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Turbulent Recycling as a Mechanism for Ion Heating in the RFP Plasma

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ABSTRACT

Particle recycling resulting from turbulent convection is identified as a mechanism for ion heating in the RFP plasma. A steady state plasma in which an outward flux is balanced by the appearance of new material within the discharge is investigated through a perturbation analysis of moments of the Boltzmann equation modified by the addition of a source term. The work done in overcoming the drag exerted on the plasma turbulent motion by the new particles is shown to be the origin of non-Ohmic ion heating that can compete in effectiveness with conventional viscous heating. It is rapid enough to account for the anomalously fast ion heating seen in RFP plasmas and it can explain the observed correlation between ion temperature and loop volts. It accounts for the anomaly found in the resistivity derived from global magnetic energy balance. It is responsible for a new component of resistivity that could become dominant in low density discharges, and it is incompatible with a fully relaxed Taylor state, which could be approached only if recycling were to vanish.

Introduction

The reversed field pinch is a plasma confinement scheme based on a near minimum energy magnetic configuration [1] that depends on fluctuations to sustain it against resistive decay [2]. Ion heating has been found to be anomalously large and fast in the HBTX reversed field pinch at Culham [3], and this behaviour has been associated with fluctuation activity [4,5,6]. There is moreover evidence from fluorescent scattering [7] for the existence of a population of neutral deuterium extending far into the HBTX discharges which could provide a source of new plasma to balance the outward flow and loss to the walls implied by radial diffusion [8] as well as a mechanism for ion heating [9].

This paper explores the consequences of incorporating a source term into the fluctuation model of the RFP. As in previous work e.g. [10,11,12], the influence of fluctuating quantities in standard equations describing steady state conditions is investigated. The equations under consideration are the moments of the Boltzmann equation [13] into which a source term $S(\underline{x},t)\delta(\underline{v})$ is introduced to represent the number of particles appearing at position $\underline{x}$ and time t per unit volume and per unit time. The δ-function embodies the assumption that the particles are born at rest. By integration over velocity, the moment equations for continuity, momentum balance, and thermal energy balance are obtained. Attention is restricted to periodic cylindrical geometry to represent toroidal plasma in which mean variables are axisymmetric, viz. variation of the mean values occurs only in the radial direction. Variables are separated into mean and fluctuating parts and the average, denoted by $\langle ... \rangle$, is taken over a mean flux surface, and over a time interval long compared to a typical fluctuation period yet short relative to the pulse length. Thus for a quantity P,

$$\langle P \rangle \equiv \frac{1}{T}\int_t \left[\frac{1}{A}\int_s P \, d\sigma\right] dt,$$

where T is the appropriate time interval and A is the area of the flux surface. If $P = P_o + \tilde{P}$, the sum of a mean and a fluctuating part, then $\langle P \rangle = P_o$. SI units are used throughout.

The Moment Equations

With a source term $S(\underline{x},t)\delta(\underline{v})$ added to the Boltzmann equation for the particle distribution function, the moment equations in density $n(\underline{x},t)$ and fluid velocity $\underline{u}(\underline{x},t)$ are obtained in the usual way. The zeroth moment, the continuity equation, for either ions or electrons, is

$$\frac{\partial n}{\partial t} + \nabla\cdot(n\underline{u}) = S \qquad ...(1)$$

The first moment equation describes the momentum balance for ions,

$$n\, m_i \left(\frac{\partial}{\partial t} + \underline{u}_i\cdot\nabla\right)\underline{u}_i = en(\underline{E} + \underline{u}_i \times \underline{B}) - \nabla p_i - \nabla\cdot\pi_i + \underline{R}_{ie} - m_i\underline{u}_i S \qquad ...(2a)$$

and for electrons,

$$n\, m_e \left(\frac{\partial}{\partial t} + \underline{u}_e\cdot\nabla\right)\underline{u}_e = -en(\underline{E} + \underline{u}_e \times \underline{B}) - \nabla p_e + \underline{R}_{ei} - m_e\underline{u}_e S. \qquad ...(2b)$$

Here ions are treated as singly charged (Z=1), and viscous force, $\nabla\cdot\underline{\pi}_i$, is attributed only to the ions. $\underline{R}_{ie}$ and $\underline{R}_{ei}$ represent the usual collisional transfer of momentum between ions and electrons leading to Spitzer resistivity, $\eta_s = m_e/(e^2 n\tau_{ei})$ where τ_{ei} is the electron-ion collision time.

The novel terms involving source strength S, describing the rate of change of momentum due to the birth of new plasma, have the character of frictional drag, and give rise to an effective resistivity,

$$\eta_t \equiv \frac{m_e S}{e^2 n^2} = \frac{m_e}{e^2 n \tau_p} \qquad ...(3)$$

where $\tau_p \equiv n/S$ is the particle confinement time. When η_t is expressed in terms of τ_p, the analogy between η_s and η_t is apparent; electron momentum is dissipated through electron-ion collisions in the former and through particle loss in the latter.

The second moment equations describe thermal energy balance for ions,

$$n\left(\frac{\partial}{\partial t} + \underline{u}_i\cdot\nabla\right)\frac{3}{2}KT_i = Q_i - p_i\nabla\cdot\underline{u}_i + \tfrac{1}{2}m_i u_i^2 S - \frac{3}{2}KT_i S - \nabla\cdot\underline{h}_i - \pi_i\cdot\nabla\underline{u}_i \qquad ...(4a)$$

and for electrons,

$$n\left(\frac{\partial}{\partial t} + \underline{u}_e\cdot\nabla\right)\frac{3}{2}KT_e = Q_e - p_e\nabla\cdot\underline{u}_e + \tfrac{1}{2}m_e u_e^2 S - \frac{3}{2}KT_e S - \nabla\cdot\underline{h}_e \qquad ...(4b)$$

Temperatures $T_{i,e}$, heat flow $\underline{h}_{i,e}$ and Ohmic heating $Q_{i,e}$ have all been introduced in the usual way.

The new terms in these equations are $\frac{1}{2}m_{i,e}u_{i,e}^2\,S$ and $-\frac{3}{2}KT_{i,e}S$. The first is the rate of conversion of fluid kinetic energy into thermal energy due to the drag exerted by new particles, which in this way contrive actually to heat the plasma. The second is thermal energy loss associated with recycling.

Fluctuations and the Continuity Equations

Variables are now separated into mean and fluctuating parts and averages over time and mean flux surface are constructed. Steady state is imposed by setting $\frac{\partial}{\partial t}\langle\ldots\rangle = 0$.

Applying this operation first to the continuity equation, eq(1), requires that $n = n_o + \tilde{n}$, $u = u_o + \underline{\tilde{u}}$, and $\frac{\partial n_o}{\partial t} = 0$. For electrons

$$\partial\tilde{n}/\partial t + \nabla\cdot(n_o\underline{u}_{eo} + \tilde{n}\underline{u}_{eo} + n_o\underline{\tilde{u}}_e + \tilde{n}\underline{\tilde{u}}_e) = S, \qquad \ldots(5)$$

while for ions, adopting the arbitrary convention, without loss of generality, that the mean ion velocity vanishes, viz $\underline{u}_{io} = 0$,

$$\partial\tilde{n}/\partial t + \nabla\cdot(n_o\underline{\tilde{u}}_i + \tilde{n}\,\underline{\tilde{u}}_i) = S. \qquad \ldots(6)$$

Averaging and ignoring terms of higher order results in

$$\nabla\cdot(n_o\underline{u}_{eo}) + \nabla\cdot\langle\tilde{n}\,\underline{\tilde{u}}_e\rangle = S \text{ for electrons}$$

and

$$\nabla\cdot\langle\tilde{n}\,\underline{\tilde{u}}_i\rangle = S \text{ for ions.}$$

By virtue of the averaging process, mean quantities vary only in a direction normal to a mean flux surface i.e. radially. Moreover, mean electron velocity i.e. mean plasma current is assumed to lie in a mean flux surface and possess no radial component, i.e. $\underline{u}_{eo}\cdot\nabla r \equiv 0$. Thus

$$\underline{u}_{eo}\cdot\nabla n_o = 0 \quad\text{and}\quad \nabla\cdot\underline{u}_{eo} = 0. \qquad \ldots(7)$$

Then, since the source strength S has been assumed to be the same for both ions and electrons, it follows from quasi-neutrality of the plasma that

$$\langle \tilde{n}\ \underline{\tilde{u}}_e \rangle = \langle \tilde{n}\ \underline{\tilde{u}}_i \rangle \equiv \underline{\Gamma}, \qquad \ldots(8)$$

and for either electrons or ions

$$\nabla \cdot \underline{\Gamma} = S. \qquad \ldots(9)$$

Multiplying eq(6) by $\tilde{n}$, averaging ignoring higher order terms and assuming steady state disposes of all terms but

$$n_o \langle \tilde{n}\ \nabla . \underline{\tilde{u}} \rangle + \underline{\Gamma} . \nabla n_o = 0. \qquad \ldots(10)$$

The subscript on $\underline{\tilde{u}}$ is omitted in eq(10) because the same procedure applied to the equation of electron continuity, eq(5), together with the use of eq(8) and the fact that $\underline{u}_{eo}$ has no radial component while gradients of mean quantities like n_o and $\langle \tilde{n}^2 \rangle$ are exclusively radial, produces an equation identical to eq(10) for the electrons as well as the ions.

A flux of fluid can exist even though the mean velocity is zero, and it can be seen that $\langle \tilde{n}\ \underline{\tilde{u}} \rangle$, the correlation between density fluctuations and fluid velocity fluctuations, describes such a flux. Eq(10) shows that the non-vanishing of $\langle \tilde{n}\ \underline{\tilde{u}} \rangle$ relies on the compressibility of the fluid since $\langle \tilde{n}\ \underline{\tilde{u}} \rangle = 0$ if $\nabla \cdot \underline{\tilde{u}} = 0$.

<u>Fluctuations and the Momentum Balance Equations</u>

Turning to the momentum balance, we form the dot product of u_i with eq(2a), of $\underline{u}_e$ with eq(2b), and add to obtain the local instantaneous non-ohmic rate of power dissipation, namely

$$\begin{aligned} \underline{E} \cdot \underline{j} - \eta\ j^2 = {} & n \frac{\partial}{\partial t} (\tfrac{1}{2} m_i u_i^2 + \tfrac{1}{2} m_e u_e^2) \\ & + n[\underline{u}_i \cdot \nabla (\tfrac{1}{2} m_i u_i^2) + \underline{u}_e \cdot \nabla (\tfrac{1}{2} m_e u_e^2)] \\ & + \underline{u}_i \cdot \nabla p_i + \underline{u}_e \cdot \nabla p_e + \underline{u}_i \cdot \{\nabla \cdot \pi_i\} \\ & + 2\ (\tfrac{1}{2} m_i u_i^2 + \tfrac{1}{2} m_e u_e^2)\ S. \end{aligned} \qquad \ldots(11)$$

The last term, involving S, is the rate of non-ohmic power absorption due to the source of new plasma. It has been written as the sum of two halves to emphasize that one half furnishes the kinetic energy absorbed by the newly born particles and the second half is the work done to overcome the drag exerted by the newly born zero momentum particles on the fluid.

An analogy has been suggested [14] between this process and a block dropped onto a moving belt. The block is brought up to the speed of the belt through friction, which generates heat due to the initial sliding between the block and the belt. The belt drive does more work than merely to accelerate the block: it also furnishes the energy that appears as the heat produced by friction. Consider the process from the frame of reference in which the belt is stationary. Then the block of mass M projected along the stationary belt with initial velocity V is brought to rest by friction when its initial kinetic energy $\frac{1}{2}MV^2$ has all been converted into heat. In the frame of reference in which the block is initially at rest and the belt moves with velocity V, friction is seen to accelerate the block until its velocity is the same as that of the belt, at which point it has acquired kinetic energy $\frac{1}{2}MV^2$. But the amount of frictional heating, being independent of the frame of reference from which the process is viewed, remains $\frac{1}{2}MV^2$. Thus the work done in accelerating a test particle up to any particular directed velocity is precisely twice the energy associated with that directed motion, as equations (4) and (11) already indicate. This is the physical process by which the appearance of new cold material channels energy into plasma heating. The same argument applies even if the new material has finite temperature, provided only that the particle velocities are truly random and average to zero.

The mean local rate of non-ohmic power dissipation reveals the significance of fluctuations to this process. Once again, variables n, $\underline{E}$, $\underline{j}$, $\underline{u}_i$, $\underline{u}_e$, p_i and p_e are separated into their mean and fluctuating parts and averages taken, and the assumption of steady state implies $\frac{\partial}{\partial t}\langle\ldots\rangle = 0$. The outcome is

$$
\begin{aligned}
\langle \underline{E}\cdot\underline{j} - \eta j^2\rangle = {} & \nabla\cdot[\langle \tilde{p}_i \underline{\tilde{u}}_e\rangle + \langle \tilde{p}_e \underline{u}_e\rangle \\
& + (\langle \tfrac{1}{2}m_i \tilde{u}_i^2\rangle + \langle \tfrac{1}{2}m_e \tilde{u}_e^2\rangle + \tfrac{1}{2}m_e u_{eo}^2)\underline{\Gamma}] \\
& - \langle \tilde{p}_i \nabla\cdot\underline{\tilde{u}}_i\rangle - \langle \tilde{p}_e \nabla\cdot\underline{\tilde{u}}_e\rangle \\
& + (\langle \tfrac{1}{2}m_i \tilde{u}_i^2\rangle + \langle \tfrac{1}{2}m_e \tilde{u}_e^2\rangle + \tfrac{1}{2}m_e u_{eo}^2)\, S \\
& + \langle \underline{\tilde{u}}_i\cdot(\nabla\cdot\pi_i)\rangle \qquad \ldots(12)
\end{aligned}
$$

The term involving viscosity can be written identically as

$$\langle\tilde{\underline{u}}_i\cdot(\nabla\cdot\pi_i)\rangle \equiv \nabla\cdot\langle\tilde{\underline{u}}_i\cdot\pi_i\rangle - \langle\pi_i\cdot(\nabla\tilde{\underline{u}}_i)\rangle \equiv V_2 - V_1.$$

The term on the left hand side is the total work done by the fluid in overcoming viscosity. This is equal to V_2 the rate at which this energy flows away, plus $-V_1$, the rate at which viscosity leads to local heating. The latter will appear subsequently in the ion thermal energy balance (cf. eq(15) et seq). In what follows, we choose to ignore the rate of energy flow due to viscosity and so set $V_2= 0$.

It is natural to enquire to what extent it is correct to identify $\langle\underline{E}\cdot\underline{j} - \eta j^2\rangle$ with the work being done by the dynamo, that is $-\langle\tilde{\underline{u}}x\tilde{\underline{B}}\rangle\cdot\underline{j}_o$. Using Poynting's theorem, the latter can be shown to be

$$-\langle\tilde{\underline{u}}x\tilde{\underline{B}}\rangle\cdot\underline{j}_o = \langle\underline{E}\cdot\underline{j}-\eta j^2\rangle + \nabla\cdot\langle\tilde{\underline{E}}x\tilde{\underline{B}}\rangle/\mu_o + \eta\langle\tilde{j}^2\rangle$$

$$\simeq \langle\underline{E}\cdot\underline{j}-\eta j^2\rangle + \nabla\cdot\langle\tilde{\underline{E}}x\tilde{\underline{B}}\rangle/\mu_o \qquad \ldots(13)$$

since $\eta\langle\tilde{j}^2\rangle$ is the smallest of the terms on the right hand side by at least one order of magnitude. Thus $\langle\underline{E}\cdot\underline{j}-\eta j^2\rangle$ constitutes only part of the total dynamo power, which eq(13) shows is made up of this local non-ohmic dissipation, available for heating, in addition to an energy flux term $\nabla\cdot\langle\tilde{\underline{E}}x\tilde{\underline{B}}\rangle/\mu_o$, recognizable as the transport of dynamo power in the form of Poynting's vector.

The right hand side of eq(12) sums all the mechanical processes by which the dynamo energy is absorbed locally. The divergence terms account for inflow and outflow of energy through turbulent convection, and those terms explicitly involving S are the drag force heating discussed above. Noting that pressure $\tilde{p} = \tilde{n}$ KT and using eq(10), it can be seen that

$$-[\langle\tilde{p}_i\nabla\cdot\tilde{\underline{u}}_i\rangle + \langle\tilde{p}_e\nabla\cdot\tilde{\underline{u}}_e\rangle] = K(T_i + T_e)\underline{\Gamma}\cdot\nabla n_o/n_o \qquad \ldots(14)$$

corresponding to a thermal expansion/compression of outflowing plasma. We believe this exchange of energy between heat and fluid motion is a hitherto unrecognized component of the dynamo.

The dynamo furnishes non-ohmic heating to ions and electrons alike. While the electrons acquire energy ohmically as well, ions have no other source of heating but the dynamo since equipartition heating of ions by electrons is slow by comparison with the particle confinement time typical of RFP discharges. Scrutiny reveals that if S were to vanish, so would every term on the right hand side of eq(12), except that involving the viscosity. This implies that particle recycling, embodied in S can be an important contribution to ion heating.

We note that this non-Ohmic ion heating accounts for the well known difference between the resistivities calculated on the basis of magnetic energy balance and magnetic helicity balance [15] and that the difference is related to the departure from the fully relaxed Taylor state [4].

Integrating eq(13) over the whole plasma volume V_a produces

$$-\int_{V_a} \langle \tilde{\underline{u}} x \tilde{\underline{B}} \rangle \cdot \underline{j}_o \, d^3x = \int_{V_a} \langle \underline{E} \cdot \underline{j} - \eta j^2 \rangle \, d^3x$$

since

$$\int_{V_a} \nabla \cdot \langle \tilde{\underline{E}} x \tilde{\underline{B}} \rangle / \mu_o \, d^3x \equiv - \int_{V_a} \langle \tilde{\underline{E}} \cdot \tilde{\underline{j}} \rangle \, d^3x \equiv 0 \; [16]$$

Moreover, since $\mu_o \underline{j}_o = \mu \underline{B}_o$, where μ comes from the RFP equation, $\nabla \times \underline{B} = \mu \, \underline{B}$,

$$-\int_{V_a} \langle \tilde{\underline{u}} x \tilde{\underline{B}} \rangle \cdot \underline{j}_o \, d^3x = -\int_{V_a} \langle \tilde{\underline{u}} x \tilde{\underline{B}} \rangle \cdot \mu \underline{B}_o / \mu_o \, d^3x.$$

If μ were independent of position, i.e. if the plasma were in a fully relaxed Taylor state, μ could be taken outside the integral which would then vanish [17]. Thus if,

$$\int_{V_a} \langle \underline{E} \cdot \underline{j} - \eta j^2 \rangle \, d^3x \neq 0$$

then μ cannot be independent of position and the plasma cannot be in a fully relaxed state. The implication of the foregoing argument is that dynamo non-Ohmic heating is incompatible with the fully relaxed state in a reversed field pinch plasma.

Fluctuations and the Equations of Thermal Energy Balance

We turn finally to the equations of thermal energy balance, eq(4), and ignore heat conduction by setting $\underline{h}_e = \underline{h}_i = 0$. Setting $Q_i = 0$ and $Q_e = \eta j^2$, ignoring radiation losses, recognizing $p = nKT$, introducing fluctuating density and fluid velocity (but not temperature), setting $\frac{\partial}{\partial t}(KT) = 0$ for steady state, and averaging, results in, for the ions

$$\langle \tilde{n} \tilde{\underline{u}} \rangle \cdot \nabla (^3/_2 \, KT_i) = - \langle \tilde{p}_i \nabla \cdot \tilde{\underline{u}}_i \rangle + \tfrac{1}{2} m_i \langle \tilde{u}_i^2 \rangle \, S - {}^3/_2 \, KT_i S - V_1. \quad \ldots(15a)$$

and for electrons, making use of eq(7) and the facts that the mean electron velocity $\underline{u}_{eo}$ is perpendicular to the gradient of temperature

∇KT_e,

$$\langle \tilde{n}\tilde{\underline{u}}\rangle\cdot\nabla(^3/_2\ KT_e) = \langle \eta j^2\rangle - \langle \tilde{P}_e\nabla\cdot\tilde{\underline{u}}_e\rangle + \tfrac{1}{2}m_e u_{eo}^2\ S + \tfrac{1}{2}m_e\langle\tilde{u}_e^2\rangle S - {}^3/_2\ KT_e S.$$

...(15b)

With eq(3) and eq(10), the foregoing can be written, for ions:

$$(KT_i)\underline{\Gamma}\cdot[^3/_2\ \nabla(KT_i)/(KT_i) - \nabla n_o/n_o] = \tfrac{1}{2}\ m_i\langle\tilde{u}_i^2\rangle\ S - {}^3/_2\ KT_i S - V_1.$$

and for electrons

$$(KT_e)\underline{\Gamma}\cdot\ [^3/_2\ \nabla(KT_e)/(KT_e) - \nabla n_o/n_o] = (\eta_s + \tfrac{1}{2}\eta_t)\ \langle j^2\rangle - {}^3/_2\ KT_e S.$$

Rewriting the above in cylindrical polar coordinates, introducing thermal velocities through $KT_{e,i} = \tfrac{1}{2}mv_{e,i}^2$, and approximating $\langle j^2\rangle$ by $e^2n^2u_{eo}^2$ leads, for ions, to

$$\Gamma(^3/_2\ \frac{1}{T_i}\frac{\partial T_i}{\partial r} - \frac{1}{n_o}\frac{\partial n_o}{\partial r}) = {}^3/_2\ S\ (^2/_3\ \langle\tilde{u}_i^2\rangle/v_i^2 - 1) - \frac{V_1}{KT_i}$$

$$\Gamma\ (\Lambda_{T_i} - {}^3/_2\ \Lambda_n)/(\Lambda_{T_e}\Lambda_n) = {}^3/_2\ S(^2/_3\ \langle\tilde{u}_i^2\rangle/v_i^2 - 1) - \frac{V_1}{KT_i}$$

...(16a)

and for electrons

$$\Gamma(^3/_2\ \frac{1}{T_e}\frac{\partial T_e}{\partial r} - \frac{1}{n_o}\frac{\partial n_o}{\partial r}) = {}^3/_2\ S(^2/_3(2\eta_s/\eta_t + 1)\ \frac{u^2{}_{eo}}{v^2{}_e} - 1)$$

$$\Gamma(\Lambda_{T_e} - {}^3/_2\ \Lambda_n)/(\Lambda_{T_e}\ \Lambda_n) = {}^3/_2\ S(^2/_3\ (2\eta_s/\eta_t + 1)\ \frac{u_{eo}^2}{v^2{}_e} - 1)$$

...(16b)

where terms like $\frac{1}{T}\frac{\partial T}{\partial r}$ have been written as $-1/\Lambda_T$, Λ_T being a measure of the width of the T profile, the negative sign reflecting the fact that in general $T_e(r)$, $T_i(r)$, and $n(r)$ have $\frac{\partial}{\partial r} < 0$.

Figure 1 illustrates schematically the relative profiles believed to exist in a standard HBTX RFP plasma, and it can be seen that their relative widths are ordered like

$$\Lambda_{T_i} < \Lambda_n < \Lambda_{T_e}$$

Thus, away from the axis or the wall, the left hand side of eq(16a) is negative, implying that

$$\langle\tilde{u}_i^2\rangle < {}^3/_2\, v_i^2 + 2V_1/m_i S.$$

Similarly the left hand side of eq(16b) is positive, from which it follows that

$$(\tfrac{1}{2}\eta_t + \eta_s)j^2 > {}^3/_2\, KT_e S.$$

On axis, where, it will be recalled, S≥0 and Γ=0, the left hand side of both eq(16a) and (16b) vanish, hence

$$\tfrac{1}{2}m_i\langle\tilde{u}_i^2\rangle S - V_1 = {}^3/_2\, KT_i(0)\, S \qquad \ldots(17)$$

and

$$(\tfrac{1}{2}\eta_t + \eta_s)j^2 = {}^3/_2\, KT_e(0)\, S \qquad \ldots(18)$$

Eq(17) expresses the close relation between ion temperature and fluid velocity fluctuations. Ignoring viscous effects the ion thermal velocity and the fluid fluctuation velocity have the same rms value.

On axis, Γ=0 and eq(12) becomes

$$\langle \underline{E}\cdot\underline{j} - \eta j^2\rangle = [KT_e(0) + KT_i(0)]S + (m_i\langle\tilde{u}_i^2\rangle + m_e\langle\tilde{u}_e^2\rangle + m_e\dot{u}_{eo}^2)S - V_1 \qquad \ldots(19)$$

and all quantities take their on-axis values. Since $\langle\tilde{u}_e^2\rangle < u_{eo}^2$ from $\langle\tilde{j}^2\rangle < j_o^2$, eqn (19) becomes,

$$\langle \underline{E}\cdot\underline{j} - \eta j^2\rangle = [KT_e(0) + KT_i(0)]S + (m_i\langle\tilde{u}_i{}^2\rangle + m_e\, u^2{}_{eo})S - V_1.$$

Substituting eq(17) for $m_i\langle\tilde{u}_i^2\rangle S$, eq(18) for $KT_e(0)S$, and $j_o^2/e^2 n_o^2$ for u_{eo}^2, and ignoring η_t as negligible compared to the Spitzer resistivity η_s, we obtain

$$4\,KT_i(0)\;S = \underline{E}(0)\cdot\underline{j}(0) - {}^5/_3\,\eta_s\,j^2(0)$$

Combining the last equation with eq(18) leads to an expression for the temperature ratio on axis:

$$\frac{T_i(0)}{T_e(0)} = \frac{3}{8}\left(\frac{E_o}{\eta_s j_o} - \frac{5}{3}\right) \qquad \ldots(20)$$

which compares favourably with experimental data from HBTX [18] as shown in Figure 2. Here, to maintain constant toroidal current I_ϕ, loop volts V_ϕ were increased as a graphite tile was inserted into the edge of the plasma, intercepting flux [19]. Reversal F and pinch Θ remained approximately constant during this process, and $\eta_s j_o$ was held fixed (at $1.85\mathrm{Vm}^{-1}$) accordingly.

The total energy balance, thermal plus fluid, can now be written by adding equations (12), (15a) and (15b):

$$\langle\underline{E}\cdot\underline{j}\rangle = \nabla\cdot\{{}^5/_2\,K(T_e + T_i)\underline{\Gamma} + (\tfrac{1}{2}\,m_i\langle\tilde{u}_i{}^2\rangle + \tfrac{1}{2}\,m_e\langle\tilde{u}_e{}^2\rangle)\,\underline{\Gamma} + \tfrac{1}{2}\,m_e u_{eo}{}^2\underline{\Gamma}$$

$$+ \langle\tilde{u}_i\cdot\pi_i\rangle\} \qquad \ldots(21)$$

This says, in steady state the mean rate at which energy is being put in at a point equals the flow of energy away from the point. Of the terms within the divergence brackets, the last is due to viscosity, the middle terms describe the kinetic energy associated with the mean electron current and with the electron and ion mean fluctuation velocities, while the first term consists of ${}^3/_2$ KT, the thermal energy content, plus ${}^2/_2$ KT compressional energy from terms like $\nabla\cdot\langle\tilde{p}\;\tilde{\underline{u}}\rangle$.

Discussion

The aim of this paper has been to investigate the role of particle recycling in turbulent plasma behaviour, and its influence, especially upon ion heating in the RFP, has been clearly exhibited. Viscosity terms have been carried through the calculation, but as they do not appear to couple to the recycling, their effect has thus far been ignored.

We now enquire into the relative importance of recycling and viscosity for ion heating. Their relative contribution on axis can be investigated through eq (17):

$$Q_{recyc} + Q_{visc} = {}^3/_2\,KT_iS \qquad \ldots(22)$$

where $Q_{recyc} = \tfrac{1}{2}m_i\langle\tilde{u}_i^2\rangle S$ and $Q_{visc} \equiv -V_1$. To describe Q_{visc}, the rate of viscous ion heating, we use parallel ion viscosity which is the dominant contribution for high temperature magnetised plasma with approximately equal electron and ion temperatures [20] viz

$$Q_{visc} = \mu_{\parallel} \left|\nabla\cdot\underline{u}_i\right|^2 . \qquad ...(23)$$

In the collisional regime, the form for $\mu_{\parallel}$ is [21]

$$\mu_{\parallel}^{c} = nKT_i\tau_{ii}$$

τ_{ii} being the ion-ion collisional time. On-axis, the continuity equation for the fluctuation components gives

$$\left|\nabla.\tilde{\underline{u}}_i\right|^2 = \omega^2\left|\tilde{n}/n_o\right|^2 \qquad ...(24)$$

where ω is the angular frequency of the fluctuation. In a steady state, the total power input balancing the power loss and in general

$$Q_{total} = Q_{loss} = {}^3/_2\ KT_i/\tau_{Ei}$$

where τ_{Ei} is the ion energy confinement time. The fractional contribution from viscous heating is

$$\frac{Q_{visc}}{Q_{loss}} = \frac{nKT_i\tau_{ii}\omega^2\left|\tilde{n}/n_o\right|^2}{{}^3/_2\ KT_i/\tau_{Ei}}$$

$$= {}^3/_2\ (\tau_{Ei}/\tau_p)(\omega\tau_{ii})(\omega\tau_p)\left|\tilde{n}/n_o\right|^2 \qquad ...(25)$$

with τ_p being the particle confinement time. The collisional condition requires $\omega\tau_{ii} < 1$ and $\omega\tau_p < 1$. When energy loss arises from convection, as is the case being investigated, $\tau_{Ei} \sim \tau_p$. The ratio Q_{visc}/Q_{loss} is accordingly always less than unity showing that heating rate arising from viscous damping of the turbulent velocity is ineffective to compete against the loss rate arising from the turbulent convection. On the other hand, the fractional contribution from recycling,

$$\frac{Q_{recyc}}{Q_{loss}} = {}^2/_3\ \left|\tilde{u}_i/v_i\right|^2 \qquad ...(26)$$

can account for the power balance with ion temperature being determined by the power in the fluctuating velocity.

To extend the analysis into the collisionless regime, we need an appropriate coefficient of viscosity. The collisional form of $\mu_{\|}^{c}$ increases without limit as collisions become less frequent, which is patently inphysical. This can be corrected by dividing $\mu_{\|}^{c}$ by the square of the Knudsen number for the system [22]:

$$K_n = \lambda_{mfp}/2\pi R = v_i \tau_{ii}/2\pi R$$

giving a viscosity coefficient in the collisionless regime,

$$\mu_{\|}^{k} = \mu_{\|}^{c}(2\pi R/\lambda_{mfp})^2$$

where R the major radius of the toroidal confinement assembly.

An alternate approach to avoid an unbounded increase in μ and one that appears to be appropriate for the description of fluctuation phenomena, has been devised by Rusbridge [23]. It consists of averaging the viscous heating rate over the distribution of collision times, thus

$$Q_{visc} = nKT_i \ (\nabla\cdot\tilde{\underline{u}}_i)_o \int_o^\infty (\nabla\cdot\tilde{\underline{u}}_i) \exp(-t/\tau_{ii})\, dt$$

where the $(\nabla\cdot\tilde{u}_i)$ under the integral is evaluated at time t. In terms of Fourier modes, $\nabla\cdot\tilde{\underline{u}}_i = (\nabla\cdot\tilde{\underline{u}}_i)_o \cos\omega t$, and one obtains

$$Q_{visc} = nKT_i \left|\nabla\cdot\tilde{\underline{u}}_i\right|^2 \tau_{ii}/(1 + (\omega\tau_{ii})^2).$$

In this case, the coefficient of viscosity becomes

$$\mu_{\|}^{R} \equiv \mu_{\|}^{c} \ /(1 + (\omega\tau_{ii})^2).$$

In all cases, the ratio of viscous to recycling heating is

$$\frac{Q_{visc}}{Q_{recyc}} = \frac{-V_1}{\frac{1}{2}m_i\langle\tilde{u}_i^2\rangle S} = \frac{\mu_{\|}\left|\tilde{u}_i\right|^2}{\frac{1}{2}m_i\left|\tilde{u}_i^2\right| S\Lambda_u^2} \cong \frac{\mu_{\|}}{\frac{1}{2}m_i\Lambda_u^2 S}$$

where Λ_u is the scale length associated $\nabla\cdot\tilde{\underline{u}}_i$. Using the mean field Ohms Law, $\underline{E}_o = \eta \underline{j}_o - \langle\tilde{\underline{u}}x\tilde{\underline{B}}\rangle$ and eq(24), we obtain the on-axis value of Λ_u,

$$\Lambda_u \sim \frac{\left|\underline{E}_o - \eta\underline{j}_o\right|}{\omega g\left|\tilde{n}/n_o\right|\left|\tilde{B}/B_o\right|B_o}.$$

Here, all the quantities take their on-axis values, g is the correlation coefficient associated with the dynamo fluctuations, i.e.

$$g \equiv \frac{\langle \tilde{u} x \tilde{B} \rangle}{|\tilde{u}||\tilde{B}|}.$$

Accordingly, for all cases,

$$\frac{Q_{visc}}{Q_{recyc}} = \frac{\mu_{\parallel}}{nKT_i\tau_{ii}} \frac{g^2(\omega\tau_{ii})^2}{H^2} \qquad \ldots(27)$$

with $H = \sqrt{(\tau_p/\tau_{ii})} \left|\underline{E}_o - \eta \underline{j}_o\right| / (v_i \left|\tilde{n}/n_o\right| \left|\tilde{B}/B_o\right| B_o)$.

In the collisional regime where $\mu_{\parallel}^c$ is appropriate, the foregoing becomes

$$\frac{Q_{visc}}{Q_{recyc}} = \frac{g^2}{H^2} (\omega\tau_{ii})^2. \qquad \ldots(28a)$$

With the Knudsen number correction, the result is

$$\frac{Q_{visc}}{Q_{recyc}} = \frac{g^2}{H^2} (\omega\tau_{ii})^2 (2\pi R/\lambda_{mfp})^2, \qquad \ldots(28b)$$

while the Rusbridge form for fluctuations leads to

$$\frac{Q_{visc}}{Q_{recyc}} = \frac{g^2}{H^2} \frac{(\omega\tau_{ii})^2}{1 + (\omega\tau_{ii})^2}. \qquad \ldots(28c)$$

Typical on-axis HBTX parameters, e.g. T_i = 100 eV, v_i = 1 x 10^5 ms^{-1}, τ_{ii} ~ 2 x 10^{-4} s, τ_p ~ 1 x 10^{-4} s, $|\tilde{n}/n_o|$ ~ 0.1, $|\tilde{B}/\tilde{B}_o|$ ~ 0.01, B_o ~ 0.2 T, and the axial dynamo $|\underline{E}_o - \eta \underline{j}_o|$ ~ 5 Vm^{-1}, produces H = 0.35. For T_i= 400 eV, H becomes ≃ 0.5. Also, the power spectra of the magnetic fluctuations has a broad peak at f ~ 10 kH_z or $\omega\tau_{ii}$ ~ 10.

With the value of H determined from known parameters, the relative contribution to ion heating arising from viscosity and recycling now depends on the correlation coefficient g and the normalised frequency $\omega\tau_{ii}$. The line given by Q_{visc}/Q_{recyc}= 1 in the plot of g versus $\omega\tau_{ii}$ separates the two domains where each of these ion heating mechanism dominates. This is illustrated in Figure (3a) for the collisional form of $\mu_{\parallel}$ where the shaded region corresponds to dominant viscous heating.

Viscous heating can be seen to dominate at frequency $\omega\tau_{ii} > H$ over an increasing range of g as $\omega\tau_{ii}$ increases. However, as pointed out above, when $\omega\tau_{ii} > 1$, the coefficient $\mu_{\parallel}^{c}$ is inappropriate.

In the case of the Knudsen number correction, the factor $(2\pi R/\lambda_{mfp})^2$ in eq(28b) decreases with increasing T_i and the domain within which heating arising from recycling dominates that from viscosity expands as shown in Figure (3b).

In the case of the Rusbridge correction, given by eq(29c), recycling dominates ion heating totally for frequency $\omega\tau_{ii} < H/(1 - H^2)^{1/2}$ but at higher frequencies viscosity is seen to make a contribution. The line $Q_{visc} = Q_{recyc}$ is asymptotic to $g = H$, so for values of correlation coefficient $g < H$, recycling dominates viscosity as an ion heating mechanism for all frequencies, as shown in Figure (3c).

We can use eq(22) to estimate the actual value of the correlation coefficient g on-axis in HBTX reversed field pinch discharges:

$$\frac{Q_{visc}}{Q_{recyc}} \simeq \frac{{}^3/_2\ KT_i S}{\frac{1}{2} m_i \langle \tilde{u}^2 \rangle S} - 1 = g^2 (h/H)^2 - 1 \qquad \ldots(29)$$

where $h^2 = {}^3/_2\ (\tau_{ii}/\tau_p)\ (n_o/\tilde{n})^2$. Equating this successively to Q_{visc}/Q_{recyc} for $\mu_{\parallel}^{c}$, $\mu_{\parallel}^{k}$ and $\mu_{\parallel}^{R}$ results in, for the collisional case:

$$g^2 = \frac{H^2}{h^2 - (\omega\tau_{ii})^2}$$

for the Knudsen correction case:

$$g^2 = \frac{H^2}{h^2 - (\omega\tau_{ii})^2\ (2\pi R/\lambda mfp)^2}$$

and for the Rusbridge correction case:

$$g^2 = \frac{H^2}{h^2 - \dfrac{(\omega\tau_{ii})^2}{1 + (\omega\tau_{ii})^2}}$$

These values of g are represented by dashed lines in Figures 3a, 3b and 3c.

In all cases, it can be seen that g = 0.02 at low frequencies and so lies in the region where recycling is the dominant ion heating mechanism. In HBTX1B, the dominant magnetic fluctuations have $\omega\tau_{ii}$ ~ 10. When the collisional form of $\mu_{\parallel}$ is used, g passes into the viscous heating domain near $\omega\tau_{ii}$ = 10 and rapidly goes to unity. In the Knudsen corrected case, g passes into the viscous heating region near $\omega\tau_{ii}$ = 40 and again goes rapidly to unity. In the Rusbridge case, g remains near 0.02 for all frequencies so recycling heating prevails always.

This rather modest degree of correlation g ~ 0.02 between the velocity fluctuations and the magnetic field fluctuations must be readily achievable, and may account for the apparent insensitivity to initial conditions observed when a RFP is set up. It also suggests that attempts to observe this correlation directly will prove difficult.

An estimate of the strength of the source term S on the HBTX axis can be made through eq(18) assuming η_t is negligibly small:

$$^3/_2 \quad KT_eS = \eta_s j^2$$

For T_e (0) = 200kA, assuming j(0) ~ = $4 \frac{I_\phi}{(\pi a^2)}$ produces S(0) = $9.3x10^{22} m^{-3} s^{-1}$.

At the same time , S = n_n n_e <σv>, the product of the neutral density n_n, the electron density n_e, and the rate at which neutrals are ionized, <σv>. Although the latter is a function of position through its dependence on electron temperature, over the range 50eV < T_e< 1000eV, it turns out to be almost independent of T_e in the case of hydrogen [24]. It can accordingly be treated as a constant value, <σv> = $2.5x10^{-14}m^3s^{-1}$ for hydrogen, throughout the discharge. The density of neutrals, measured at various radii by fluorescent scattering in HBTX1B, projects to an axial value $n_e(0)$ = $1x10^{16}m^{-3}$ [25]. Taking an on-axis electron density of $n_e(0)$ = $5x10^{19}m$ accordingly gives S(0) = $1.25x10^{22}m^{-3}s^{-1}$, in modest agreement with the value estimated above. The higher value, corresponding to the shorter particle confinement time $\tau_p = n_e/S$ ~ 0.5 ms, might be thought the more plausible.

A particle diffusion coefficient can be introduced through Fick's Law $\underline{\Gamma} = D\nabla n_e$. Using eq(9) $\nabla.\underline{\Gamma}$ = S, and a parabolic electron density distribution relates S and D on the axis:

$$S(0) = 4 \frac{n_e(0)}{a^2} D(0)$$

from which $D(0) = 29m^2s^{-1}$, in favourable agreement with the axial value deduced from a Monte Carlo calculation based on the measured neutral distribution [26].

Having identified particle recycling as the principal mechanism for ion heating in an RFP plasma, one may ask whether it can indeed furnish energy as fast as required to account for the anomalous heating rate seen in HBTX. We compute an ion energy confinement time on axis, taken to be

$$\tau_{Ei} = \frac{^3/_2\ nKT_i}{\langle \underline{E}.\underline{j}-\eta j^2\rangle_{ion,\ axis}} .$$

Extracting the power going specifically to the ions on axis from eq(19) viz

$$\langle \underline{E}\cdot\underline{j}-\eta j^2\rangle_{ion,\ axis} = KT_iS + m_i\langle\tilde{u}_i^2\rangle S - V_1$$

and using eq(17) to substitute for $m_i\langle\tilde{u}_i^2\rangle S$, we find

$$\tau_{Ei} = \frac{n}{s}\,\frac{^3/_2\ KT_iS}{4KT_iS + V_1}$$

or

$$\tau_{Ei} \sim \frac{3}{8}\tau_p .$$

Thus, on axis, the ion energy confinement time, that is, the time required to heat the ions, is comparable to the particle recycling time n/S, which, as we have seen, is a fraction of a millisecond.

Though it proves to be negligible in standard HBTX discharges, the recycling resistivity η_t, which is different from that associated with helicity loss [10], becomes important at low density and even imposes a limit on the electron streaming parameter, so discouraging electron runaway. On axis, eq(16b) reduces to

$$\eta_t = \frac{2/3\ \epsilon^2}{1-2/3\ \epsilon^2}$$

where the electron streaming parameter $\epsilon \equiv u_{eo}/v_e$. Using typical axial current density of $4x10^6Am^{-2}$ for a 200kA toroidal current, $T_e(0) = 200$ eV and axial electron density $n = 5x10^{19}m^{-3}$, we find $\epsilon = 0.06$ and $\eta_t/\eta_s = 0.002$, confirming that in normal circumstances η_t is very small compared to Spitzer resistivity.

Low density discharges in HBTX had toroidal current of 80 kA and axial electron temperature as high as 800 eV. The possible importance of η_t in this regime is revealed by the accompanying TABLE in which again the axial current density is taken to be about 4 times the average value.

$n(m^{-3})$	$1x10^{19}$	$1x10^{18}$	$0.8x10^{18}$	$0.6x10^{18}$	$0.5x10^{18}$
ϵ	0.06	0.6	0.75	1.0	1.2
η_t/η_s	0.004	0.6	1.2	2.6	25

It can be seen that in this example, when $n < 1x10^{18}m^{-3}$, the recycling resistivity becomes more prominent than the Spitzer. Moreover the streaming parameter is limited to $\epsilon \lesssim \sqrt{3/2}$.

Neglecting the small contribution from viscous heating, the on-axis velocity fluctuation level can be inferred from the axial ion temperature using eq(17). Taking a typical value for $T_i(0) \sim 200eV$ gives $|\tilde{u}_i| \sim 1.7x10^5 ms^{-1}$.

Armed with this, the degree of correlation needed between velocity and density fluctuations to account for the observed diffusion rate can be deduced. Let the relative density fluctuations on axis be $f = |\tilde{n}|/n$. Then flux $|\underline{\Gamma}| = D|\nabla n| \sim D\ n/a$, where a is the plasma radius, 0.25m in HBTX. But by definition, $|\underline{\Gamma}| = \gamma\ |\tilde{n}||\tilde{\underline{u}}| = \gamma fn\ |\tilde{\underline{u}}|$ where γ is a correlation coefficient. Thus $\gamma f\ n\ |\tilde{\underline{u}}| = D\ n/a$ and $\gamma f = D/(a\ |\tilde{\underline{u}}|)$. Since we showed before that on axis $D \sim 30m^2s^{-1}$, $\gamma f \sim 0.7x10^{-3}$ and if the relative density fluctuation $f \sim 0.1$ then the correlation between velocity and density fluctuations needed to account for the axial diffusion rate is only $\gamma \sim 0.7\%$.

The arguments embodied in this paper are constructed around the notion of a recycling process in which plasma loss to the wall is balanced by an influx of neutrals that provide new, cold plasma by ionisation within the discharge. The evidence is that neutrals are not cold except near the walls. Yet replacing the cold source by a hot one in the Boltzmann equation would leave the zeroth and the first moment equations and the conclusions drawn from them unaltered, owing to the isotropy of the hot source velocity distribution. Hot neutrals would however contribute new terms to the equations of the second moment. The ionization of hot neutrals appears to yield additional heating to the ions. This heating, energy, however, is offset by the energy originally absorbed from the ions by the neutrals through collisions, charge exchange and recombination. These processes constitute extra loss mechanisms for particles and energy via transport of neutrals. This is ignored in our model.

Conclusions

Particle recycling has been identified as a mechanism whereby dynamo fluctuations can account for ion heating and transport by turbulent convection in the reversed field pinch. A steady state plasma in which radially outward particle flux arising from and driven by fluctuations is balanced by the appearance of new plasma internal to the discharge has been modelled by constructing Braginskii-type equations from moments of the Boltzmann equation with a source term $S(\underline{x},t)\delta(\underline{v})$. The work done in overcoming the drag exerted on the plasma motion by the new particles provides non-Ohmic ion heating that can compete in effectiveness with and dominate conventional viscous heating in balancing turbulent convection loss.

It is rapid enough to account for the anomalously fast rate of ion heating seen in RFP plasmas, and it explains the observed correlation between ion temperature and loop volts. It accounts for the difference between the resistivities deduced from magnetic energy balance and helicity balance arguments, and it can be shown to be incompatible with a fully relaxed Taylor state, which could be approached only if recycling were to vanish.

The particle source gives rise to a new component of resistivity, η_t, additional to the Spitzer value and to that associated with helicity edge loss. Being inversely proportional to density squared, η_t could become dominant in low density discharges and impede electron runaway.

Under typical conditions in HBTX plasma, the source strength deduced from the power balance near the axis is 10^{22} to $10^{23} m^{-3} s^{-1}$, corresponding to a plausible particle confinement time of a fraction of a millisecond, and to a diffusion coefficient on axis of $D \sim 30 m^2 s^{-1}$, in reasonable agreement with fluorescent scattering measurements.

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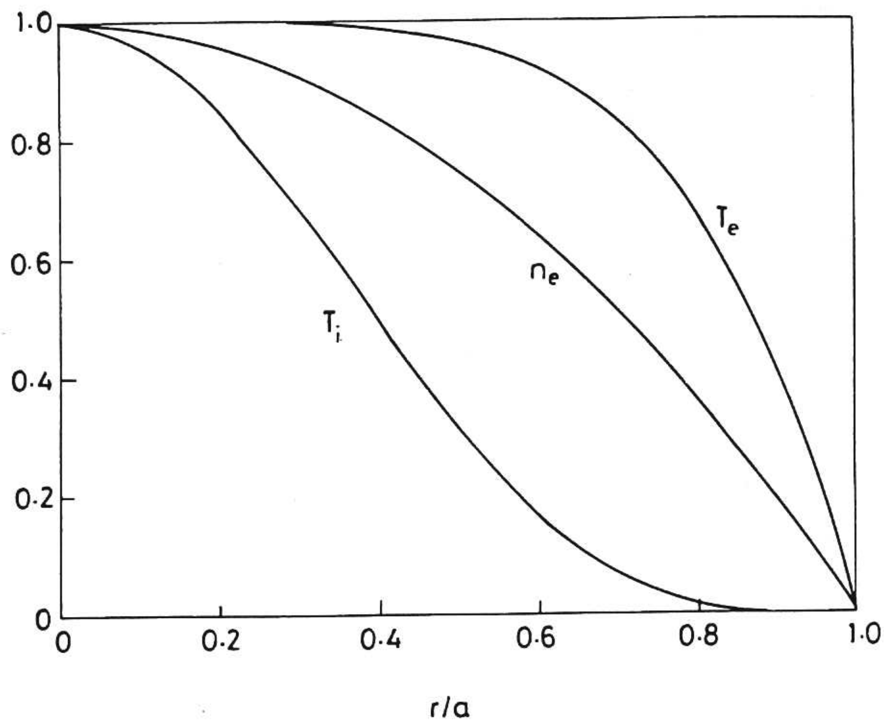

Fig. 1 Schematic of radial profiles of electron and ion temperatures, and density in typical HBTX discharges.

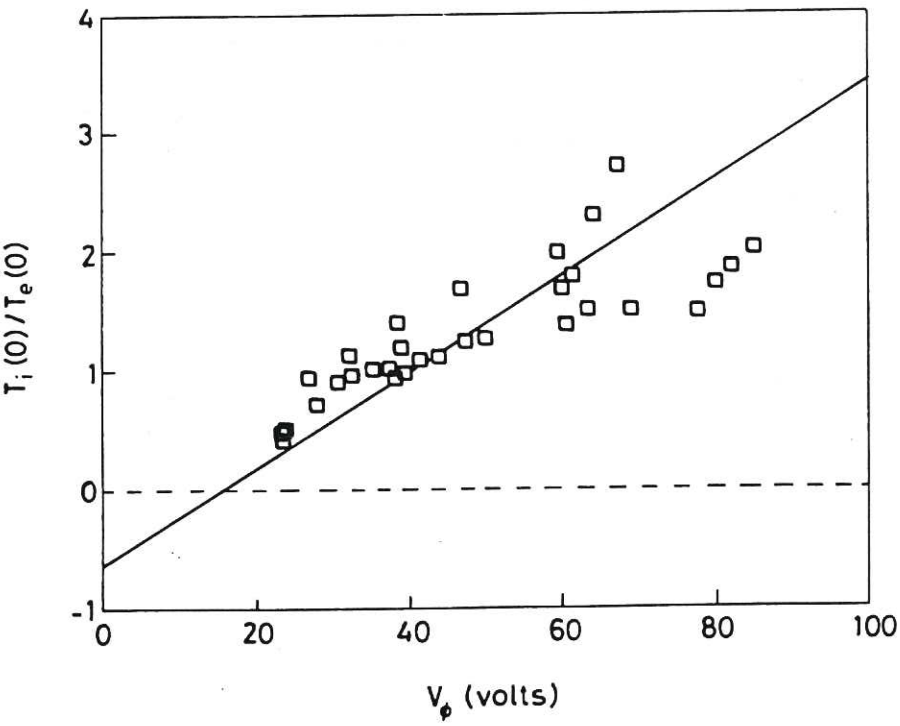

Fig. 2 Experimental ratio of axial temperature to axial electron temperature measured as a function of loop volts. The latter was altered by inserting an obstruction (paddle) while all other parameters were held constant. The line is the theoretical curve given by eq (20).

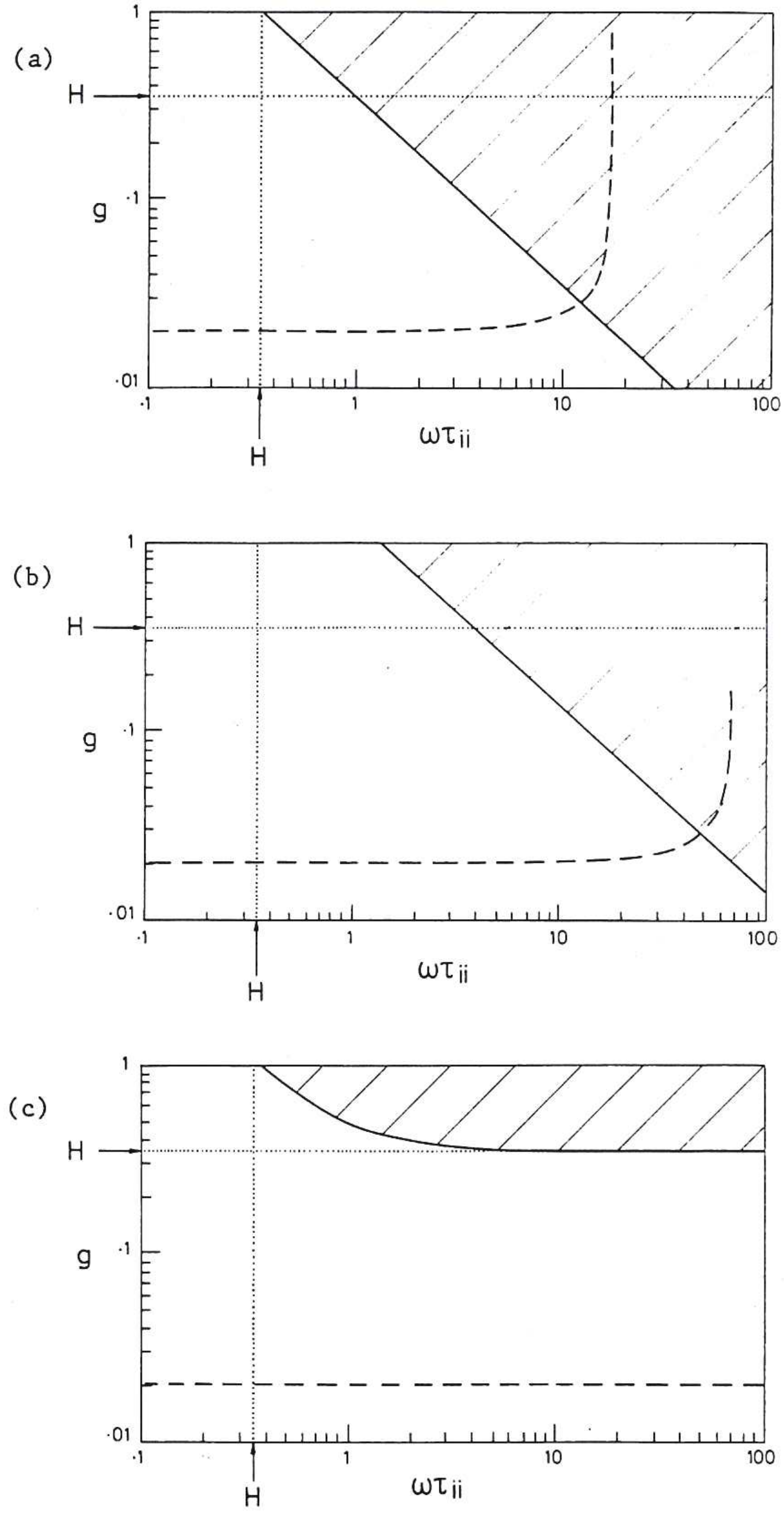

Fig. 3 Relative importance of viscosity and recycling to ion heating. Ordinate is correlation coefficient g associated with dynamo; abscissa is fluctuation frequency ω normalized to the ion-ion collision time τ_{ii}. Viscosity dominates recycling in the hatched region. Dashed line is the locus of g determined from experiment. In HBTX the dominant magnetic field fluctuations has $\omega\tau_{ii} \sim 10$.

(a) Collisional case

(b) Collisionless case: Knudsen correction

(c) Collisionless case: Rusbridge stratagem