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FLUTTER INTERPRETATION OF THERMAL CONDUCTION WITH APPLICATION TO ALCATOR

I. COOK
F. A. HAAS
A. THYAGARAJA

CULHAM LABORATORY
Abingdon Oxfordshire

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FLUTTER INTERPRETATION OF THERMAL CONDUCTION WITH APPLICATION TO ALCATOR

I. Cook, F.A. Haas, A. Thyagaraja

Culham Laboratory, Abingdon, Oxon, OX14 3DB, UK

(Euratom/UKAEA Fusion Association)

Abstract

A previously presented magnetic fluctuation interpretation of thermal conduction in tokamaks is extended to the high field, high density regime of ALCATOR. The earlier results for the thermal conductivity are unchanged, but additional convective flux and turbulent heating terms are identified. Matching with ALCATOR scaling leads to the conjecture that the typical magnetic field fluctuation levels varies inversely with the average electron density.

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1. INTRODUCTION

In this paper we extend our previous magnetic fluctuation interpretation of thermal conduction [1] to the high field, high density regime of ALCATOR. Under these conditions, experiment [2] indicates the density fluctuations to be significant. Hence, it is necessary to take account of both density and temperature fluctuations in our analysis. This apart, our approach follows that described in our earlier paper. We reiterate that our work is an interpretation of transport rather than a theory, in the sense, that we do not attempt to determine the causes of the turbulence. In an earlier investigation, Molvig et al [3] also sought to explain thermal conduction in ALCATOR as due to magnetic fluctuations. Unlike our analysis, however, which is based on the two-fluid equations, their study is related to Rechester and Rosenbluth's [4] treatment of stochastic individual particle orbits.

We show in Section 2 that the thermal conduction is unaffected by density fluctuations. In the course of the analysis, however, we identify two additional features, namely, convective flux and turbulent heating. In Section 3 we attempt to interpret the experimental results of Gondhalekar et al [5]. Following this, in Section 4 we relate our work to the so-called ALCATOR scaling, and are led to the conjecture that the magnetic field fluctuation level varies inversely with the average electron density. The discussion in Section 5 considers the relevance or otherwise of the temperature fluctuation measurements of Arunasalam et al [6] and Cavallo et al [7] to our interpretation. This section also contains a brief discussion of Drake et al's [8] microtearing mode description of anomalous transport. Section 6 gives our conclusions.

2. MEAN ENERGY BALANCE EQUATIONS

Following our earlier work [1] we take the energy equation for the electrons to be

$$\frac{3}{2} n \left(\frac{\partial T_e}{\partial t} + \vec{v}_e . \vec{\nabla} T_e\right) + n\, T_e \vec{\nabla} . \vec{v}_e = \nabla . (\bar{K}_{\parallel e} \nabla_{\parallel} T_e + \bar{K}_{\perp e} \nabla_{\perp} T_e)$$

$$+ 3 \left(\frac{m_e}{m_i}\right) \frac{n}{\tau_e} (T_i - T_e) + S_e \qquad (1)$$

where $\bar{K}_{\parallel e}$ and $\bar{K}_{\perp e}$ are the neoclassical coefficients of thermal conduction[(1)] and S_e represents the various sources and sinks. We make the very reasonable assumption that $\bar{K}_{\parallel e} \gg \bar{K}_{\perp e}$. All quantities in Eq. (1) are to be regarded as instantaneous values, not averages.

As before, taking the basic (mean) state to be cylindrical, we express any physical quantity as the sum of a mean and a fluctuating part. Thus, for example, the magnetic field is written as

$$\vec{B} = \vec{B}_o(r) + \Delta\vec{B}\ (r, \theta, z, t) \equiv \vec{B}_o(r) + \varepsilon\, \delta\vec{B} \qquad (2)$$

where ε is a small parameter denoting the level of field fluctuations. Substituting these forms in Eq. (1) and carrying out the averaging procedure defined previously [1] we obtain

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(1) For $\bar{K}_{\perp e}$ and $\bar{K}_{\perp i}$ we take the usual expressions. For $\bar{K}_{\parallel e}$ and $\bar{K}_{\parallel i}$, however, we use the forms previously discussed [1]

$$\frac{1}{r}\frac{d}{dr}\left(rK_{\perp e}\frac{d\,T_{oe}}{dr}\right) + S^*_{oe}\ (r) = 0 \qquad (3)$$

where, unlike in our earlier work the fluctuations in the coefficients $\bar{K}_{\parallel e}$, $\bar{K}_{\perp e}$ have also been taken into account. The effective perpendicular thermal conductivity, $K_{\perp e}$, is given by

$$K_{\perp e} = \left\langle \bar{K}_{\perp e} \right\rangle + \left\langle \bar{K}_{\parallel e} \right\rangle \varepsilon^2 \Gamma_e \qquad (4)$$

where

$$\Gamma_e = \left\langle \left(\frac{\delta B_r}{B_o}\right)^2 \right\rangle + \left\langle \frac{\delta B_r}{B_o}\, \Delta_\parallel\, (\delta T_e) \right\rangle \left(\frac{d\,T_{oe}}{dr}\right)^{-1} \qquad (5)$$

$\Delta_\parallel$ being the derivative along $\vec{B}_o$. The effective source, S^*_{oe}, can be expressed as

$$S^*_{oe} = S_{oe}(r) + 3\,\frac{m_e n_o}{m_i \tau_e}\,(T_{io} - T_{eo}) - \frac{3}{2}\frac{1}{r}\frac{d}{dr}\left(r\left\langle \delta p_e\ \delta v_{er}\right\rangle\right)\varepsilon^2$$
$$- \left\langle \delta p_e\ \nabla\cdot\delta\vec{v}_e \right\rangle \varepsilon^2 + \frac{3}{2}\left\langle \delta\sigma_e\ \delta T_e \right\rangle \varepsilon^2 + \text{higher-order terms}, \qquad (6)$$

where $\delta\sigma_e$ is the fluctuating electron source. The higher-order terms comprise both third-order correlations and small second-order contributions which arise from the fluctuations in S_e, the equipartition term and $\bar{K}_{\perp e}$. Thus we see that inclusion of δn and fluctuations in the coefficients $\bar{K}_{\perp e}$, $\bar{K}_{\parallel e}$, makes no difference to the effective electron thermal conduction term in Eq. (3). From Eq. (6), however, we note that the effective source is modified, S^*_{oe} now including two additional terms. The first, $-\frac{3}{2}\frac{1}{r}\frac{d}{dr}\left(r\left\langle \delta p_e\ \delta v_{er}\right\rangle\right)\varepsilon^2$ is a form of turbulent thermal convection. The second, $-\left\langle \delta p_e\ \nabla\cdot\delta\vec{v}_e\right\rangle\varepsilon^2$, is the work done by the fluctuations, and is therefore, a form of turbulent heating.

The mean energy equation for the ions can be similarly reduced, leading to

$$\frac{1}{r}\frac{d}{dr}\left(r\overline{K}_{\perp i}\,\frac{d\,T_{oi}}{dr}\right) + S^{*}_{oi}(r) = 0 \tag{7}$$

with S^{*}_{oi} given by

$$S^{*}_{oi} = \frac{1}{r}\frac{d}{dr}\left(r\overline{K}_{\parallel i}\,\varepsilon^{2}\Gamma_{i}\,\frac{d\,T_{oi}}{dr}\right) - \left\langle \delta p_{i}\;\nabla.\vec{\delta v}_{i}\right\rangle \varepsilon^{2}$$

$$-\frac{3}{2}\frac{1}{r}\frac{d}{dr}\left(r\left\langle \delta p_{i}\;\delta v_{ir}\right\rangle\right)\varepsilon^{2} + \frac{3}{2}\left\langle \delta\sigma_{i}\;\delta T_{i}\right\rangle\varepsilon^{2} + 3\frac{m_{e}\,n_{o}}{m_{i}\,\tau_{e}}\left(T_{eo} - T_{io}\right)$$

$$+ S_{oi}(r) + \text{higher-order terms}, \tag{8}$$

where

$$\Gamma_{i} = \left\langle\left(\frac{\delta B_{r}}{B_{o}}\right)^{2}\right\rangle + \left\langle\left(\frac{\delta B_{r}}{B_{o}}\right)\Delta_{\parallel}\;(\delta T_{i})\right\rangle\;\left(\frac{d\,T_{oi}}{dr}\right)^{-1} . \tag{9}$$

From our earlier work we have the relation

$$\frac{\delta B_{r}}{B_{o}} = -\;L_{p}\,\Delta_{\parallel}\left(\frac{\delta_{p}}{p_{o}}\right) \tag{10}$$

where

$$p_{o} = n_{o}\;(T_{io} + T_{eo}) \tag{11}$$

and

$$L_{p}^{-1} = \frac{1}{p_{o}}\;\frac{dp_{o}}{dr}\;. \tag{12}$$

Using Eqs. (5) and (10) together with the normalisation $\left\langle\left(\frac{\delta B_{r}}{B_{o}}\right)^{2}\right\rangle = 1$, and taking cognisance of the identity

$$\frac{p_o}{n_o L_p} - \frac{d\,T_{oi}}{dr} - \frac{d\,T_{oe}}{dr} = (T_{oi} + T_{oe})\frac{1}{n_o}\,\frac{dn_o}{dr} \qquad (13)$$

we obtain

$$\frac{1}{r}\frac{d}{dr}\left(r\bar{K}_{\parallel i}\ \varepsilon^2\Gamma_i\ \frac{d\,T_{oi}}{dr}\right)$$

$$= -\frac{\varepsilon^2}{r}\ \frac{d}{dr}\left\{r\bar{K}_{\parallel i}\left[(1 + \xi(r)\ \frac{T_{oe} + T_{oi}}{n_o}\ \frac{dn_o}{dr} + \Gamma_e\ \frac{d\,T_{oe}}{dr}\ \right]\right\}$$

(14)

where

$$\xi(r) = \left\langle \frac{\delta B_r}{B_o}\ \Delta_{\parallel}\left(\frac{\delta n}{n}\right)\right\rangle \left(\frac{1}{n_o}\ \frac{dn_o}{dr}\right)^{-1} . \qquad (15)$$

Eq. (14) differs from the equivalent form derived in our earlier paper only through $\xi(r)$; this quantity represents the correlation between density and magnetic fluctuations. As in the case of the electrons, the effective ion source, S^*_{oi} , also contains two additional terms, namely, ion convection and turbulent heating.

From Eq. (7) we observe that the effective ion thermal conduction is just the neoclassical. As regards the effective ion source, the limited scope of our analysis only allows us to make a few qualitative remarks. For present tokamaks $\xi(r)$ is of order unity, and hence as previously noted [1], the first term in S^*_{oi} is unimportant. Some experiments [9] suggest that, in contrast to the electrons, the convective ion heat flux could be comparable with the neoclassical ion thermal conduction. In an experimental investigation of power balance for the ions, such an effect would appear as a measured anomaly of the ion perpendicular conductivity.

Crude estimates indicate this to be a few times the neoclassical value. Turning to the turbulent heating, we remark that this process must be necessarily dissipative. Since ion viscosity is large compared to electron viscosity, it is reasonable to expect that ion turbulent heating is much greater than the corresponding electron heating. Thus in principle, turbulence provides a mechanism for directly heating the ions.

3. COMPARISON WITH ALCATOR

For Alcator [5] we take the following parameters: $R = 54\text{cm}$, $a = 10\text{cm}$, $B_{tor} \simeq 60\text{kg}$, $I_p \equiv 300\text{kA}$, $q(a) \sim 2.0$, $n = 2 \times 10^{14}\text{cm}^{-3}$, $T_i \simeq T_e \simeq 1\text{keV}$, $\frac{\Delta B_\theta}{B_\theta} = 0.3\%$. The magnetic fluctuation level quoted is associated with the observed $m = 2/n = 1$, $m = 3/n = 2$ modes. The mean free paths, λ_e and λ_i, are typically of order $2.0 \times 10^3\text{cm}$ and $3.0 \times 10^3\text{cm}$, respectively. Thus the appropriate Knudsen numbers ($Kn = \lambda/2\pi R$) are $(Kn)_e \simeq 0(6)$ and $(Kn)_i \simeq 0(10)$. Since these values are greater than unity, we follow the procedure discussed previously [1], that is, we set $\bar{K}_{\parallel e} = Cn\,(2\pi R)^2 \tau_e^{-1}$. In comparing with PLT, TFR and TOSCA, we found that taking $C = 0.25$ gave good qualitative agreements; in the present paper we propose to assume the same value. From the fluctuation level for $\frac{\Delta B_\theta}{B\theta}$ and taking $q \simeq 2.0$, we estimate that $\frac{\Delta B_r}{B_{tor}} \simeq 3.0 \times 10^{-4}$. Thus we deduce that

$$K_{\perp e} \simeq \frac{1}{4}\, n \;\frac{(2\pi R)^2}{\tau_e}\; \varepsilon^2 \simeq 3.0 \times 10^{17}\text{cm}^{-1}\ \text{sec}^{-1} \qquad (16)$$

or equivalently, the thermal diffusivity is

$$\chi_{\perp e} = \frac{K_{\perp e}}{n} \simeq 1.5 \times 10^3\text{cm}^2\ \text{sec}^{-1} \ . \qquad (17)$$

Unfortunately, to our knowledge, the ALCATOR group has not quoted a value for $\chi_{\perp e}$ in the literature. We note, however, that the above value is in good agreement with the 'measured' values quoted by Berezovskii et al [10] and Waltz and Guest [11] for other tokamaks.

Consideration of the collisionality indicates that ALCATOR operates roughly in the banana-plateau transition. Thus it is of interest to compare $\chi_{\perp e}$ with the ion thermal diffusivity obtained from the formulae for these two regimes. We find that $\chi_{\perp i\,plateau} \simeq 6.0 \times 10^3 cm^2\ sec^{-1}$ and $\chi_{\perp i\,banana} \simeq 1.0 \times 10^3 cm^2\ sec^{-1}$. Comparing these values with Eq. (17), we find general agreement with the observation that at densities of about $2.0 \times 10^{14} cm^{-3}$, the dominant thermal conduction loss changes from electronic to ionic.

It is instructive at this point to ask whether the observed perpendicular electron thermal conductivity can be interpreted in terms of classical parallel thermal conductivity (that is, no Knudsen correction to $K_{\parallel e}$) and the measured magnetic fluctuation level of $0(10^{-4})$. We find $K_{\parallel e}$ (Braginskii) $\sim 10^{27} cm^{-1}\ sec^{-1}$ and this leads to $\chi_{\perp e}$ typically of order $5 \times 10^5 cm^2\ sec^{-1}$. This is $100 \times \chi_{\perp i\,neoc}$ which is entirely in contradiction with experiment.

Some qualitative conclusions regarding the significance of convection, can be inferred from the experimental results of Gondahalekar et al. In particular, we refer to Fig. 9 of their paper which compares $n_e(r)$ and $T_e(r)$ profiles at $B_{tor} = 3.5T$ and 7.7T, the mean density and plasma current being held constant. Whereas the density profile is virtually unaltered, the electron temperature profiles are markedly different. At the smaller value of q ($B_{tor} = 3.5T$), the temperature profile is lower and flatter compared with the temperature profile for $B_\phi = 7.7T$. We would expect MHD activity to be greater for the lower q discharge, being due to low m, n modes of order 10^{-4}. From the fact that the n_e profiles are almost identical, it is reasonable to infer that particle diffusion is not seriously influenced by these modes. Yet clearly, the electron temperature distribution is substantially altered. This suggests that electron energy transfer is primarily by conduction as interpreted in our model, and that convective transport for <u>electrons</u> is unimportant. Unfortunately a corresponding conclusion cannot be drawn for the ions in the absence of published ion temperature profiles.

4. SCALING LAWS

Gondhalekar et al have presented evidence for the variation of the experimentally determined $\chi_{\perp e}$ with the mean density, $\bar{n}_e$[(2)]. For low densities ($n_e < 3 \times 10^{14} \mathrm{cm}^{-3}$) they state that $\chi_{\perp e}$ (experiment) $\propto \frac{1}{\bar{n}_e}$. In this density range the electron Knudsen number is 0 (10), and we expect on the basis of our model that $\bar{K}_{\parallel e} \simeq \frac{1}{4} \frac{(2\pi R)^2 n}{\tau_e}$. The $\chi_{\perp e}$ predicted by our theory is:

$$\chi_{\perp e} = \frac{K_{\perp e}}{n} \simeq \frac{1}{4} \frac{(2\pi R)^2}{\bar{\tau}_e} \left\langle \left(\frac{\Delta B_r}{B_o} \right)^2 \right\rangle \tag{18}$$

where $\bar{\tau}_e\ (\bar{T}_e, \bar{n}_e)$. Furthermore, following our earlier work the dominant magnetic fluctuations are taken to be the low m, n modes. Since $\bar{\tau}_e \propto (\bar{n}_e)^{-1}$, $\chi_{\perp e}$ will be proportional to $(\bar{n}_e)^{-1}$ if

$$\sqrt{\left\langle \left(\frac{\Delta B_r}{B_o} \right)^2 \right\rangle} \propto \frac{1}{\bar{n}_e} \, . \tag{19}$$

It is of interest to note that in TFR $\sqrt{\left\langle \left(\frac{\Delta B_r}{B_o} \right)^2 \right\rangle} \sim 0.1\%$, $\bar{n}_e$ being

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(2) Note that the bar over n_e, T_e denotes that these quantities have been averaged over the minor cross-section. This convention does not apply to $\bar{K}_{\parallel e}$, $\bar{K}_{\perp e}$, etc.

of order $10^{13}cm^{-3}$, whereas in ALCATOR $\sqrt{\left\langle\left(\frac{\Delta B_r}{B_o}\right)^2\right\rangle} \sim 0.01\%$ with $\bar{n}_e$ of order $10^{14}cm^{-3}$. This suggests that an experimental investigation of the variation of magnetic fluctuation level with $\bar{n}_e$ (and $\bar{T}_e$, B_ϕ, I_p) should be undertaken in a given machine. Such a study could verify, or otherwise, the basic ideas of our work and their relation to the so-called ALCATOR scaling. It would also have profound implications for the future course of reactor design. To see this we give a brief discussion of the scaling laws for energy confinement.

The total energy loss in a plasma can be represented by a global loss rate, or an inverse confinement time, $\bar{\tau}_E^{-1}$. The loss occurs through three channels: radiation, conduction, convection. Each of these may be crudely represented by its own loss rate, that is, we have $\frac{1}{\bar{\tau}_{rad}}$, $\frac{1}{\bar{\tau}_{cond}}$, $\frac{1}{\bar{\tau}_{conv}}$. In general, we expect them to depend on various global parameters (e.g. $\underline{T}_p$, $\bar{n}_e$, B_ϕ ....) and their profiles. Theory should provide the actual dependences for the different parameter regimes. As an example, let us consider $\frac{1}{\bar{\tau}_{cond}}$. Clearly, conductive loss through the ion channel alone provides us with a typical conduction loss rate (neoclassical) which we denote by $\frac{1}{\bar{\tau}_{i\,neoc}}$. Since $\bar{\tau}_{cond} / \bar{\tau}_{i\,neoc}$ is non-dimensional, it can only be a function of all the dimensionless quantities which characterise a given discharge. Some of these quantities are well-known, e.g. $\frac{a}{R}$, β, $\omega_{ce}\tau_e$, $\bar{n}_e\lambda_d^3$, etc. We have seen, however, that other dimensionless quantities like the ratio of the average magnetic fluctuation energy to the mean field energy, $\left\langle\left(\frac{\Delta B}{B_o}\right)^2\right\rangle$, play a fundamental role in electron thermal conduction. If the external circuitry and mass and energy sources have zero noise, then turbulent quantities of this type can only result from saturation of intrinsic instabilities, and therefore ought to be functions of the mean plasma properties and profiles already listed. As yet, of course, we do not know whether the turbulence is self-generated or driven; for the present we have to regard the turbulence quantities as independent parameters characterising the state of the plasma. In principle, they are no different from the other well-known plasma parameters, in the sense that they are susceptible to experimental measurement. Thus $\bar{\tau}_{cond} / \bar{\tau}_{i\,neoc}$ is a function of q, β, etc. and turbulence properties like $\left\langle \left(\frac{\Delta B}{B_o} \right)^2 \right\rangle$. In the absence of a complete theory of turbulence in plasmas which is valid for all parameter ranges, analysis alone cannot be expected to give useful information on the so-called scaling-law for conduction. Although experiment can determine $\bar{\tau}_{cond} / \bar{\tau}_{i\,neoc}$ as a function of all the relevant non-dimensional parameters in certain ranges, there is no a priori guarantee that extrapolation to other ranges is valid. Indeed, it is well known in ordinary fluid turbulence that transport coefficients (e.g. drag-coefficient in aerodynamics) are not smooth or uniform functions of the relevant dimensionless parameters characterising the system.

The present uncertainty over the scalings can only be precisely resolved by experiment and relevant theoretical analysis. In our view it is unrealistic to expect to be able to produce scalings valid under reactor conditions based solely on current experiments in the absence of a more profound understanding of plasma transport than exists at present.

5. DISCUSSION

Our theory is an attempt to correlate $\frac{\Delta B}{B_o}$ with $K_{\perp e}$. In our earlier work, we used momentum balance to find a relation between $\frac{\Delta B}{B_o}$ and $\frac{\Delta p}{p_o}$. The experiments of Arunasalam et al [6] led us to suppose that $\frac{\Delta p}{p_o} \sim \frac{\Delta T_e}{T_e}$, and the magnitude of $\frac{\Delta T_e}{T_e}$ predicted was the order of that observed, namely, $\frac{\Delta T}{T_e} \sim 10\%$. Cavallo et al [7] have recently questioned the validity of the interpretation of the original temperature fluctuation measurements [6]. Cavallo et al claim that the observed $\frac{\Delta T_e}{T_e}$ could be entirely due to fluctuations associated with the signal-to-noise ratio in the apparatus, and that this could not be ruled out on the basis of the data published by Arunasalam et al. While the status of the claim of Cavallo et al is not clear, it is evident that this matter has to be thoroughly re-examined.

Our position regarding the controversy is the following:

(1) Our theory predicts the correlations between $\frac{\Delta B}{B_o}$, $\frac{\Delta p}{p_o}$ and $\frac{\Delta B}{B_o}$, $\frac{\Delta T_e}{T_e}$ (see Eqs (2) and (15) of reference [1]). It is important that they should be checked experimentally. Since radiation is the dominant loss channel at the edge, it is necessary that our relations be checked in the central region of tokamak. This is likely to be difficult, however, since $\frac{\Delta B}{B_o}$, $\frac{\Delta T_i}{T_i}$ and $\frac{\Delta T_e}{T_e}$ are hard to measure in the region of interest. Nonetheless, for a complete validation or otherwise of our fluctuation interpretation of heat transfer, such measurements are essential.

(2) The present work shows that the assumption $\frac{\Delta n}{n} \ll \frac{\Delta T_e}{T_e}$ is not necessary to estimate the anomalous thermal conduction.

(3) While Cavallo et al have raised an important point which needs to be clarified, in our view, they have not demonstrated conclusively that the fluctuations observed by Arunasalam et al

in the frequency range 0 - 450kHz are not turbulent temperature fluctuations. The experimental results of Cavallo et al on TFR refer to the frequency range 1 - 100MHz.

As we pointed out in our earlier work the turbulent fluctuations which contribute dominantly to the thermal transport are most likely to have low m, n and low frequencies, say 0 - 100kHz for PLT. Since in our analysis $\langle(\Delta B)^2\rangle$ is finite and very much smaller than B^2, we expect the power-spectrum of field or pressure fluctuations to decrease with frequency beyond the range 0 - 100kHz. Thus while the observations of Cavallo et al in the range 1 - 100MHz may well be consistent with their interpretation, we believe it to be irrelevant to our interpretation of anomalous transport.

Following our initial paper, Drake et al [8] have advanced an alternative theory of anomalous transport. Their approach differs from ours in four essential respects:

(1) They propose a mechanism for the magnetic fluctuations, namely, saturated microtearing (high - m) modes localised at rational surfaces.

(2) Their formula for the effective perpendicular electron thermal conductivity is not derived from the calculations presented in their paper.

(3) It follows from (2) that no Knudsen correction is applied, nor is any attempt made to assess the relative importance of the various mode numbers.

(4) The governing equations take no account of like-particle collisions. In our approach $\bar{K}_{\parallel e}$, which arises from electron collisions only, is fundamental both in the mean and the fluttered equations.

6. CONCLUSIONS

We have extended our previous magnetic fluctuation interpretation of thermal conduction to the high field, high density regime of ALCATOR. In the course of this we have taken account of density as well as temperature fluctuations. Our earlier results for thermal conduction are unaltered. However, new thermal convection and turbulent heating terms have been identified. While these are not discussed in detail, in principle their effects can be separated experimentally from those of conduction. In matching with ALCATOR scaling we are led to conjecture that the magnetic fluctuation level, $\left\langle \left(\frac{\Delta B}{B_o}\right)^2 \right\rangle^{\frac{1}{2}}$, varies like $\frac{1}{\bar{n}_e}$. This effect could be sought in existing tokamaks.

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