$V_\phi\Phi_0(r) = \int_V \underline{E}_0\cdot\underline{B}_0 d^3x + \int_V <\underline{\tilde{E}}\cdot\underline{\tilde{B}}> d^3x + H(r) \tag{10}$$

which, in analogous to eq (3), is the steady state helicity balance equation for a sub-volume. From eq (2) and the analogous identity to that demonstrates above, $\int_V \underline{E}_0\cdot\underline{B}_0\, d^3x = V_\phi\Phi_0(r)$, eq (10) reduces to the sub-volume analogy to eq (4).

$$H(r) = -\int_V \langle \underline{\tilde{u}} \times \underline{\tilde{B}} \rangle \cdot \underline{B}_0 \, d^3x - \int_V \eta \langle \underline{\tilde{j}} \cdot \underline{\tilde{B}} \rangle \, d^3x \tag{11}$$

which relates helicity transport to the reversal sustainment by $\underline{\tilde{u}} \times \underline{\tilde{B}}$ fluctuations.

When the plasma is bounded by a perfectly conducting wall (with gaps to allow mean field penetration), the radial component of $\underline{\tilde{B}}$ is $\tilde{B}_r' \equiv 0$ at the plasma surface, s_a. It can be shown from $\underline{B} = \nabla \times \underline{A}$ that the last term in eq(9) vanishes when $\tilde{B}_r'=0$. In the steady state, $\langle \tilde{\Phi}(a) \dot{\tilde{\psi}}(a) \rangle$ is small on s_a as seen in experiments. Thus, according to eq(9), the helicity transport $H(r=a)=0$ at the plasma boundary recovering the earlier result (equation 4).

In a steady state RFP, the sustaining loop voltage is continually injecting helicity into the inner region of the plasma but extracting from the outer region where the toroidal field is reversed. Transport is therefore required to balance this extraction and also the dissipation in this outer region. Such balance can be summarised by the following equations which is obtained by substituting $\underline{E}_0$ and $\underline{\tilde{E}}$ from eqs(1) and (2) into eq(10) and neglecting the small term $\int_V \eta \langle \underline{\tilde{j}} \cdot \underline{\tilde{B}} \rangle \, d^3x$:

$$V_\phi \Phi(r) = \int_V \eta \underline{j}_0 \cdot \underline{B}_0 \, d^3x + H(r) \tag{12a}$$

or, with eq (11),

$$V_\phi \Phi(r) = \int_V \eta \underline{j}_0 \cdot \underline{B}_0 \, d^3x - \int_V \langle \underline{\tilde{u}} \times \underline{\tilde{B}} \rangle \cdot \underline{B}_0 \, d^3x \qquad (12b)$$

In these forms, the steady state helicity balance equations display the injection of helicity by the sustaining loop voltage to balance the local resistive dissipation and to support the outward helicity flow. The generation of 'dynamo' and helicity transport is equivalent. In the outer region where $\underline{B}_0 \cdot \underline{\phi}$ is reversed, it is the flow represented by H which balances the local dissipation. It is important to realize that approximately no helicity is dissipated in the transport process as given by eq(5). This condition is independent of the profile of $\mu = \underline{j}_0 \cdot \underline{B}_0 / B^2$. In contrast, energy dissipation is μ profile dependent.

The global magnetic energy balance in a steady state RFP [22,12,7] is

$$V_\phi I_\phi(a) = \int_{V_a} \langle \underline{E} \cdot \underline{j} \rangle \, d^3x$$

and with fluctuation,

$$V_\phi I_\phi(a) = \int_{V_a} \underline{E}_0 \cdot \underline{j}_0 \, d^3x + \int_{V_a} \langle \underline{\tilde{E}} \cdot \underline{\tilde{j}} \rangle \, d^3x \qquad (13)$$

where $I_\phi(a)$ is the plasma toroidal current. Following similiar derivations as shown earlier and recongnising that

$$\int_{V_a} \underline{E}_0 \cdot \underline{j}_0 \, d^3x = \int_{V_a} \frac{V_\phi}{2\pi R} (\underline{j}_0 \cdot \underline{\phi}) d^3x = V_\phi I_\phi(a)$$

identically and $\int_{V_a} \eta\langle\tilde{j}^2\rangle\, d^3x = 0$, the analogy to eqns (4) and (5) for the energy balance are:

$$\int_{V_a} \langle\tilde{\underline{E}}\cdot\tilde{\underline{j}}\rangle\, d^3x = \int_{V_a} \eta\langle\tilde{j}^2\rangle\, d^3x - \int_{V_a} \langle\tilde{\underline{u}}\times\underline{B}_0\cdot\tilde{\underline{j}}\rangle\, d^3x \equiv 0 \qquad (14)$$

and

$$\int_{V_a} \langle\tilde{\underline{u}}\times\underline{B}_0\cdot\tilde{\underline{j}}\rangle\, d^3x = 0 \qquad (15)$$

With eqns (1) and (14) substituted into (13), the energy balance becomes

$$V_\phi I_\phi(a) = \int_{V_a} \eta j_0^2\, d^3x - \int_{V_a} \langle\tilde{\underline{u}}\times\tilde{\underline{B}}\rangle\cdot\underline{j}_0\, d^3x \qquad (16)$$

In the energy balance above, the integral which relates the energy throughout coupled to the fluctuations $\tilde{\underline{u}}\times\tilde{\underline{B}}$ is not necessarily zero and depends on the μ profile [7]. The power input from the sustaining loop voltage has to balance not only the resistive dissipation in heating the electrons but also the energy loss from the fluctuations. Recent studies (Carolan et al [23]) suggested that this loss of energy from the fluctuations is responsible for heating the ions.

As in the discussion on helicity transport, the energy balance equation for a sub-volume bounded by a mean flux surface inside the plasma needs to be derived first in order to reveal the role of fluctuations in transporting energy. The mean magnetic energy inside a mean flux surface is

$$\langle W \rangle = \int_V \langle B^2/2\mu_0 \rangle \, d^3x \tag{17}$$

The time rate change of the mean magnetic energy, after substituting $\underline{\dot{B}} = -\nabla \times \underline{E}$ is

$$\langle \dot{W} \rangle = -\int_V \langle \underline{E} \cdot \underline{j} \rangle \, d^3x - \frac{1}{\mu_0} \int_S \langle \underline{E} \times \underline{B} \rangle \cdot (d\underline{\theta} \times d\underline{\phi})$$

In terms of the mean and fluctuating parts and with eq(2)

$$\langle \dot{W} \rangle = V_\phi(r) I_\phi(r) - V_\theta(r) I_\theta(r) - \int_V \underline{E}_0 \cdot \underline{j}_0 d^3x - \int_V \langle \underline{\tilde{E}} \cdot \underline{\tilde{j}} \rangle \, d^3x$$

$$- \frac{1}{\mu_0} \int_S \langle \underline{\tilde{E}} \times \underline{\tilde{B}} \rangle \cdot (d\underline{\theta} \times d\underline{\phi}) \tag{18}$$

where $V_\theta(r) = \oint_S \underline{E}_0 \cdot d\underline{\theta}$, $V_\phi(r) = \oint \underline{E}_0 \cdot d\underline{\phi}$, $I_\phi(r) = \oint \underline{B}_0 \cdot d\underline{\theta}$ and $I_\theta(r) = \oint \underline{B}_0 \cdot d\underline{\phi}$. In steady state where $\langle \dot{W} \rangle = 0$, $V_\theta = 0$ and V_ϕ is a constant, the energy balance for a sub-volume becomes

$$V_\phi I_\phi(r) = \int_V \underline{E}_0 \cdot \underline{j}_0 d^3x + \frac{1}{\mu_0} \int_S \langle \underline{\tilde{E}} \times \underline{\tilde{B}} \rangle \cdot (d\underline{\theta} \times d\underline{\phi}) + \int_V \langle \underline{\tilde{E}} \cdot \underline{\tilde{j}} \rangle \, d^3x \tag{19}$$

With eq (2) and the identity, $\int_V \underline{E}_0 \cdot \underline{j}_0 \, d^3x = V_\phi I_\phi\ (r)$, substituting into eq (19), the sub-volume analogy to eq (14) is

$$\frac{1}{\mu_0}\int_S \langle \underline{\tilde{E}}\times\underline{\tilde{B}}\rangle\cdot(d\underline{\theta}\times d\underline{\phi}) = \int_V \langle \underline{\tilde{u}}\times\underline{B}_0\cdot\underline{\tilde{j}}\rangle\, d^3x - \int_V \eta\langle \tilde{j}^2\rangle\, d^3x \tag{20}$$

To complete the details of the energy transport, the momentum equation for the fluctuating $\underline{\tilde{u}}$ is dotted with $\underline{\tilde{u}}$, averaged and integrated over the volume to give

$$\int_V \rho\,\frac{\partial}{\partial t}\langle \tilde{u}^2/2\rangle\, d^3x = \int_V \langle \underline{\tilde{j}}\times\underline{B}_0\cdot\underline{\tilde{u}}\rangle\, d^3x + \int_V \langle \underline{j}_0\times\underline{\tilde{B}}\cdot\underline{\tilde{u}}\rangle\, d^3x$$

$$- \int_V \langle \nabla\tilde{p}\cdot\underline{\tilde{u}}\rangle\, d^3x - [\text{loss}] \tag{21}$$

where [loss] represents possible energy loss from $\underline{\tilde{u}}$ such as viscous damping. In a steady state

$$\int_V \rho\,\frac{\partial}{\partial t}\langle \tilde{u}^2/2\rangle\, d^3x = 0$$

After re-arranging and combining $\int_V \langle \tilde{p}\nabla.\underline{\tilde{u}}\rangle d^3x$ into [loss],

$$\int_V \langle \underline{\tilde{u}}\times\underline{B}_0\cdot\underline{\tilde{j}}\rangle\, d^3x = -\int_V \langle \underline{\tilde{u}}\times\underline{\tilde{B}}\rangle\cdot\underline{j}_0\, d^3x - [\text{loss}] - \int_S \langle \tilde{p}\underline{\tilde{u}}\rangle\cdot(d\underline{\theta}\times d\underline{\phi}) \tag{22}$$

With this, eq(20) becomes

$$- \int_V \langle \underline{\tilde{u}}\times\underline{\tilde{B}}\rangle\cdot\underline{j}_0\, d^3x = \int_V \eta\langle \tilde{j}^2\rangle\, d^3x + [\text{loss}] + U(r) \tag{23}$$

where the surface integrals which represents the transport of energy by fluctuations are denoted by

$$U(r) = \frac{1}{\mu_0} \int_s \langle \underline{\tilde{E}} x \underline{\tilde{B}} \rangle \cdot (d\underline{\theta} x d\underline{\phi}) + \int_s \langle \tilde{p} \underline{\tilde{u}} \rangle \cdot (d\underline{\theta} x d\underline{\phi}) \tag{24}$$

With eqs(1), (19) and (20) the steady state energy balance for a sub-volume can be written as

$$V_\phi I_\phi(r) = \int_V \eta j_0{}^2 \, d^3x - \int_V \langle \underline{\tilde{u}} x \underline{\tilde{B}} \rangle \cdot \underline{j}_0 \, d^3x \tag{25a}$$

or, when $\int_V \eta \langle \tilde{j}^2 \rangle \, d^3x$ is neglected and with eq (23),

$$V_\phi \; I_\phi(r) = \int_V \eta j_o{}^2 \, d^3x + U(r) + [\text{loss}] \tag{25b}$$

Similar to eqs(12a) and (12b), the above equations display the injection of energy by the sustaining loop voltage to balance the local dissipations and to support the outward energy flow carried by fluctuations. Although $\tilde{B}_r = 0$ on the plasma surface s_a, U(r) and [loss] are not necessarily zero and, according to eq (23),

$$- \int_{V_a} \langle \underline{\tilde{u}} x \underline{\tilde{B}} \rangle \cdot \underline{j}_0 \, d^3x = U(a) + [\text{loss}] \neq 0 \tag{26}$$

If, in addition to $\tilde{B}_r(a) = 0$, $\tilde{u}_r(a) = 0$ on s_a, then U(a)=0 and $\int_{V_a} \langle \underline{\tilde{u}} x \underline{\tilde{B}} \rangle \cdot \underline{j}_0 d^3x = -[\text{loss}]$. It can be seen from eq (5) that the integral in eq(26) is zero when the μ profile is uniform ie the fully relaxed state configuration:

$$\int_{V_a} \langle \underline{\tilde{u}} \times \underline{\tilde{B}} \rangle \cdot \underline{j}_0 \, d^3x = \frac{\mu}{\mu_0} \int_{V_a} \langle \underline{\tilde{u}} \times \underline{\tilde{B}} \rangle \cdot \underline{B}_0 \, d^3x = 0$$

Furthermore, eg(26) indicates that the loss of uniform μ profile is related to the non-resistive energy loss from the fluctuations in redistributing the helicity input. This concept departs from those which relate $\nabla\mu$ to helicity transport or dynamo generation. It should be noted that a sustained discharge cannot have a fully relaxed state if reversal sustainment activity is to vanish when $\nabla\mu=0$.

THE FULLY RELAXED MINIMUM ENERGY STATE

In the fully relaxed state, μ is uniform across the plasma radius and the integral in eq (16) which contain $\langle \underline{\tilde{u}} \times \underline{\tilde{B}} \rangle$ or eq(26) vanishes. This indicates either there is no transport of helicity and energy (ie $\langle \underline{\tilde{u}} \times \underline{\tilde{B}} \rangle \cdot \underline{B}_0 \equiv \langle \underline{\tilde{u}} \times \underline{\tilde{B}} \rangle \cdot \underline{j}_0 \equiv 0$) or there is no energy loss associated with the process in transporting helicity. In a steady state plasma, helicity input and dissipation must be balanced everywhere to maintain the field configuration stationary. It is obvious that the input is not uniform across the plasma radius of a steady state RFP. It is also unlikely that the dissipation is uniform. In fact, helicity input from the sustaining loop voltage is negative at the outer region where the toroidal field is reversed. Therefore transport of helicity must have taken place. Hence, in a sustained RFP, the only consistent situation for attaining the fully relaxed minimum energy state is when there is no energy loss associated with the fluctuations $\tilde{\underline{u}} \times \tilde{\underline{B}}$ in redistributing the helicity input. The μ profile observed in experiments [4,5,6] is not uniform but decreasing towards the plasma edge. Such departure can only be caused by energy losses, other than resistive dissipation, such as viscous damping or direct energy loss at the plasma edge.

It may appear that since plasma resistivity can affect the plasma current density, it would in turn affect the μ profile. Provided that the resistivity is not too high to render the concept of relaxation [3], the variation in the resistivity profile cannot change the μ profile if the plasma is to relax to the minimum energy state when there is no other constraint in addition to the global helicity invariant. Suppose there is an increase in the resistivity towards the plasma outer region. To maintain the same $\underline{j}_0$ as in the relaxed state, it is required that $\langle \tilde{\underline{u}} \times \tilde{\underline{B}} \rangle$ increases in this region. This is achieved by increasing the 'negative' $\langle \tilde{\underline{u}} \times \tilde{\underline{B}} \rangle$ in the inner region. The changes in $\langle \tilde{\underline{u}} \times \tilde{\underline{B}} \rangle$ in the two regions are related by eq (5). Such change will also extract more energy from the inner to the outer region. If there is no constraint to limit $\langle \tilde{\underline{u}} \times \tilde{\underline{B}} \rangle$, the relaxation will attain a uniform μ, in spite of any spatial variations in resistivity. On the other hand, if the μ profile at the outer region is to drop because of an increase in resistivity, then from eqs (5) and (16) more energy is extracted from the centre than that released to the outer region. This excess energy, if not lost, will increase the fluctuation activity to enhance $\langle \tilde{\underline{u}} \times \tilde{\underline{B}} \rangle$ until μ is restored to a uniform profile. When μ is uniform, there is no excess energy to increase the fluctuations further. The concept of relaxation or single global helicity invariant breaks down when there exists a limiting surface across which not sufficient helicity transport is permitted by the internal activity. In this case, the plasma cannot have complete relaxation with one global helicity invariant. The limiting surface divides the plasma into two regions which relax separately with their own global helicity invariant. As there is still finite helicity flow across the surface, the two regions are not totally independent. If there is no helicity transport, relaxation ceases.

HELICITY LEAKAGE

The global helicity balance equation has been shown to be unaffected by the fluctuations [7,11,12,13] in maintaining the field configuration. In the HBTX experiment [15], the resistivity calculated from the helicity balance is higher than the Spitzer resistivity based on the observed electron temperature and concentration of impurity. Such differences should be attributed to the leaking of helicity from the plasma edge region where magnetic field lines intercept a material surface. The helicity leakage when field lines come into contact with the vacuum vessel or objects such as the limiters, should be described by a surface integral which has hitherto been discarded. When the field lines do not close within the plasma volume but extend into an exterior volume v_b (eg a limiter), the global helicity balance equation is

$$V_\phi \Phi(a) = \int_{v_a} \langle \underline{E} \cdot \underline{B} \rangle d^3x - \int_{v_b} \nabla \chi_o \cdot \underline{B}_o \, d^3x$$

$$= \int_{v_a} \eta \underline{j}_o \cdot \underline{B}_o \, d^3x + \int_{s_a} \chi_o \underline{B}_o \cdot \underline{n} \, d^2x \qquad (27)$$

where $\nabla\chi_o = -\dot{\underline{A}}_o - \underline{E}_o$ represents an electrostatic potential and $\underline{n}$ the unit normal vector to the surface, s_a. The surface integral vanishes when the field lines are closed on themselves within v_a ie, $\underline{B}_o \cdot \underline{n} = 0$ on s_a.

When there is helicity leakage induced by the non-vanishing $\underline{B}_o \cdot \underline{n}$, eq (27) says that higher toroidal loop voltage is needed to sustain the plasma. For an obstruction such as a limiter with depth d inserted into the plasma and a projection area wd normal to the field lines, the surface integral can be written as

$$\int_{s_a} \chi_o \underline{B}_o \cdot \underline{n} \; d^2x = \Delta\chi B_o wd \tag{28}$$

where $\Delta\chi$ is the potential difference between the points of exit and entry of the field lines to s_a and B_o is taken to be uniform across the surfaces. This potential difference at the boundaries is generated by the obstruction of the current flow along the field lines (see later). A higher loop voltage is then required to balance not only the resistive dissipation of helicity within the plasma volume but also the leakage of helicity owing to the interception of field lines by the obstruction. The behaviour of the loop voltage now depends on two different processes. The first is resistive dissipation which depends on the electron temperature and the concentration of impurities. The second is the leakage loss, which is proportional to both the potential drop across the obstruction and the magnetic flux intercepted. According to eqs(27) and (28), this additional loop voltage is

$$\Delta V_\phi = \frac{V_\phi}{\Phi(a)} - \frac{1}{\Phi(a)} \int_{V_a} \eta \underline{j}_0 \cdot \underline{B}_0 \; d^3x$$

$$= \Delta\chi \; B_0 \; wd/\Phi(a) \tag{29}$$

To test such dependence, an experiment was conduced in HBTX. In this experiment [24], a graphite tile was inserted into the plasma. It was found that the loop voltage increases with the insertion depth as shown here in figure 1. Furthermore, when the tile is rotated by 90°, the increase in loop voltage is changed by a factor of 3 consistent with the change in the projection width. These results confirm the area dependence as given by eq (29).

In the HBTX experiment, it was found [25,26] that the toroidal loop voltage reduces as plasma is centred by an external vertical field. Such behaviour can be explained in terms of helicity loss through magnetic field lines intersecting the vacuum vessel. The magnetic flux crossing the wall is related to the plasma displacement (Δ) by

$$\int |\underline{B}_0 \cdot \underline{n}| \; d^2x = 8\pi R \Delta B_\theta(a)$$

where B is approximated by $B_\theta(a)$. Including such loss, the global helicity balance becomes

$$V_\phi \Phi(a) = \int_{V_a} \eta \underline{j}_0 \cdot \underline{B}_0 \; d^3x + \overline{\Delta\chi} \; 4\pi R \Delta B_\theta(a)$$

and the increase in the loop voltage is

$$\Delta V_\phi = \overline{\Delta\chi}\ (\frac{4R\theta}{a^2})\Delta \tag{30}$$

where $\theta = \pi a^2 B_\theta(a)/\Phi(a)$ is the pinch ratio and $\overline{\Delta\chi}$ is the averaged potential difference between the points of exit and entry of by the field lines. This is consistent with the observed dependence of V_ϕ on Δ which is reproduced here in figure 2. The experimental results gave value of $\overline{\Delta\chi}$ = 18 volts (an interpretation will be given later).

The leakage of helicity leads to increased fluctuation activity with an associated increase in helicity transport and affects the global energy balance because more energy throughout is coupled to the fluctuations. Therefore, the global effect on the energy balance is significant even though the local energy loss (see later) directly related to the helicity leakage may be small.

ENERGY LEAKAGE

The leakage of energy from the plasma when current flow is obstructed is described by a similar surface integral as in the helicity balance. With the surface integral retained, the global magnetic energy balance equation in a steady state RFP is

$$V_\phi I_\phi(a) = \int_{V_a} \langle \underline{E}\cdot\underline{j}\rangle d^3x - \int_{V_b} \nabla\chi_0\cdot\underline{j}_0 d^3x$$

$$= \int_{V_a} \eta j_o^2 \, d^3x - \int_{V_a} \langle \tilde{\underline{u}} \times \tilde{\underline{B}} \rangle \cdot \underline{j}_o \, d^3x + \int_{S_a} \chi_o \underline{j}_o \cdot \underline{n} \, d^2x \qquad (31)$$

In comparison to the surface integral in eq(27) for the helicity balance, the surface integral in eq (29) is insignificant for edge leakage except when the configuration is approaching the fully relaxed state. With a μ profile decreasing towards the plasma edge as observed in experiments, such edge energy leakage is negligible when compared to the total resistive dissipation.

There is, however, another form of energy leakage which can be related to the energy throughput coupled to the fluctuations. This can be seen in eqs (16) and (26). Since there are non-resistive energy loss when the plasma departs from the fully relaxed state, eq (26) says there must be radial $\tilde{\underline{u}}$ at the plasma surface leading to energy leakage if there is no internal non-resistive energy dissipation such as viscous damping ie [loss] = 0. Expressed in these leakage terms, the energy balance can be written as

$$V_\phi I_\phi(a) = \int_{V_a} \eta j_o^2 \, d^3x + \int_{V_a} \langle \tilde{p}\tilde{\underline{u}} \cdot \underline{n} \rangle \, d^2x + \int_{V_a} \chi_o \underline{j}_o \cdot \underline{n} \, d^2x + [\text{loss}] \qquad (32)$$

THE CURRENT SATURATION MODEL

When an object is immersed in a plasma, a sheath or boundary region appears between the material surface and the plasma. The formation of a sheath is initiated from the imbalance of the electron and ion flux to the surface. To maintain equal electron and ion fluxes to the surface, a sheath with a potential drop across it between the plasma and the material surface is developed to retard the higher influx specie. In general, when particle fluxes arise mainly from thermal motion, the electron flux is higher than ion flux resulting in a negative surface potential with respect to the plasma; this is called the floating potential. For Maxwellian electrons, the electron density inside the sheath varies exponentially as $n_e = n_o \exp(e\phi/kT_e)$ where n_o is the plasma density and ϕ is the potential with respect to the plasma potential outside the sheath. The size of the sheath is of the order of the Debye length. When there is drift in addition to the thermal motion, the sheath structure and size can be quite different. In a magnetised plasma, the influx of electron is mainly along the field lines whereas the influx of ions can be in all directions because of the difference in their Larmor radii. This effect increases the collecting surface area for ions.

When magnetic field lines intersect material objects the current flow along the field lines is obstructed if the surfaces of the obstacles do not emit charged particles, even though they may emit neutrals. Under this condition, the obstruction merely provides a surface for electrons and ions to recombine and cannot inject current into the plasma. Depending on the ratio of the plasma current density to the ion saturation current density ($j_s = en_oc_s$ where $c_s = (kT_e/m_i)^{1/2}$ is the ion sound speed), one of the two situations can exist to maintain the parallel plasma current flow:

1. When $j/j_s < 1$, a nett current flow through the obstacle can be

generated by the differences in the sheath potential drop and the electron flux to the two sides.

2. When $j/j_s > 1$, the plasma current has to flow around the obstruction (R S Pease, private communication). This can arise from electron flux along field lines neutralized by ion flux across field lines into the sheath.

If the current density is zero, the sheath structure and the potential drop on both sides of the obstruction is identical and there is no nett potential difference across the two sheaths. From one dimensional sheath analysis (see, for example, [27]), the floating potential (ϕ_f) which is the same on both sides is given by

$$j_s = \frac{en_o\bar{v}}{\pi^{1/2}} \exp\left(\frac{e\phi_f}{kT_e}\right) \tag{33}$$

where $\bar{v} = (2kT_e/m_e)^{1/2}$ is the electron thermal speed.

When $j/j_s < 1$, the formation of the sheath is not affected by the electron drift speed; the electron flux still arises mainly from thermal motion. In this case, the difference in the sheath potential drops ($\Delta\chi$) is related to the current density as in double probe theory (see, for example, [27]) ie

$$\frac{j}{j_s} = \tanh\left(\frac{e\Delta\chi}{kT_e}\right) \tag{34}$$

When $j/j_s > 1$, the formation of the sheath on the electron drift side is strongly affected. The electron flux entering the sheath now consists of two components; one from thermal motion as discussed above and another from the electron drift which gives the plasma current. The electron density inside the sheath does not drop as rapidly as $\exp(e\phi/kT_e)$ because part of the electron density which corresponds to the current carrying component increases as they slow down to maintain the same plasma current density. Inside the sheath, the difference between the ion and electron density which give rises to the non-neutrality is reduced. Consequently, the sheath becomes larger and more ions can be collected across the field lines. Once these ions enter the sheath, they are accelerated towards the surface by the sheath potential to neutralise the electron influx along the field lines. Thus, a parallel plasma current density higher than the ion saturation current density is made possible by increasing the ion collecting surface on the electron drift side of the obstruction.

Assuming that the electron current density in excess of the saturation current density will be balanced by the additional ion cross field influx as described above, an approximation for the potential drop can be obtained by equating the current through the obstruction to the saturation current ie

$$\frac{en\bar{v}}{\pi^{1/2}} \left[\exp\left(\frac{e\phi_f + e\Delta\chi/2}{kT_e}\right) - \exp\left(\frac{e\phi_f - e\Delta\chi/2}{kT_e}\right)\right] = 2\frac{en\bar{v}}{\pi^{1/2}} \exp\left(\frac{e\phi_f}{kT_e}\right)$$

$$\sinh\left(\frac{e\Delta\chi}{2kT_e}\right) = 1$$

$$\Delta\chi \simeq 1.8\ (kT_e/e) \qquad (35)$$

This means that the helicity loss caused by the non-vanishing $\underline{B}_o \cdot \underline{n}$ depends on the amount of flux intercepted and the local electron temperature. The value of $\Delta\chi$ obtained in HBTX experiment when field lines intersect the wall is 18 volts indicating that the edge electron temperature is 10 eV. For a small tile which intercepts a fixed amount of flux, the resultant ΔV_ϕ depends on the radial position of the tile in accordance to the radial distribution of kT_e. Applying this technique to HBTX [28], the radial ΔV_ϕ distribution suggests a broad temperature profile as shown in figure 3. The peak kT_e is about 50% higher than that measured by Si-Li detector with pulse height analysis suggesting that $\Delta\chi \sim 2.6\ kT_e/e$ may be a better approximation.

LOOP VOLTAGE ANOMALY

The loop voltage anomaly observed in RFP experiments can be explained in terms of dissipation and transport of helicity. The behaviour of the anomaly can be examined through eq (16) and (26). The first integral in eq(16) represents the classical resistive dissipation. The second integral or eq(26) which is the energy throughput coupled to fluctuations $<\tilde{\underline{u}} \times \tilde{\underline{B}}>$ in transporting helicity produces the voltage anomaly. Its behaviour depends on both the energy loss mechanism and the helicity transport. The loop voltage anomaly increases with helicity transport which can be enhanced by leakage or non-uniform dissipation. A limiter which intercepts the field lines and obstructs current flow leads to helicity leakage. An increase in resistivity near the edge of the plasma can cause higher helicity dissipation locally. In both cases, the local resistive dissipation of energy represents only a tiny fraction of the total resistive dissipation because of the small current flow near the edge. Nevertheless the loop voltage is increased anomalously by the 'amplified' energy loss due to the enhanced helicity transport in the bulk plasma.

CONCLUSIONS

The mean correlation between fluctuations in fluid velocity and in magnetic field, $\langle \tilde{\underline{u}} \times \tilde{\underline{B}} \rangle$, thought to furnish the emf needed to sustain current parallel to $\underline{B}_0$ has been investigated.

It has been shown that the integral $\int_V \langle \tilde{\underline{u}} \times \tilde{\underline{B}} \rangle \cdot \underline{B}_0 \, d^3x$ describes the rate of change of helicity in a volume v mainly by virtue of a flow through the bounding surface. No dissipation of helicity is involved, for when v represents the volume of the whole plasma the integral vanishes, implying that it contributes nothing to the rate of change of helicity in the plasma as a whole. This dissipation-free helicity transport, being independent of $\mu = \mu_0 j_0 / B_0$, persists whether or not the plasma is fully relaxed, and is identified with the generation of "dynamo" effect.

However, $\int_V \langle \tilde{\underline{u}} \times \tilde{\underline{B}} \rangle \cdot \underline{j}_0 \, d^3x$, the (non-resistive) energy loss associated with $\langle \tilde{\underline{u}} \times \tilde{\underline{B}} \rangle$, does not vanish even when the integral extends over the whole plasma volume, except when the plasma is fully relaxed, ie. $\nabla\mu=0$.

Together, these results show that $\nabla\mu$ is related to the non-resistive loss of energy by the fluctuations involved in helicity transport, but independent of the helicity transport itself, in contrast to models that seek to relate $\nabla\mu$ directly to the helicity transport (dynamo). By increasing these non-resistive losses, the plasma resistivity can increase the departure from the fully relaxed state although the resistive dissipation on its own will not.

Loss of magnetic flux through the bounding surface of the plasma enhances the helicity transport which increases the loop volts necessary to sustain the toroidal current. The increase is shown to be proportional to the flux intercepted and the local electron temperature. When the helicity leakage due to plasma wall contact is increased, as done on HBTX by controlling the plasma equilibrium shift using a vertical field or by the insertion of a graphite tile, a higher toroidal loop voltage is observed as expected from the theory. The associated extra energy input compensates the additional energy loss from the fluctuations in supporting increased helicity transport to balance the leakage. Even though only a small fraction of the total energy input is dissipated in the edge region, it can control the global plasma behaviour through enhanced helicity dissipation and transport in the plasma as a whole.

Two channels through which power flows into the plasma are identified, the first balancing resistive dissipation and accounting for electron heating, the second driving the fluctuations responsible for the helicity transport and perhaps the ion heating.

ACKNOWLEDGEMENTS

The author is grateful to M K Bevir, H A B Bodin, D E Evans, C G Gimblett and J B Taylor for their helpful discussions and D E Evans for invaluable comments and assistance with this manuscript.

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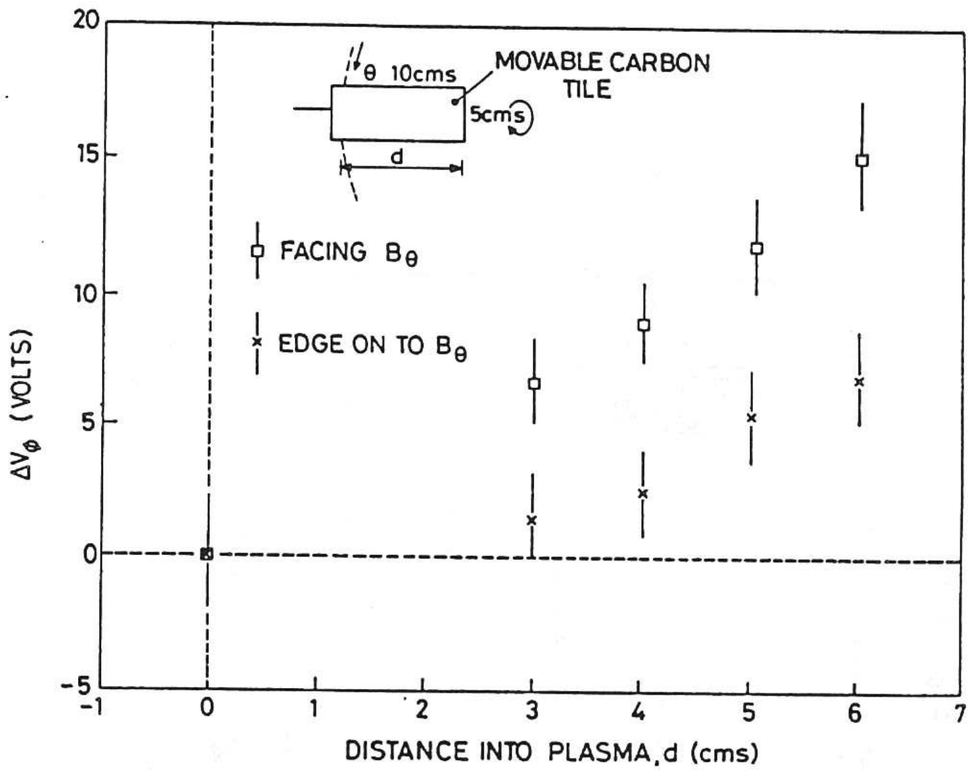

Figure 1 The increase in toroidal loop voltage (ΔV_ϕ) with the insertion distance (d) of a graphite tile. The ratio of the projection width of the tile for the two orientations, face on and edge on, is about 3.

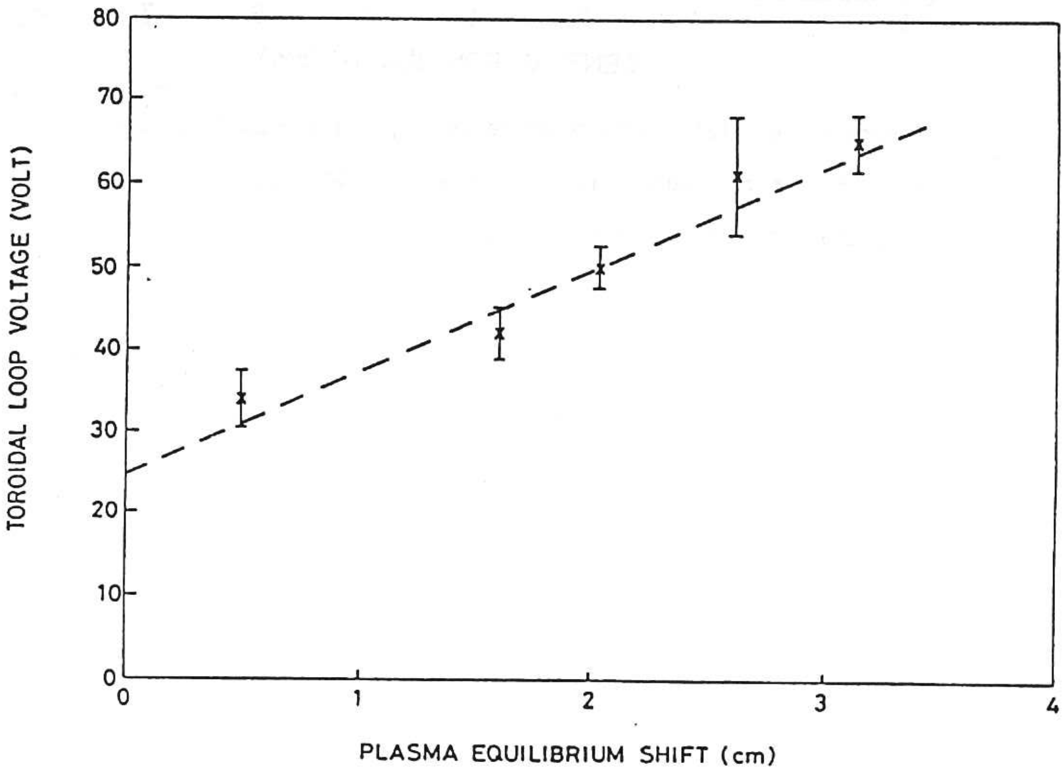

Figure 2 The variation of toroidal loop voltage (V_ϕ) with plasma equilibrium displacement (Δ). The dotted line is a linear fit to the data with $\Delta\chi$=18 volts.

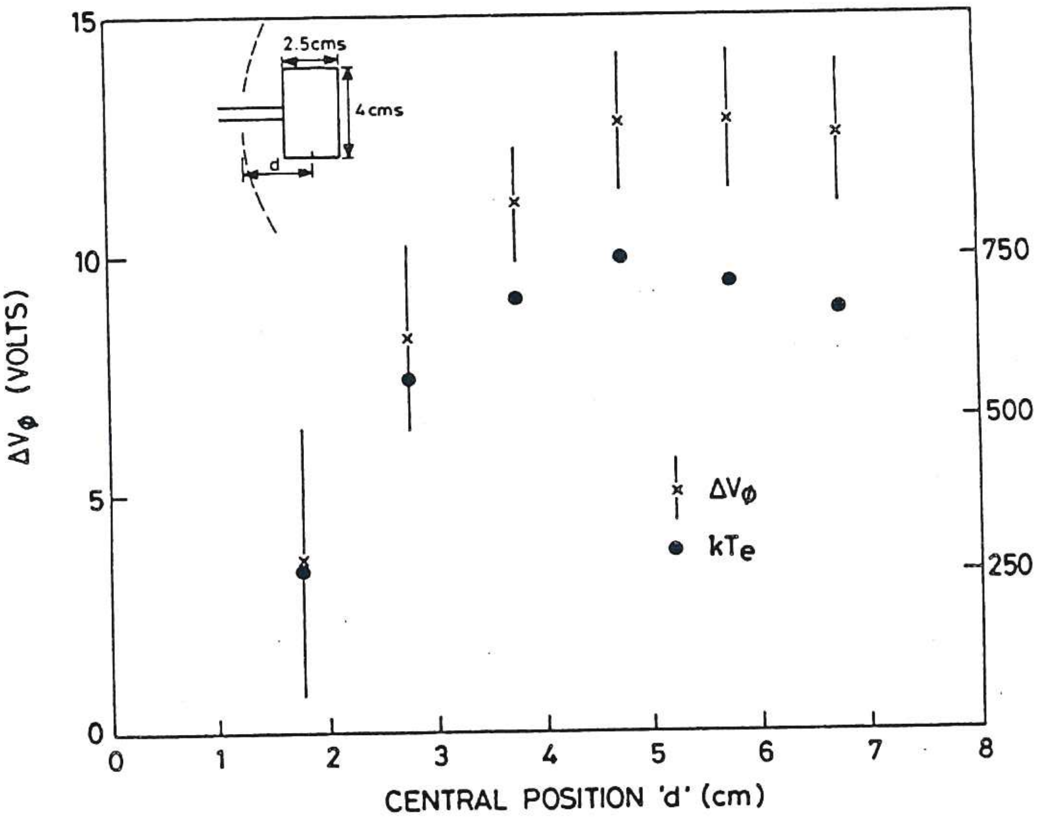

Figure 3 The electron temperature (kT_e) profile according to eqs (29) and (35) and the radial distribution of ΔV_ϕ obtained by the insertion of a small graphite tile with fixed area.