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**ELECTRON ENERGY TRANSPORT AND MAGNETIC
CURVATURE DRIVEN MODES**

By

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ABSTRACT

A transport coefficient for anomalous electron thermal conduction is constructed on the basis of the so-called Principle of Profile Consistency. It is assumed that the relevant modes in plasma regimes where a substantial fraction of the electron population is magnetically trapped produce magnetic reconnection at a microscopic level and are driven by the combined effects of the plasma pressure gradient and the magnetic field curvature. Consequently, the scaling for the electron energy confinement time exhibits a strongly favorable dependence on the radius of magnetic curvature.

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A class of anomalous transport coefficients that can reproduce the results of a variety of experiments on magnetically confined plasmas was introduced in earlier work [1] and was based on the so-called "Principle of Profile Consistency" [2]. One of the main assumptions involved is that the current density profile relaxes toward a class of ("canonical") profiles that resemble a Gaussian. The validity of this hypothesis is in fact supported by results from recent experiments carried out on the Alcator C machine featuring the injection of pellets in high density plasmas [3]. Specifically, the electron temperature profile, which is closely coupled to that of the current density, has been observed to recover its canonical profile following the abrupt perturbation due to the pellet penetration into the plasma column over times shorter than 1/100th of the electron energy confinement time.

In order to arrive at a specific form of the electron thermal conductivity, one must also specify the main characteristics of the type of microscopic process (e.g., microinstability) that allows this current density to relax to its canonical profile. In previous papers [1-3] the type of process considered for this was a form of current-driven microscopic magnetic reconnection, and the adopted representative dimensionless parameter was

$$C_s = v_{eff} d_e^2 \times \frac{v_{ee}}{v_{the}^2} \quad (1)$$

where $v_{eff} d_e^2$ is the relevant anomalous magnetic diffusion coefficient with $d_e = c/\omega_{pe}$, v_{ee} is the electron-electron collision frequency, and v_{the} the electron thermal velocity. In particular, v_{eff} was taken to be

$$v_{eff} = \frac{v_s}{v_{the}} \times \omega_{pi} \quad (2)$$

where v_s is the ion sound velocity and ω_{pi} the ion plasma frequency.

In the present note we consider a different type of process which is relevant in regimes where a substantial fraction of the electron population is magnetically trapped and which involves modes driven by the combined effects of the plasma pressure gradient and of the unfavorable magnetic field curvature. Comprehensive linear studies of this general class of instabilities indicate that under realistic operating conditions for present day experiments, modes of this type are likely to be present over a relatively wide range of frequencies and wave numbers [5].

Consequently, instead of using C_s as the representative "number" with v_{eff} given by Eq. (2), we can adopt

$$C_N = \frac{v_s}{(r_{po}R)^{1/2}} d_e^2 \left(\frac{a_i}{R}\right)^{1/2} \times \frac{v_{ee}}{v_{the}^2} \left(\frac{a_i}{R}\right) \quad (3)$$

Here a_i represents the width of the plasma current channel [3], and

$$\frac{I}{r_{po}} = \left. \frac{d \ln p}{dr} \right|_{\max}$$

with r_{po} being the scale distance of the maximum pressure gradient and R the radius of magnetic curvature. Thus

$$C_N = \frac{1}{L_1} \times \frac{1}{T_e^2} \times \left(\frac{Z_i}{A_i}\right)^{1/2} \times \left(\frac{a_i}{R}\right)^{3/2} \quad (4)$$

where $L_1 \equiv (r_{po}R)^{1/2}$, and Z_i and A_i are the ion charge and mass number.

The electron heat flow is expressed as

$$q_e = -\kappa \frac{dT_e}{dr} \quad (5)$$

where $J_{\parallel}$ is the plasma current density. Considering for simplicity the case where the prevalent loss of plasma thermal energy occurs through the electron channel, the relevant expression for κ can be derived following the procedure given in Ref. 3.

In particular, when only ohmic heating is present, we obtain

$$\kappa = \frac{\epsilon_N}{24\pi} \frac{I(r)}{r^2 J_{\parallel}} \frac{a_I \bar{v}}{8\pi} \quad (6)$$

where $I(r) \equiv \pi dr^2 J_{\parallel}(r)$, $a_I \equiv a(6q_0^0/q_s)^{1/2}$, a is the radius of the plasma column, the current density distribution is represented [3] approximately by

$$J_{\parallel} \approx J_0^0 \exp\left(-\frac{q_s}{q_0^0} \frac{r^2}{a^2}\right) \quad (7)$$

and

$$\bar{v} \equiv \frac{T_e}{e} C_N^{1/2} \quad (8)$$

Notice that

$$\bar{v} \propto \frac{1}{L_1^{1/2}} \left(\frac{Z_1}{A_1}\right)^{1/4} \left(\frac{a_I}{R}\right)^{3/4}$$

and, recalling that the electron thermal energy balance equation is, in this case,

$$V_{\parallel} I(r) = 4\pi^2 R \kappa \frac{d \ln J_{\parallel}}{d \ln r} \quad (9)$$

we obtain

$$V_{||} = \epsilon_N \times \left(\frac{R}{4a}\right) \times \left(\frac{q_s}{6q_0}\right)^{1/2} \times \bar{V} \quad (10)$$

Here ϵ_N is a numerical coefficient that is found by fitting the results of a consistent set of confinement experiments. Thus we find that the loop voltage scales as

$$V_{||} \propto \left(\frac{R}{a_I}\right)^{1/4} \frac{1}{L_{||}^{1/2}} \left(\frac{Z_i}{A_i}\right)^{1/4} \quad (11)$$

Notice that the scaling for $V_{||}$ here exhibits no dependence on density. However, experimental results indicate that this should be a weakly increasing function of density. We could argue that such an effect may be caused by the contribution of the (collisional) ion thermal conductivity neglected for simplicity in the present calculation. On the other hand, we note that the scaling associated with the form of C_g [3] given in Eq. (1) is

$$V_{||} \propto n^{1/5} \left(\frac{R}{a_I}\right) \left(\frac{Z_i}{A_i}\right)^{2/5} \quad (12)$$

and that this introduces a weak density dependence even when ion thermal conduction effects are ignored.

We can obtain a scaling for the electron temperature (assuming that the electrical resistivity $\eta_{||}$ is classical) by noting that

$$\frac{V_{||}}{R} \sim \eta_{||} J_{||} \propto \frac{Z_{eff}}{T_e^{3/2}} J_{||}$$

This implies

$$\langle T_e \rangle = B_p^{2/3} L_1^{1/3} \left(\frac{A_1}{Z_1} \right)^{1/6} \left(\frac{R}{a_1} \right)^{1/2} Z_{eff}^{2/3} \quad , \quad (13)$$

where $B_p = I_{||}/a$ is the poloidal field at the edge of the plasma column and $\langle T_e \rangle$ is the volume-averaged electron temperature. We can then obtain the following scaling for the electron energy containment time

$$\tau_e \sim \frac{\langle n T_e \rangle R}{V_{||} \langle J_{||} \rangle} = n a^{7/6} R^{5/3} \times \frac{Z_{eff}^{2/3}}{B_p^{1/3}} \left(\frac{A_1}{Z_1} \frac{r_{po}}{a} \right)^{5/12} \quad . \quad (14)$$

Finally, we consider the case where auxiliary heating (e.g., by neutral beam injection) prevails over ohmic heating. We assume that the maximum plasma pressure that can be well confined is described by the parameter β_p ("beta-poloidal") and that this is related to the total injected power by

$$W_T^e = \epsilon_N \times \left(\frac{R}{4a_1} \right) \times \bar{v} I_{||} \times (1 + \gamma_d \beta_p) \quad . \quad (15)$$

This expression is a natural extension of Eq. (10). Here γ_d is a degradation coefficient that depends on the departure of energy deposition profile from the canonical profile, which is characteristic of ohmic heating, and $I_{||}$ is the total plasma current. This in turn leads to

$$\kappa \approx \frac{\epsilon_N}{18\pi} G_w(r) \frac{I_{||}^2 a_1 \bar{v}}{r_{J_1}^2 8\pi} \beta_p \gamma_d \quad . \quad (16)$$

where

$$G_w(r) = 2\pi^2 R \frac{\int dr^2 S_x^e(r)}{W_T^e} \quad (17)$$

is a source profile function, with S_x^e indicating the power per unit volume

deposited on the electrons. Thus the electron energy confinement time scales as

$$\tau_e \propto \frac{1}{V} B_p a^2 \propto I_{\parallel} R \left(\frac{A_1}{Z_1}\right)^{1/2} a^{1/2} \left(\frac{r_{po}}{a}\right)^{1/4}.$$

Correspondingly, the electron temperature scales as

$$T_e \propto \frac{S_x^e}{n} \frac{I_{\parallel} a}{V} \propto \frac{S_x^e}{n} I_{\parallel} a L_{\perp}^{1/2} \left(\frac{A_1}{Z_1}\right)^{1/2} \left(\frac{R}{a}\right)^{3/4}.$$

Finally, we observe that although the anomalous transport model proposed in earlier work [1-3] has led to numerical simulation results in remarkable agreement with the experimental observations for ohmically heated plasmas, the relevant scaling does not exhibit the strong major radius dependence proposed in recent papers [6]. In this note, the model presented exhibits a favorable dependence of confinement time on major radius while preserving the profile-consistent properties of the earlier model. In general, we would caution that conclusions and directions for future experiments inferred from empirical data should be carefully assessed because there exists an intrinsic arbitrariness [7] in deducing energy confinement time scalings.

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