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Abstract

An application is made of the authors' magnetic fluctuation interpretation of anomalous electron heat conduction to the collisional regime of slow RFP experiments. The gross thermal properties of Zeta and Eta-Beta II are investigated for the quiescent period. Given magnetic fluctuation levels of between 0·1 and 3·0%, energy conduction loss can be interpreted in terms of classical parallel electron thermal conductivity. For Eta-Beta II the ions are predicted to be hotter than the electrons at the end of quiescence. The magnitude of this effect, which is purely a consequence of the fluctuation theory, depends on the assumptions made.

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1. INTRODUCTION

Using a two-fluid model, we have recently [1] given a theoretical interpretation of temperature fluctuations and heat transfer in tokamaks, both the ions and the electrons being treated as neoclassical. Essentially, we assumed the magnetic field $\vec{B}$ to comprise a mean $(\vec{B}_o(r))$ and a fluctuating part $(\Delta\vec{B})$ such that $\Delta B/B_o \sim 10^{-3}$. We were able to show that the perpendicular electron thermal conductivity depends strongly on the fluctuation level, whereas the ions are largely unaffected, and as found experimentally, the ion thermal conductivity is neoclassical. Our prediction of the electron thermal conductivity was found to be reasonably consistent with the fluctuation levels observed in TFR, PLT and TOSCA, provided that the mean free path is taken to be the major circumference ($2\pi R$) in the formula for the parallel thermal conductivity. The reason for this correction is not far to seek. The collisionality of a plasma with respect to the parallel thermal conductivity depends on the Knudsen numbers $(Kn)_i \sim \frac{v_i^{th}\tau_i}{2\pi R}$ and $(Kn)_e \sim \frac{v_e^{th}\tau_e}{2\pi R}$. For the tokamaks studied, these numbers are much greater than unity, thus indicating that the mean free path loses its significance as a step-length as it approaches $2\pi R$. The objective of the present paper is to test the fluctuation interpretation of heat transfer in experiments operating in the <u>collisional</u> regime, that is, where the above Knudsen numbers are less than unity and where the above correction does not need to be made. To this end we consider the energy balance in the quiescent phase of slow RFP experiments.[1] As far as this work is concerned we assume the quiescent phase to be that period ($\lesssim$ 1ms) of the discharge during which the magnetic/current fluctuations are greatly reduced.

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[1] In RFP experiments it might be argued that the relevant Knudsen numbers should be based on a length of order $2\pi a$. The experimental results, however, appear to be consistent with $2\pi R$ as the appropriate length.

We should point out that our model is specifically restricted to examining the effects of conductive heat transfer as modified by magnetic turbulence. While radiation is crudely allowed for, convection and turbulent heating are ignored. There is no a priori justification for this neglect in the case of the quiescent phase of slow RFP experiments. Only experiment can confirm or otherwise that these effects are unimportant compared with those considered. In the absence of such evidence it is legitimate to ask what conduction alone would imply. In reference [1] and the present work we assume $\frac{\delta n}{n} \ll \frac{\delta T_e}{T_e}, \frac{\delta T_i}{T_i}$. This is consistent with our neglect of convection and turbulent heating. In more recent work, however, we have found this assumption to be inessential in deriving the conductive terms used in this paper. Furthermore, since convection is generally a loss mechanism while turbulent heating is a positive volume source of energy, it follows that there is at least partial cancellation between these effects. Both turbulent convection and heating depend on the correlations between density and velocity fluctuations, and are expected to be of comparable magnitude. If tokamaks are any indication [2, 3], these effects are unimportant for the electrons, while for the ions they may be comparable in magnitude to the neoclassical conduction. Since the present model is very crude, such effects should not change the qualitative features of the results reported here.

On grounds of simplicity we assume that the field fluctuation level is constant throughout the quiescent period. In the experiments considered here the electron temperature increases by 50 - 100% during the quiescent phase. Due to the paucity of the published data it is impossible to make precise or detailed comparisons with theory. However, we show that the flutter required to interpret the available results lies in the range $0\cdot 1\% \leq \Delta B/B \leq 3\cdot 0\%$.

2. REVIEW OF EXPERIMENTAL DATA

We begin by reviewing the three sets of data to be investigated. The first two are based on measurements made in Zeta (minor radius $a = 50$ cm) and the third on recent observations in Eta-Beta II ($a = 12$ cm).

Set I These experiments [4] were carried out for a toroidal field $B_\phi \simeq 3$ kG and a plasma current $I_p \simeq 400$ kA. The time-variation of electron temperature and density measured on axis is shown in reference [5]. The temperature is seen to rise from 100 eV to 150 - 200 eV during the 'quiescent phase' which lasts about one millisecond; in this period the density remains essentially constant at the value $n \simeq 5 \times 10^{13}\,\mathrm{cm}^{-3}$. Although error bars are not shown the data are presumably subject to considerable uncertainty.

Set II These somewhat earlier experiments [6] were carried out for $B_\phi \simeq 750$G and $I_p \simeq 330$ kA. The temperatures quoted refer to average values across the minor radius. During the quiescent phase (1 ms), which was defined such that $0\cdot 1\% < \delta I_p/I_p < 1\%$, the temperature rose from 20 eV to 50 eV with a density typically $n \simeq 10^{14}\,\mathrm{cm}^{-3}$.

Set III The most recent data for the quiescent phase [7] has been obtained in Eta-Beta II for a field $B_\phi \simeq 4$ kG. During this period, which lasts 0·6 ms, the plasma current fell from 180 to 120 kA; in calculations based on this case we shall take $I_p \simeq 150$ kA. The central temperature rose from 50 eV to 100 eV while the electron density remained roughly constant at about $n = 5 \times 10^{14}\,\mathrm{cm}^{-3}$.

For I and II, the toroidal current fell sufficiently slowly in the quiet phase for us to take I_p as constant. Although no information on the radial density, current and temperature profiles has been published, we shall make crude allowances for them, demonstrating that these effects could be significant.

3. THEORETICAL EQUATIONS

Following Newton [5], for cases I and II, we shall assme the total radiated power during quiescence to be $\lesssim 10\%$ of the ohmic input. In the case of Eta-Beta II, however, the radiated power is presently unknown, and hence we obtain results for several different assumed levels of radiation. We shall also make the approximation $\underline{v}\cdot\nabla T_e \ll \partial T_e/\partial t$, that is, the radial convective transport is negligible.

Specifically then, we make the above assumptions and derive zero-dimensional time-dependent energy equations for the ions and electrons. Following earlier work [1, 9] the effective perpendicular electron thermal conductivity is flutter dependent and involves the parallel electron thermal conductivity; we take both the parallel electron and ion conductivity to be classical. With these ideas, our objective is to determine the upper and lower limits on $\Delta B/B$ required to give the observed temperatures.

The time-dependent mean electron and ion energy equations are [1],

$$\frac{3}{2} n \frac{\partial T_e}{\partial t} = \frac{1}{r}\frac{\partial}{\partial r}\left(r K_{\perp e} \frac{\partial T_e}{\partial r}\right) + \frac{3m_e}{m_i}\frac{n}{\tau_e}(T_i - T_e) + \eta j^2 H(r) \qquad (1)$$

$$\frac{3}{2} n \frac{\partial T_i}{\partial t} = \frac{1}{r}\frac{\partial}{\partial r}\left(r \overline{K}_{\perp i} \frac{\partial T_i}{\partial r}\right) - \frac{3m_e}{m_i}\frac{n}{\tau_e}(T_i - T_e) - \frac{1}{r}\frac{\partial}{\partial r}\left(\varepsilon^2 r \overline{K}_{\parallel i} \frac{T_e + T_i}{n}\frac{\partial n}{\partial r}\right)$$

$$- \frac{1}{r}\frac{\partial}{\partial r}\left(\varepsilon^2 r \overline{K}_{\parallel i} \Gamma_e \frac{\partial T_e}{\partial r}\right) \qquad (2)$$

where ε denotes the R.M.S. value of the magnetic fluctuations $(\Delta B_r/B)$ and $\Gamma_e(r)$ is the corresponding form factor. The effective perpendicular electron thermal conductivity $K_{\perp e}$ is given by

$$K_{\perp e} = \overline{K}_{\perp e} + \varepsilon^2 \Gamma_e(r) \overline{K}_{\parallel e} , \tag{3}$$

where $\overline{K}_{\perp e}$ and $\overline{K}_{\parallel e}$, together with $\overline{K}_{\perp i}$ and $\overline{K}_{\parallel i}$, are taken to be the classical or Braginskii [8] values. The validity of this assumption will be discussed later. The term $\eta j^2 H(r)$ represents the effective energy input, $(1 - H(r))$ being the fraction of ohmic heat which is lost by radiation. The last two terms in Eq. (2) are characteristic of our version of the flutter model for transport [1]; they represent a flutter dependent equipartition or ion heating. Thus, although Callen [9] has derived the same formula for $K_{\perp e}$, his theory does not include these terms; this difference results from his neglect of parallel momentum balance.

To proceed, we replace Eqs. (1) and (2) by two zero-dimensional model equations; the form factor $\Gamma_e(r)$ is absorbed into ε to give an effective fluctuation level. Thus we write

$$\frac{3}{2} n \frac{dT_e}{dt} = -\left(\overline{K}_{\perp e} + \varepsilon^2 \overline{K}_{\parallel e}\right) \frac{T_e}{a^2} + 3 \frac{m_e}{m_i} \frac{n}{\tau_e} (T_i - T_e) + \eta j^2 H \tag{4}$$

$$\frac{3}{2} n \frac{dT_i}{dt} = -\overline{K}_{\perp i} \frac{T_i}{a^2} - 3 \frac{m_e}{m_i} \frac{n}{\tau_e} (T_i - T_e) + \frac{\overline{K}_{\parallel i}}{a^2} \varepsilon^2 (T_i + T_e) + \frac{\overline{K}_{\parallel i}}{a^2} \varepsilon^2 T_e \tag{5}$$

where $\overline{K}_{\perp e}\left(\sim \frac{\overline{K}_{\parallel e}}{\omega_{ce}^2 \tau_e^2}\right)$ is small compared with $\overline{K}_{\perp i}$, and of course, with $\varepsilon^2 \overline{K}_{\parallel e}$, since $(\omega_{ce}\tau_e)^{-2} \ll \varepsilon^2$. In writing Eq. (5) we have assumed $\frac{1}{n}\frac{\partial n}{\partial r} \sim \frac{1}{T_e}\frac{\partial T_e}{\partial r}$; this approximation bears particularly on the last two terms (ion flutter heating). Since we have no knowledge as to the relative importance of these terms, we shall bracket our calculations by retaining both, or keeping only $\overline{K}_{\parallel i}\, a^{-2} \varepsilon^2 T_e$. Eqs. (4) and (5) are, of course, only a crude representation of the electron and ion energy balance. A more sophisticated approach would be to assume radial profiles for T_e, T_i and n with time dependent coefficients, and to substitute in Eqs. (1) and (2). Subsequent integration from the axis to the plasma boundary would lead to improved zero dimensional equations. In view of the lack of detailed data, however, we feel that this would be unwarranted.

We now introduce the dimensionless temperatures $x = \frac{T_e}{T_o}$ and $y = \frac{T_i}{T_o}$, where T_o is the electron temperature at the beginning of the quiescent phase ($t = 0$). We also define $s = t/t^*$, where t^* is the period of quiescence. Thus Eqs. (4) and (5) can be expressed as

$$\frac{3}{2}\frac{dx}{ds} = -\frac{\varepsilon^2 E x^{7/2}}{\mu} + R\mu\,\frac{y - x}{x^{3/2}} + \frac{G}{\mu}\,\frac{1}{x^{3/2}} \tag{6}$$

$$\frac{3}{2}\frac{dy}{ds} = -F\mu y^{\frac{1}{2}} + R\mu\,\frac{x - y}{x^{3/2}} + \frac{\varepsilon^2 S}{\mu}\,(y + 2x)\,y^{5/2} \tag{7}$$

where

$$\left.\begin{aligned}
E &= \frac{3\cdot 2\, T_o \tau_o t^* \Theta^{5/2}}{m_e a^2} \\
F &= \frac{2\left(\frac{m_i}{m_e}\right)^{\frac{1}{2}} T_o \tau_o t^*}{(\omega_{ce}\tau_o)^2 m_e a^2 \Theta^{\frac{1}{2}}} \\
G &= \frac{\eta_o I_p^2\, H t^* \Phi^2}{n T_o (\pi a^2)^2 \Theta^{5/2}} \\
S &= E\left(\frac{m_e}{m_i}\right)^{\frac{1}{2}} \\
\text{and}\quad R &= 3\left(\frac{m_e}{m_i}\right)\frac{t^*}{\tau_o \Theta^{3/2}}
\end{aligned}\right\} \tag{8}$$

The self-collision time τ_o is defined to be

$$\tau_o = \frac{3}{4} \frac{(m_e)^{\frac{1}{2}}}{\sqrt{2\pi}} \frac{(T_o)^{\frac{3}{2}}}{n \Lambda e^4} , \quad (9)$$

where T_o is in ergs and Λ denotes the Coulomb logarithm evaluated at T_o. The parameters Θ, μ and Φ are defined by

$$\left.\begin{aligned} \Theta &= \frac{T_i}{T_{i\ peak}} = \frac{T_e}{T_{e\ peak}} \\ \mu &= \frac{n}{n_{peak}} \\ j &= \Phi\ (\pi a^2)^{-1} I_p , \end{aligned}\right\} \quad (10)$$

where $\Phi = 1$ for a flat current profile and $\Phi = \sqrt{3/2}$ for a linearly falling current. The suffix 'peak' denotes the value on axis at $s = 0$. Different choices of Θ, μ and Φ enable us to make a crude allowance for profile effects.

4. COMPARISON WITH EXPERIMENT

In the absence of precise knowledge of the ion temperature we shall assume $T_i = T_e$ at the start of the quiescent phase ($t = 0$).

Set I In considering this data we put $H = 0{\cdot}9$ (10% of ohmic input lost by radiation), and take $\mu = \frac{1}{2}$, $\Theta = \frac{1}{2}$, $\Phi = \sqrt{3/2}$ as roughly representative of the profile effects.

(a) We begin by solving Eqs. (6) and (7) in the absence of fluctuations ($\varepsilon = 0$). Thus the equations only include ohmic heating, electron-ion equilibration and ion conduction, the appropriate coefficients being $G = 2{\cdot}46$, $F = 6{\cdot}2 \times 10^{-3}$ and $R = 3{\cdot}3$ for

$t^* = 1\,ms$. Fig. 1 shows the electron temperature to be always greater than the ion temperature during the quiescent period, with $T_e \simeq 1{\cdot}8\,T_i$ at $s = 1{\cdot}0$. If the ion conduction is also dropped, then essentially the same curves are obtained, indicating ion-conduction to be completely negligible in the present case.

(b) We now solve the full equations for a range of fluctuation levels, namely, $\varepsilon = 0{\cdot}1$, $0{\cdot}2$ and $0{\cdot}5\%$; the coefficients E and S have the values $2{\cdot}8 \times 10^4$ and $4{\cdot}6 \times 10^2$, respectively. The results (see Figs. 2 and 3) clearly show that the observed temperature rise, $T_e \simeq 2\,T_o$, is consistent with fluctuation levels such that $0{\cdot}1\% \lesssim \varepsilon \leqslant 0{\cdot}2\%$. For these levels equilibration is incomplete, with $T_e \simeq 1{\cdot}5\,T_i$. With $\varepsilon = 0{\cdot}5\%$, however, good equilibration is attained, but the electron temperature rise does not match that observed (see Fig. 4). We note that a thermal steady-state is achieved well before the end of the phase in this case.

Set II Since the temperatures quoted for this data are average rather than peak values, we set $\Theta = \mu = 1$, but do retain $\Phi = \sqrt{3/2}$; we again take 10% of the ohmic input to be lost by radiation.

(a) As for the high-temperature Zeta-data, we begin by solving the energy balance equations in the absence of fluctuations ($\varepsilon = 0$). Setting $t^* = 1\,ms$ the appropriate coefficients are $E = 1{\cdot}4 \times 10^3$, $F = 0{\cdot}3$ and $G = 9{\cdot}4$. Fig. 5 shows the electron temperature to be always greater than the ion temperature during quiescence, with $T_e \simeq 1{\cdot}1\,T_i$ at $s = 1{\cdot}0$. Slight increases in the final electron and ion temperatures are discernible if the ion conduction term is dropped; we conclude that this term is unimportant.

(b) We now solve the full equations for the fluctuation levels $\varepsilon = 0{\cdot}2\%$ and $\varepsilon = 1{\cdot}0\%$; the coefficients R and S have the values

27·2 and 23·1, respectively. The results (Figs. 6 and 7) show that as the flutter level increases, equilibration improves. The calculated electron temperature rise is compatible with observation; the ions are always cooler than the electrons. Comparing Figs. 5 and 6 we note that a 0·2% flutter level is unimportant. Even at $\varepsilon = 1{\cdot}0\%$ the electron temperature is scarcely affected. However, the flutter heating of ions is more efficient, and consequently T_i is closer to T_e.

Set III As for Set I we take $\Phi = \sqrt{3/2}$ and $\mu = \Theta = \frac{1}{2}$. The corresponding coefficients E, F, R and S are $8{\cdot}4 \times 10^3$, $9{\cdot}1 \times 10^{-2}$, 57·8 and 138·6, respectively. As mentioned previously, the power lost through radiation is presently unknown for Eta-Beta II. Hence it is necessary to obtain results for several different assumed levels of radiation.

(a) Assuming 10% of the ohmic input to be lost by radiation (H = 0·9), we begin by solving Eqs. (6) and (7) in the absence of fluctuations. Fig. 8 shows the electron temperature to be always greater than the ion temperature during the quiescent period ($t^* = 0{\cdot}6$ ms), with $T_e \simeq 1{\cdot}5\, T_i$ at $s = 1{\cdot}0$. We note that T_e rises from 50 to 350 eV, far higher than that observed (50 to 100 eV). We further note that if ion conduction is omitted, Fig. 8 is unaffected, showing for this case that ion losses are negligible. This result is a clear indication of the existence of some anomalous heat loss, if as assumed only 10% is radiated. Our model attempts to explain this in terms of the flutter-dependent transport.

(b) We now solve the full equations for two fluctuation levels, namely, $\varepsilon = 1{\cdot}0\%$ and $\varepsilon = 3{\cdot}0\%$, with H again taken as 0·9. For the lower ε level, Fig. 9 shows T_e to saturate very rapidly at about $T_e = 2{\cdot}5\, T_o$, while T_i rises very smoothly, eventually crossing the T_e curve. In contrast to the previous cases $T_i > T_e$ at the end of quiescence ($T_i \simeq 1{\cdot}2\,T_e$). This effect is due entirely to the ion-flutter heating. When the flutter level is increased to 3% the electron behaviour is qualitatively the same, although T_e/T_o at $s = 1$ is now of order 1·5. The ions behave as before, except that the cross-over occurs earlier, and at the end of the quiescent period the ions are significantly hotter than the electrons ($T_i/T_e = 1{\cdot}4$). At $s = 1$ we note that $T_i/T_o = 2{\cdot}1$ for $\varepsilon = 3{\cdot}0\%$ (see Fig. 10), while $T_i/T_o = 2{\cdot}75$ for $\varepsilon = 1{\cdot}0\%$. That is, a decrease in flutter level leads to an increase in both ion and electron temperature rises. Finally, we have repeated the $\varepsilon = 3{\cdot}0\%$ case with the first ion-flutter heating term - that involving dn/dr - omitted (see Fig. 10). Although not given in Fig. 10, the calculation shows the electrons to be unaffected by this flutter heating term. Equilibration is greatly improved, steady-states being attained for both species.

(c) In order to estimate the effect of varying H on the final temperature rises and equilibration, several runs have been made, the results of which are given in Figs. 11 - 13. Fig. 11 presents the ion heating ratio at the end of quiescence (T_i^*/T_o) as a function of the radiation efficiency $\eta_R = H \times 100$ for $\varepsilon \equiv \frac{\Delta B}{B} = 1{\cdot}0$ and $3{\cdot}0\%$. These curves show that over a wide variation of η_R, namely, 50 to 100%, the ion temperature rise is not significantly affected. For a given η_R, the heating ratio increases with decreasing fluctuation level. Similar conclusions apply to the electron temperature curves in Fig. 12. It is important to note that these results have been obtained by setting $T_e = T_i = T_o$ at the beginning of quiescence. If more experimental information is forthcoming on this point, the curves can be readily recalculated.

Fig. 13 shows the percentage equilibration ratio at the end of quiescence $\left(\frac{T_i^* - T_e^*}{T_e^*}\right) \times 100$ as a function of η_R for the two fluctuation levels considered. Higher fluctuation levels naturally lead to larger imbalances between electron and ion temperatures. Furthermore, for both levels the imbalance increases with η_R. It is interesting to note, that for the parameters considered, the ions are always hotter than the electrons at the end of quiescence. This is a definite prediction of the flutter model. Braginskii theory, on the other hand, will inevitably lead to negative values of this ratio, whatever the magnitude of η_R.

There are, in fact, two distinct interpretations of the observed electron temperature rise in Eta-Beta II. The first, which we have given above, assumes the radiation loss to be the same as in Zeta, that is, η_R is between 50 and 100%. In this explanation the ions are always found to be hotter than the electrons for fluctuation levels consistent with the observed electron temperature rise. A second interpretation is indicated by calculations for lower fluctuation levels and very high radiation loss. Thus for $\frac{\Delta B}{B} < 1\%$ and $\eta_R \lesssim 10\%$, the computed rise in electron temperature again matches observation. In this regime, however, the electrons are necessarily hotter than the ions. Although at present we do not have the data to distinguish between these two explanations, in principle, this could be done.

5. DISCUSSION

As pointed out in the introduction, the collisionality of a plasma with respect to parallel thermal conductivity depends on the Knudsen numbers $(Kn)_i \sim \frac{v_i^{th} \tau_i}{2\pi R}$ and $(Kn)_e \sim \frac{v_e^{th} \tau_e}{2\pi R}$. We have taken classical formulae for $\overline{K}_{\parallel e}$ and $\overline{K}_{\parallel i}$, and this is justified provided $(Kn)_i$, $(Kn)_e \lesssim 1$. In fact we need only consider the latter. For the low temperature Zeta data and Eta-Beta II, we find $(Kn)_e \ll 1$, while for the high temperature Zeta results $(Kn)_e \lesssim 1$.

Computations show that the choice of profile factors does affect the precise temperature predictions. This is due to the fact that the coefficients in the energy equations are strongly temperature dependent. Qualitatively, however, the results are unchanged. We have found that our results are fairly insensitive to the amount of power radiated provided that this does not exceed 50% of the ohmic input. The flutter model is only applicable to slow experiments. For fast pinches the full dynamics of plasma and fields need to be taken into account, the flutter decomposition being invalid. Unlike our earlier work [1], we have not attempted to relate magnetic fluctuations to temperature flutter, since no data relating to the latter is as yet available.

For Zeta, given the radiation loss to be 10% of the ohmic input and magnetic fluctuations levels of between 0·1 and 1·0%, we have satisfactorily calculated the observed electron temperature rise. In the high-temperature Zeta experiments, the thermal conduction loss is mainly due to the electrons. However, in the low-temperature Zeta experiments, thermal conduction loss may be due to ions or electrons depending on the level of flutter assumed. It is of interest to note that no allowances for profiles need to be made in this case, thus removing one source of uncertainty.

For Eta-Beta II, with 10% radiation loss and ion conduction as the only conductive loss, we cannot explain the observed electron temperature rise. However, given the same radiative loss we <u>can</u> interpret the rise if we assume flutter levels of between 1·0 to 3·0% ; the ions are found to be hotter than the electrons at the end of quiescence. The actual temperature rise is relatively insensitive to the radiation level provided that it is between 10 to 50% of ohmic input. It is possible to interpret the observed rise without any reference to magnetic flutter, if the radiative loss is 85% or higher. In this case, however, the electrons are hotter than the ions.

In our view these predictions clearly indicate the need for detailed measurements of both T_i and T_e as functions of r and t during quiescence; simultaneous measurement of magnetic fluctuation levels should be carried out, as for example, in the TFR tokamak. Subject to the availability of more data the present theory could be developed to any degree of sophistication, thus enabling a detailed test to be made of the magnetic fluctuation interpretation of heat transfer in collisional plasmas. It should be noted that although Rusbridge [10] has attempted a magnetic flutter interpretation (based, however, on different principles from ours), his work relates only to the non-quiescent phase of Zeta. As far as we know, no self-consistent magnetic flutter interpretation of energy transfer in the quiescent phase (including electrons and ions) has yet been published.

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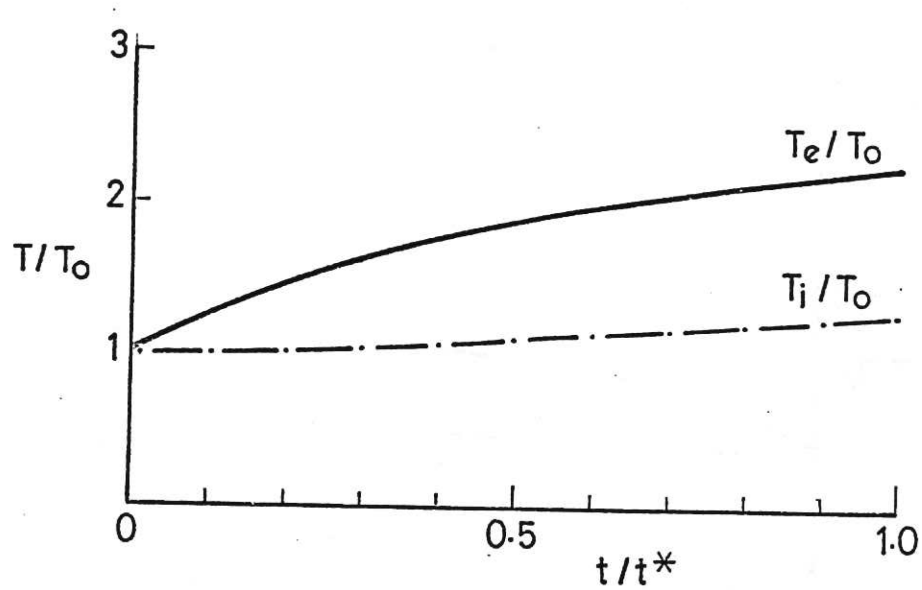

Fig.1 Temperature curves for Zeta (100 eV) without fluctuations but including equilibration and ion conduction.

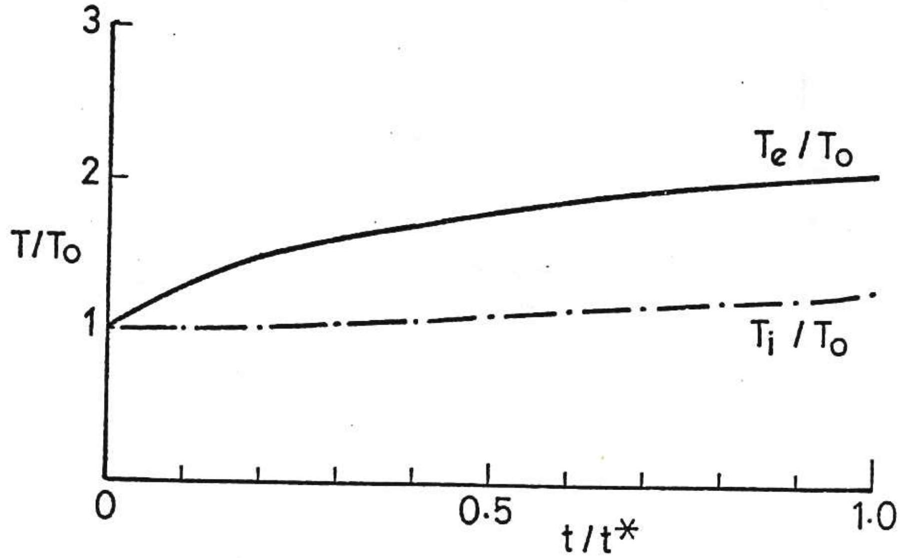

Fig.2 Temperature curves for Zeta (100 eV, ΔB/B = 0·1%) including equilibration and ion conduction.

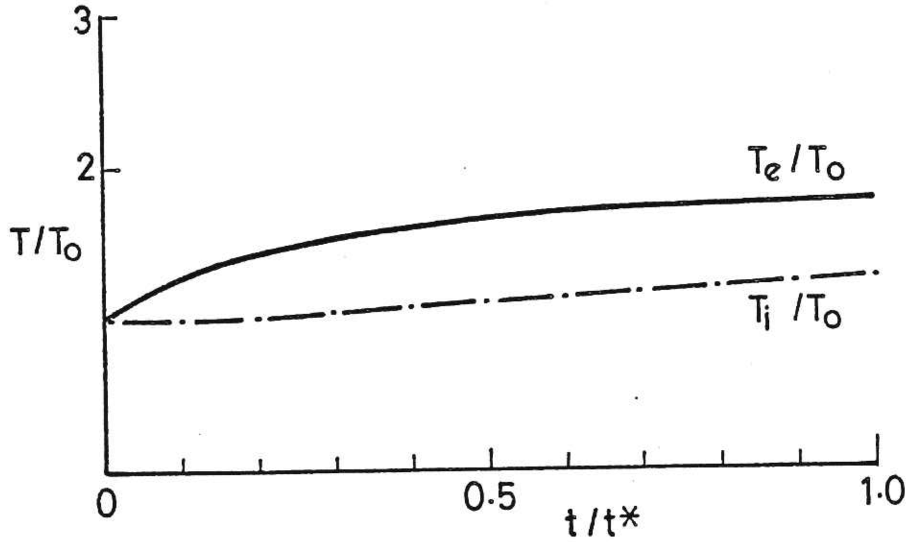

Fig.3 Temperature curves for Zeta (100 eV, ΔB/B = 0·2%) including equilibration and ion conduction.

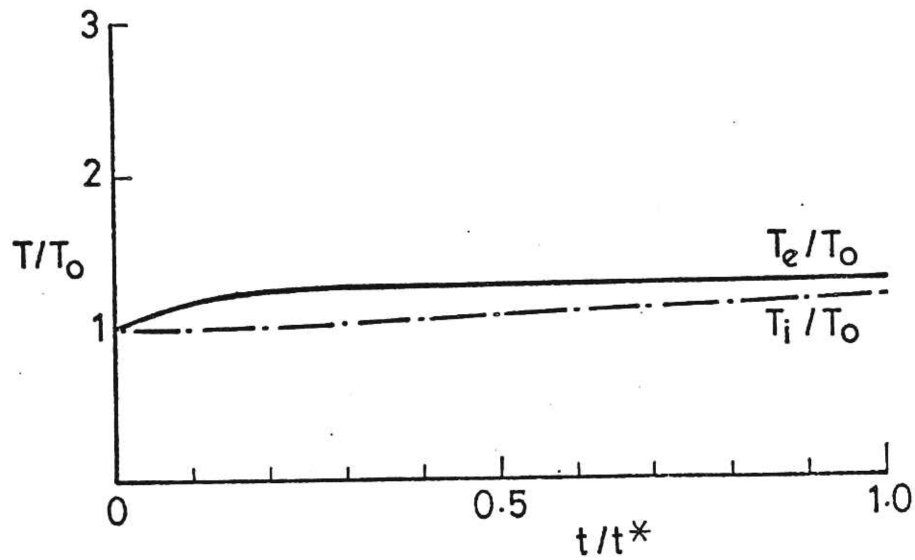

Fig.4 Temperature curves for Zeta (100 eV, ΔB/B = 0·5%) including equilibration and ion conduction.

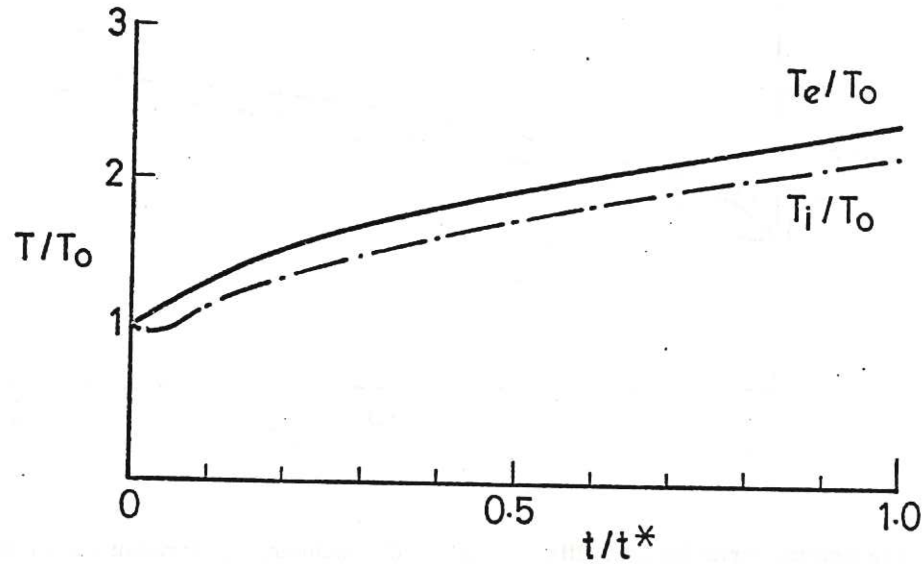

Fig.5 Temperature curves for Zeta (20 eV) without fluctuations but including equilibration and ion conduction.

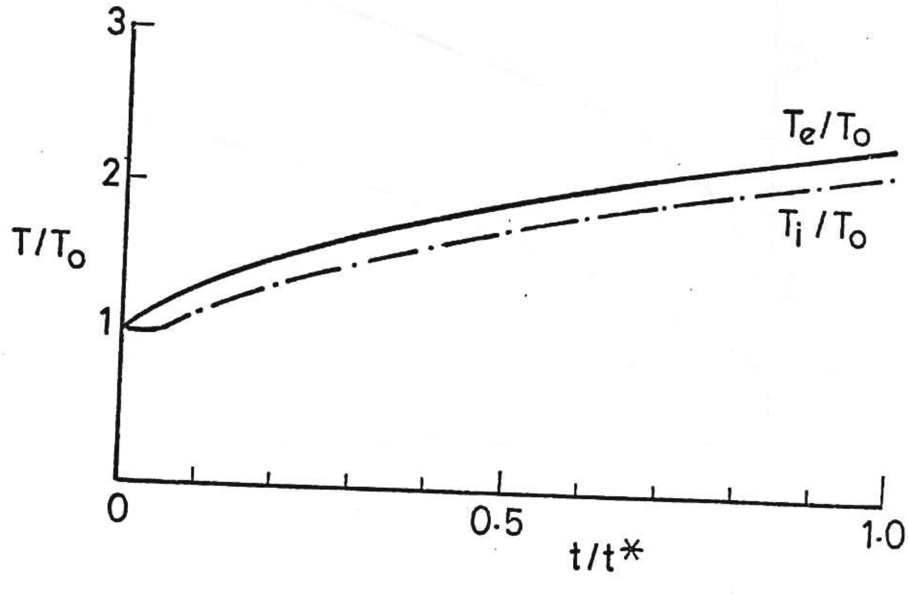

Fig.6 Temperature curves for Zeta (20 eV, ΔB/B = 0·2%) including equilibration and ion conduction.

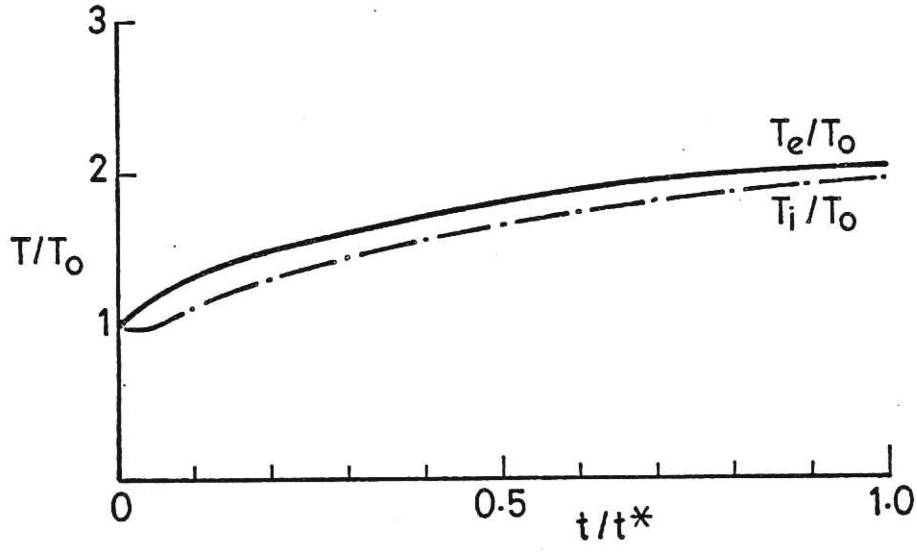

Fig.7 Temperature curves for Zeta (20 eV, $\Delta B/B = 1{\cdot}0\%$) including equilibration and ion conduction.

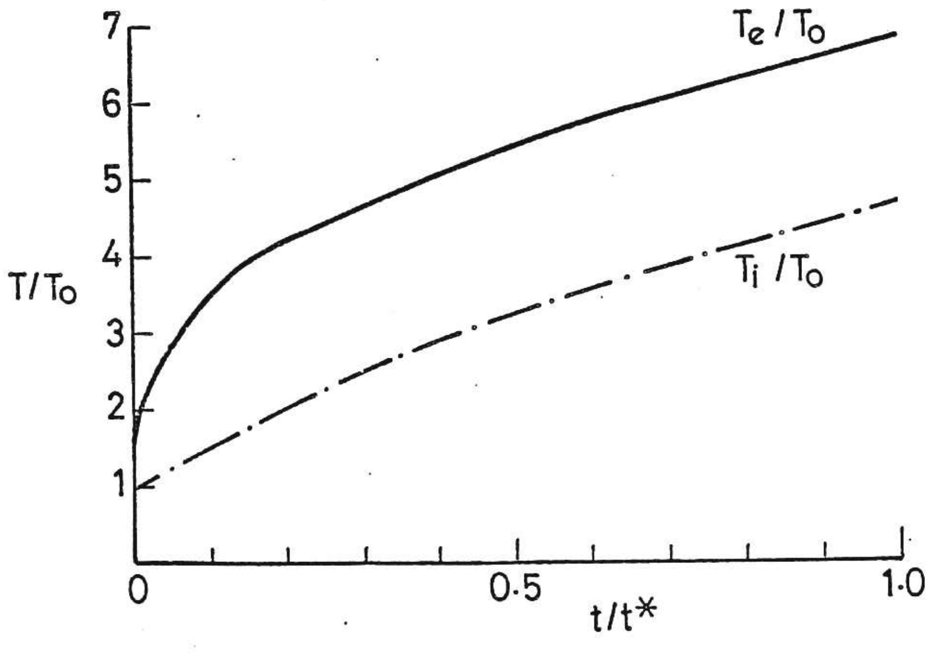

Fig.8 Temperature curves for Eta-Beta II without fluctuations but including equilibration and ion conduction.

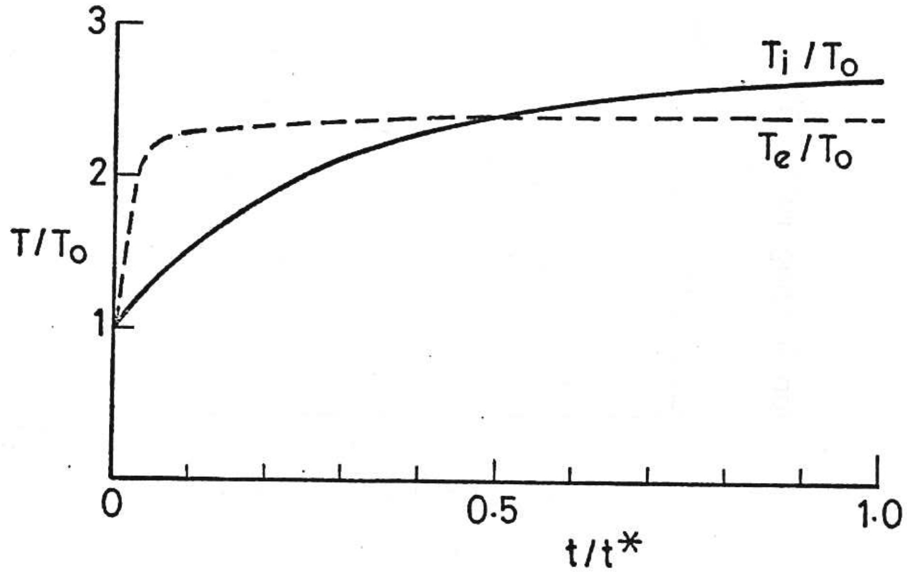

Fig.9 Temperature curves for Eta-Beta II ($\Delta B/B = 1{\cdot}0\%$, $H = 0{\cdot}9$) including equilibration and ion conduction.

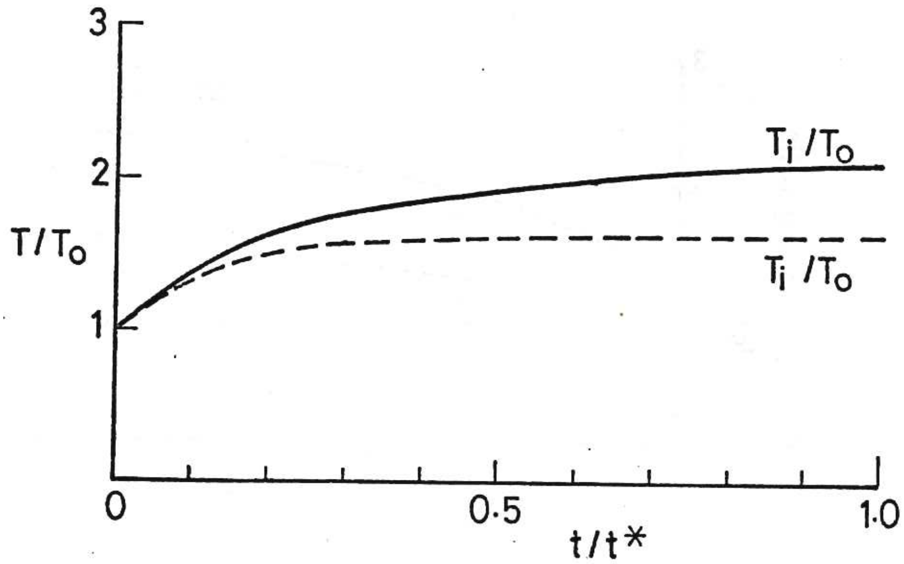

Fig.10 Ion temperature curves for Eta-Beta II ($\Delta B/B = 3{\cdot}0\%$, $H = 0{\cdot}9$) with (solid line) and without (broken line) ion flutter heating.

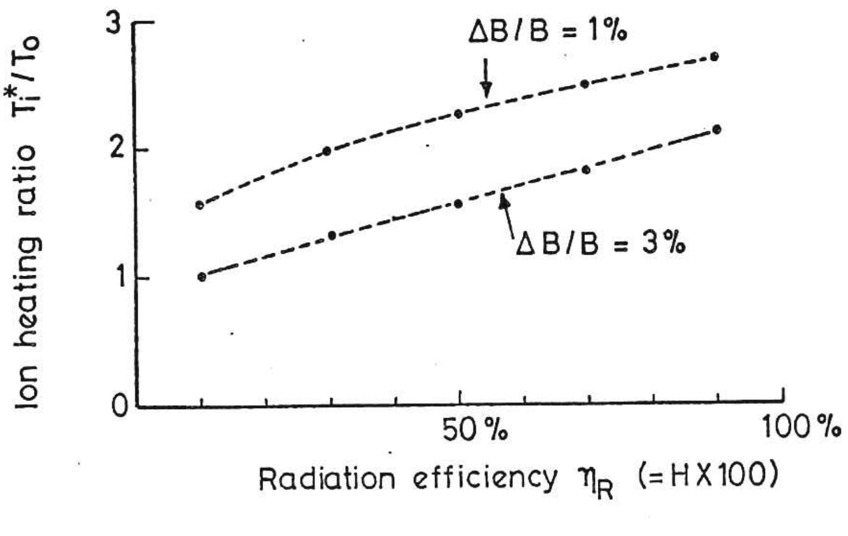

Fig.11 Ion heating ratio (T_i^*/T_o) as a function of radiation efficiency (η_R) for $\epsilon \equiv \frac{\Delta B}{B}$ = 1·0 and 3·0%.

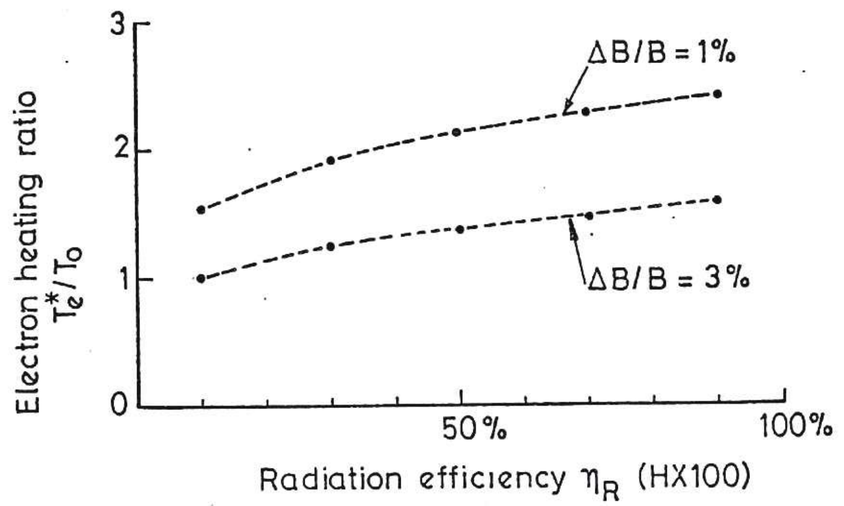

Fig.12 Electron heating ratio $\left(\frac{T_e^*}{T_o}\right)$ as a function of radiation efficiency (η_R) for $\epsilon \equiv \frac{\Delta B}{B}$ = 1·0 and 3·0%.

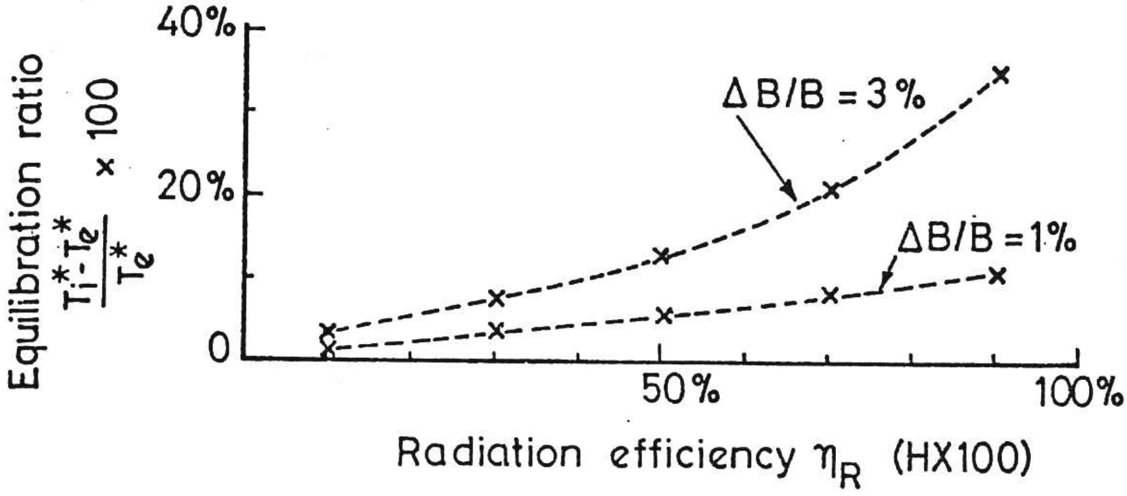

Fig.13 Equilibration ratio $\left(\frac{T_i^* - T_e^*}{T_e^*}\right)$ as a function of radiation efficiency (η_R) for ϵ = 1·0 and 3·0%.