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Observations on Energy Transport in Tokamaks: A New Formula for Turbulent Transport

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Abstract

Arguments for a new, modified Ohkawa type formula of anomalous electron energy transport due to electromagnetic microturbulence $(D_\perp \sim \omega/k_\perp^2)$ are presented and its predictions are compared with tokamak experiments.

I. Introduction

Because of its importance to the scientific feasibility of magnetically confined fusion, the subject of energy transport in tokamaks has been intensively investigated with a view to obtaining scaling laws for energy confinement.

Energy transport in tokamaks has several essential features:

1. it is anomalous, and takes place mainly through the electron channel (the dominance of electron energy transport);[1–5]

2. various auxiliary heating schemes (independent of their nature) lead to a deterioration of the energy confinement time;[6,7]

3. the electron temperature tends to keep its own "proper" profile.[8]

The second characteristic implies that the origin of the anomalous energy transport may be more or less independent of the nature of the heating source. Consequently, it is likely that the underlying mechanism for the deterioration of energy confinement time in auxiliary heating cases is probably the very mechanism responsible for the anomalous energy transport in ohmic heating.

The last characteristic implies the decoupling of the electron energy and particle transports. Since the electron and ion particle transport must be strongly coupled through ambipolarity consideration, both the dominance and the independence of electron energy transport suggests that there exists a leading anomalous process that preferentially acts on the electrons and leaves the ions essentially unaffected.

Since the low-frequency microturbulence is generally supposed to be responsible for the anomalous transport,[9] the aforementioned experimental observations would favor microturbulence based on magnetic fluctuations. The low-frequency plasma response is essentially controlled by the parallel component of Ampere's law making the ion dynamics irrelevant to the process.

The observed density dependence[10,11] of the anomalous transport is another important factor for preferring the magnetic microturbulence in a collisionless plasma; pure electrostatic fluctuation cannot have any density dependence (collisions, of course, could change the picture).

At this stage, it must be clearly stated that the anomalous transport may be a result of a variety of processes. A model based on limited processes is bound to have limitations. However, if a particular process has the potential for explaining some salient features of the experimental observations, it must be taken seriously. It is precisely in this spirit that we are arguing for the usefulness of a model based on the electromagnetic fluctuations.

We begin by making a comment on the linear stability of the e.m.f. in collisionless tokamak plasmas. It is found that the electromagnetic fluctuations are strongly damped by the effects of magnetic shear. The shear damping is usually measured by $\left\langle k_{\|}\right\rangle v_e$, where $\left\langle k_{\|}\right\rangle$ is the average shear-induced parallel wavenumber, and v_e is the electron thermal speed. Obviously electromagnetic fluctuations can grow to significant levels only if $\left\langle k_{\|}\right\rangle v_e$ were to become sufficiently small. We shall soon see what "sufficiently small" means in this context. We must admit, however, that it is quite difficult to drive the electromagnetic modes unstable in linear theory. Although several possibilities are being currently pursued, for the purpose of this article, we shall simply assume that an appropriate unstable electromagnetic mode has been found. This mode, for example, can either be an unstable mode due to the shear damping suppression by some nonlinear effect, or a forced mode driven by some other fluctuations. We proceed and explore its consequences. This exploration has to be of a "generic" or a general nature because the details of the instability are not available. In fact, the details of the linear theory (as long as enough energy is available in the mode) are not even relevant for the development of our argument.

Applying the methods of renormalized perturbation theory to the nonlinear electron gyrokinetic equation, we can derive the equation for turbulent transport of energy[12]:

$$\frac{\partial}{\partial t} nT = \nabla \cdot \left(\frac{c}{B}\right)^2 \mathrm{Im} \sum_k (\mathbf{k} \times \mathbf{b}) \Gamma_0(b) (\mathbf{k} \times \mathbf{b}) m \cdot \int dv_{\parallel} v_{\parallel}^2 G_k \nabla F_0\left(v_{\parallel}^2\right) \left\langle \delta\psi_k \delta\psi_k^* \right\rangle + \cdots \tag{1}$$

where the notation is standard; $G_k = \left(\omega - k_{\parallel} v_e + i\Gamma\right)^{-1}$, $F_0\left(v_{\parallel}^2\right) = \left(\sqrt{\pi} v_e\right)^{-1} \exp\left(-v_{\parallel}^2 / v_e^2\right)$ is the one-dimensional Maxwellian, $b = \rho_e^2 k_{\perp}^2 / 2$, ρ_e is the electron gyroradius, $\mathbf{k}$ is the wavevector, $\mathbf{b}$ is the unit vector along the ambient magnetic field $\mathbf{B} = \mathbf{b} B$, $\Gamma_0(b) = I_0(b) e^{-b}$, I_0 is the modified Bessel function, and $\delta\psi_k = \delta\varphi_k - \left(v_{\parallel}/c\right) \delta A_k$ with $\delta\varphi_k$ as the perturbed electrostatic potential and δA_k as parallel component of the perturbated vector potential. In Eq. (1) the drag term is neglected. For the time being, we neglect $\delta\varphi_k$ implying

$$\left\langle \delta\psi_k \delta\psi_k^* \right\rangle = \frac{v_{\parallel}^2}{c^2} \left\langle \delta A_k \delta A_k^* \right\rangle \simeq \frac{v_{\parallel}^2}{c^2} \frac{\left\langle \delta B_k \right\rangle^2}{k_{\perp}^2}, \tag{2}$$

where $\delta B_k = (\nabla \times \delta \mathbf{A})_k$ is the perturbed magnetic field. Equations (1) and (2) suggest a general form for the turbulent energy transport coefficient driven by the magnetic fluctuations

$$D_{\perp} \sim \frac{v_e^2}{\omega - \left\langle k_{\parallel} \right\rangle v_e + i\Gamma} \left(\frac{\delta B}{B}\right)^2, \tag{3}$$

with $\Gamma \sim k_{\perp}^2 D_{\perp}$. The existence of three different inverse time scales (ω, $\left\langle k_{\parallel} \right\rangle v_e$, Γ) in the denominator of Eq. (3) allows the possibility of a variety of scalings for $D_{\perp}$. Separation into the following two independent regimes is obvious:

1. Weak Turbulence Scaling: If the turbulence induced frequency is much smaller than the basic frequencies associated with the mode ($\Gamma \ll$, ω, $\left\langle k_{\parallel} \right\rangle v_e$), and $\omega \lesssim \left\langle k_{\parallel} \right\rangle v_e$, then one obtains the Rechester-Rosenbluth[13] type of diffusion coefficient

$$D_{\perp} \sim \frac{v_e}{\left\langle k_{\parallel} \right\rangle} \left(\frac{\delta B}{B}\right)^2. \tag{4}$$

2. Strong Turbulent Scalings: In the opposite limit of strong turbulence where Γ is not negligible compared to ω and/or $\left\langle k_{\parallel} \right\rangle v_e$, several distinct possibilities emerge:

(a) The ordering $\Gamma = k_\perp^2 D_\perp \sim \left\langle k_{||} v_e \right\rangle > \omega$ provides the basis for the Ohkawa[14–20] type scaling

$$D_\perp \sim \frac{v_e \left\langle k_{||} \right\rangle}{k_\perp^2} \tag{5}$$

which coupled with (3) gives the turbulence level to be

$$\left\langle \frac{\delta B}{B} \right\rangle \sim \frac{\left\langle k_{||} \right\rangle}{k_\perp}. \tag{6}$$

(b) A change of ordering between ω and $\left\langle k_{||} \right\rangle v_e$, i.e., $\Gamma = k_\perp^2 D_\perp \sim \omega > \left\langle k_{||} \right\rangle v_e$ leads us to a different type of scaling

$$D_\perp \sim \frac{\omega}{k_\perp^2} \tag{7}$$

with

$$\left\langle \frac{\delta B}{B} \right\rangle \sim \frac{\omega}{k_\perp v_e}. \tag{8}$$

It is interesting to observe that the saturation level $\delta B/B$ of Eq. (8) is an expression of the condition that the electron drift speed in the fluctuating fields ($v_e \delta B/B$) equals the perpendicular phase velocity of the wave. This balance is fundamentally different from the conventional balance where the nonlinear damping rate is balanced by the linear growth rate to yield $D_\perp = \gamma/k_\perp^2$.[21] Clearly the saturation mechanism leading to Eq. (8) has no dependence on γ, and results from a balance of the linear current with the nonlinear current induced by the fluctuations. In order to make $D_\perp$'s in Eqs. (4)–(8) as genuine diffusion coefficients, we must specify the typical or maximizing values of ω, $\left\langle k_{||} \right\rangle$ and $k_\perp^2$. Since almost all kind of tokamak microturbulence is considered to be due to the expansion free energy, $\omega = \omega_*$ (where ω_* is the diamagnetic drift frequency due to both the density and temperature gradients), is an obvious choice. For collisionless electron dynamics, the parallel component of Ohm's law leads to a natural $k_\perp = (\omega_{pe}/c) = \delta_e^{-1}$, the inverse of the collisionless skin depth. Finally, for the use of Eqs. (4)–(6), the average value of $k_{||}$ is taken to be the shear length $\hat{s}/qR$, where R is the major radius of the tokamak, q is the safety factor, and $\hat{s}$ is the shear parameter.

Notice that $\left\langle k_{||} \right\rangle$ play a fundamental role in both the Rechester-Rosenbluth type weak-turbulence as well as the Ohkawa type strong-turbulence scalings. We strongly feel that the inequality $\left\langle k_{||} \right\rangle v_e > \omega$ (which is needed to obtain these two scalings) is precisely the wrong limit for the electromagnetic modes to grow in a collisionless plasma. In fact, unless the effective $\langle k_{\|} \rangle$ is suppressed by some mechanism (there does exist a plausible nonlinear theory for $\langle k_{\|} \rangle$ suppression,[22,23] and $\langle k_{\|} \rangle$ is no longer measured by $\hat{s}/qR$), it will be inconsistent to assume that the dominant transport process stems from essentially magnetic fluctuations.

The newly proposed strong turbulence scaling of Eq. (7) ($\omega > \langle k_{\|} \rangle v_e$ is assumed) was motivated by the idea of restoring consistency, i.e., the "weakened" shear assumption was made so that the required modes could, in principle, grow (we stated earlier, that at this time, we do not have a well-defined instability mechanism). In the context of Eqs. (7)–(8) (and the discussion immediately following them), we propose that the diffusion coefficient (ξ is a number to be determined by a detailed nonlinear theory, or fixed by comparison with experiments)

$$D_{\perp}^{(T)} \equiv \xi \frac{c^2}{\omega_{pe}^2} \cdot \omega_* \equiv \xi \frac{c^2}{\omega_{pe}} \frac{1}{eB} \left| \frac{dT_e}{dr} \right| \left[1 + \frac{\alpha}{\eta_e} \right] \tag{9}$$

with the expected fluctuation level

$$\frac{\langle \delta B \rangle}{B} = \xi \frac{\rho_e}{L_T} \left(1 + \frac{\alpha}{\eta_e} \right) \tag{10}$$

follows the ordering inherent in the contention that the anomalous transport is primarily of magnetic origin, which is strongly suggested by the experiment. In Eqs. (9) and (10) we have chosen to explicitly display the electron temperature gradient dT_e/dr, the scale length $L_T^{-1} = (T_e^{-1} |dT_e/dr|)$ etc. ($\alpha \simeq 1$), because in most of the confinement region the electron temperature gradients are sharper than the density gradients, i.e., $\eta_e \equiv T_e^{-1} |dT_e/dr| \big/ n_e^{-1} |dn_e/dr| > 1$.

In this paper, the scaling of $D_{\perp}$ was arrived at heuristically. However, this can be rigorously derived from a theory of collisionless electromagnetic turbulence based on the Vlasov equation (in which the shear suppression is modelled by setting $\langle k_{\|} \rangle = 0$), and the parallel component of Ampere's law as long as one stipulates that somehow the required modes could grow to suprathermal levels. The interested reader can look up the details in Ref. 24.

The rest of this paper is devoted to an investigation of the energy transport in tokamaks using the turbulent diffusion coefficient $D_{\perp} = D_{\perp}^{(T)}$ of Eq. (9). A detailed comparison with the experiments has also been presented.[25,26]

Substitution of Eq. (9) in the energy transport equation (in a cylindrical model)

$$\frac{1}{r} \frac{d}{dr} r D_{\perp}^{(T)} \frac{d}{dr} n T_e = P_e(r) \tag{11}$$

leads to the dimensionless system (where, for simplicity $\eta_e > 1$ has been assumed)

$$\frac{1}{x}\frac{d}{dx}\left[\frac{x}{\tilde{n}(x)^{1/2}}\frac{d\tilde{T}_e}{dx}\frac{d}{dx}\tilde{n}(x)\tilde{T}_e(x)\right] = \kappa\tilde{P}_e(x), \tag{12}$$

$$\kappa \equiv \frac{\Lambda^2}{2\xi T_*^2} \equiv eB\omega_{pe}(0)a^3\bar{P}_e/2\xi c^2 n(0)T_*^2, \tag{13}$$

where $x = r/a$ is the normalized radial variable, $\tilde{n}(x) = n(r)/n(0)$, $n(0)$ is the central electron density, $\tilde{T}_e(x) = T_e(r)/T_*$, T_* is some arbitrary normalizing temperature, and $\tilde{P}_e(x) = P_e(r)/\bar{P}_e$, where $\bar{P}_e$ is scaled to the average net electron power density in the tokamak. The structure of Eq. (12) leads us to conclude that the central electron temperature must be of the form

$$T_e(0) = T_*\tilde{T}_e(0) = T_* F\left(\frac{\Lambda^2}{2\xi T_*^2}, g, p\right) = \left[\frac{eB\omega_{pe}(0)a^3\bar{P}_e}{c^2 n(0)}\right]^{1/2}\frac{\tilde{F}(g,p)}{\sqrt{2\xi}}, \tag{14}$$

where $\tilde{F}(g,p)$ is a form factor depending on the parameters g and p which are respectively the characteristics of the electron power, and number density profiles. Notice that the last step in Eq. (14) is merely a statement of the fact that the central temperature must not depend on our normalization choice.

Numerical solutions of Eq. (12) confirm Eq. (14), and in addition, reveal the following interesting properties:

1. The normalized temperature profile $\tilde{T}_e(x)$ remains essentially unaltered for physically reasonable density profiles. This is in agreement with the experimental observation already referred to as the decoupling between energy transport and particle transport.[8]

2. For a large variety of power density profiles $\tilde{P}_e(x)$, the variation in $\tilde{T}_e(x)$ is comparatively small. Even for the "edge heating" case, a small amount of ohmic heating power in the central region is sufficient to provide a reasonable $\tilde{T}_e(x)$ profile, thus supporting the so-called "principle of profile consistency."[8] The curves 1 and 2 in Fig. 1 compare the results of the experiments in edge heating[27] in TFTR, and of the results in normal heating. There are exceptions, however: 1) if $\tilde{P}_e(x)$ vanishes completely in the central region (like some stellerator experiments[28]), the temperature profile develops a "flat top" in the central region; 2) if the net electron power density is very small or negative in the edge region, or it is highly concentrated in the central region, a substantial deviation of $\tilde{T}_e(x)$ from its normal shape may and does occur (Curve 3 in Fig. 1).

3. The form factor $\tilde{F}(g,p)$ is generally an insensitive function of g and p; it ranges between 0.9–1.1 for the values g and p which characterize normal discharges. For some unusual discharges, e.g., those with power density highly concentrated in the center, $\tilde{F}(g,p)$ may be well above 1.4 . However, the narrower temperature profile does not mean a substantial gain in electron stored energy.

Before presenting a detailed comparison of the theoretically predicted electron temperature $T_e(r)$ with experiments, we point out several limitations of the present approach:

1. The present model can only cope with the situations where the magnetohydrodynamic (MHD) activity is either absent or has a minor effect on the electron dynamics. For example, in some cases the sawtooth activity does construct a flat-top in the temperature profile within the inversion surface. In such cases, the present theory is valid only in the region outside the inversion surface, i.e., the "confinement region."

2. Since the detailed knowledge of $\tilde{P}_e(x)$ is generally not available, Λ is taken to be a measure of $T_e(0)$ to be compared with the experiment (the parameter ξ is supposed to be machine independent)

$$\Lambda \equiv \left\{ \frac{eB\omega_{pe}(0)a^3\bar{P}_e}{c^2 n(0)} \right\}^{1/2} = 2.44 \left\{ \frac{B_{[T]}(a/R)P_{e[MW]}}{\sqrt{n_{[14]}(0)}} \right\}^{1/2} \text{[kev]}, \qquad (15)$$

where $B_{[T]}$ is the magnetic field B in Tesla, $n_{[14]}$ is the density in 10^{14}cm^{-3}, $P_{e[MW]}$ is the bulk electron power in MW, R is the major radius.

3. The present model [Eq. (1), and thereafter] is not likely to be valid for discharges where the electron distribution function departs substantially from a Maxwellian, for example, in discharges dominated by lower hybrid current drive at low density.

Guided by the aforementioned consideration, we display in Fig. 2, a comparison between the theoretically predicted (characterized by Λ) and experimentally observed [$T_e(0)$] central electron temperature.* The data is primarily collected from published literature. We have considered discharges ranging from 0.3 to 7 KeV in temperature with a variety of heating schemes such as ohmic (OH), ion cyclotron resonance (ICRH), electron cyclotron resonance

*For experiments with sawtooth activity, Λ should be reinterpreted as the value of Eq. (15) multiplied by a factor. This factor is characterized by the ratio of normalized temperature at the inversion surface to that at the center. To determine the ratio use has been made of typical theoretical profiles [for instance, the curve 2 in Fig. 1a] and the experimental value of the inversion surface.

(ECRH), and neutral injection (NI) heating. It is clear that most of the points ($\Lambda - T_e(0)$ plot) fall in a straight narrow band parallel to the $\Lambda = T_e(0)$ line. This allows us to estimate $\xi \sim 0.8$, although some points are located on the upper side of the narrow band. These latter points are known to correspond to the L-mode, and are believed to be plagued by strong magnetohydrodynamic activity. However, even these data points are not far from our predictions. The principal message of the plot is strikingly obvious: all these machines with a large variation in parameters and heating schemes are characterized by a region (the confinement region) where electron dynamics is anomalous and dominant.

In addition to comparing the specific values of Λ and $T_e(0)$, we must also compare the predicted and observed parametric dependence of the central electron temperature, in particular the scaling with the external magnetic field B, the density n, and the total electron power P_e. Before displaying our data, we would like to mention that several empirical and semi-empirical scalings of $T_e(0)$ have been suggested for tokamaks with ohmic[29–32] as well as neutral beam heating.[33] Most of these scalings are similar to the scaling contained in Λ.

For ohmically heated tokamaks with Spitzer-Harm resistivity, the local electron energy confinement is derived to be

$$\tau_{Ee}(0) \sim Z_{\text{eff}}^{-2/7} n^{9/14} B^{1/7} a^{5/7} R^{4/7} q^{4/7}(a). \tag{16}$$

The predicted density dependence of $\tau_{Ee}(0) \sim n^{9/14}$ is compared with the Alcator A data[34] in Fig. 3. The agreement is very good up to $n \sim 8 \times 10^{14}\text{cm}^{-3}$ beyond which the energy loss through the ion channel probably accounts for the deviation.

Because of difficulties involved in a direct measurement of $\bar{P}_e$, we show only one experiment on Doublet III with auxiliary heating[35] to compare our $T_e(0) \sim P_e^{1/2}$ scaling. It is easy to see from Fig. 4 that the theoretical scaling is an excellent fit to the experimental data, where $P_{\text{total}} \sim P_e$ is assumed. From the temperature scaling, the electron energy confinement time is then inferred to be $\tau_{E_e} \sim P_e^{-1/2}$.

Strong additional support for our formula of Eq. (9) has been provided by the recent work of Hiroe et al.[36] who have analyzed, in considerable detail, the energy transport data from TFTR. They report that our expression $D_\perp^T$ does an excellent job in providing understanding of the data.†

The main thrust of this paper has been to examine the salient features of the energy transport in tokamaks, and as a result motivate the setting up of the turbulent transport coefficient $D_\perp^{(T)}$ (Eq. (9)) as a consistent and natural choice for a dominantly magnetic microturbulence. The weak point of this paper lies in our inability to pinpoint the source of microturbulence. Several possibilities exist which might (in future) provide an answer to this question.‡ The agreement of the theory based on $D_\perp^{(T)}$ of Eq. (9) with a wide variety of tokamak experiments, as shown by us and others, does point to the basic soundness of our model. We believe that this newly proposed $D_\perp^{(T)}$ will probably be an essential component of a future, more complete theory of anomalous transport.

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†In Ref. 34 Hiroe et al. adopted a modified version of Eq. (9) (owing to our earlier private communications) in which a multiplier factor $(2\beta_e)^{1/4}$ is introduced, (β_e—the electron beta value) and α/η_e is dropped. The factor $(2\beta_e)^{1/4}$ arises from an attempt to derive $D_\perp$ similar to Eq. (9) on the basis of forced magnetic fluctuations induced by η_e-turbulence, an electrostatic turbulence induced by the η_e-mode,[37] in contrast to the theories by which the η_e-turbulence alone may give rise sufficient contribution to electron energy transport.[20,38,39,40] Although this derivation has not yet been justified in the context of the perturbation theory, we still cannot rule out this scaling for the forced mode (in a nonperturbative sense).[41,42]

‡We together with H. L. Berk have suggested two mechanisms, by which the η_e turbulence may drive the electromagnetic waves in collisionless plasmas. The first mechanism is the forced production of electromagnetic fluctuations with $k_\perp \sim \omega_{pe}/c$ due to the annihilation (or mode coupling) of nonlinear η_e waves.[41–43] The second mechanism is the linear destabilization of electromagnetic waves by nonlinear saturated η_e turbulence, in which η_e turbulence exactly plays the role for suppression of shear damping.[44] Drake, Guzdar, and Hassam have shown in a 2-D fluid simulation that the unstable η_e-mode at the longer wavelength near skin depth ($k_\perp \sim \omega_{pe}/c$) is important to the transport.[38–40,45] Recently, Drake, Guzdar, Hassam,[45] and independently, Bekki, Horton, Hong, Tajima[47] (the latter authors used 3-D shearless fluid simulation) emphasize the importance of the mode coupling processes (for η_e turbulence) in transferring energy into longer wavelength region, and claimed observations of the inverted cascading in their fluid simulations. Drake, Hassam, and Guzdar[46] also claimed that they obtained a formula of transport coefficient in collisionless regime, which is exactly the same one proposed in this paper.

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Figure Captions

1. (a) $\tilde{T}_e(x)$ versus $x = r/a$. Solid curves 1, 2, 3 in (a) are the normalized electron temperature profiles corresponding to the electron power density profiles represented, respectively, by solid curves 1, 2, 3 in Fig. 1(b). The dots with error bars represent the experimental results in TFTR "edge heating,"[27] corresponding to power density profile given by dashed lines in Fig. 1(b).

   (b) Electron power profile $\tilde{P}_e(x)$ versus $x = r/a$.

2. Λ versus $T_e(0)\mid_{\rm exp}$, Λ—the characteristic central electron temperature given by Eq. (15), and $T_e(0)\mid_{\rm exp}$ the experimental result.

3. $\tau_{Ee}(0)$ versus $\hat{n}_e$. Dots represent the experimental results on ALCATOR-A.[34] The dashed line represents $\hat{n}^{9/14}$ curve fitting to the data. ($\hat{n}_e = 6 \times 10^{14}{\rm cm}^{-3}$ is taken as reference point.)

4. $T_e(0)$ versus $P_{\rm total}$. Dots are experimental results on Doublet-III.[35] The dashed line is the theoretical curve $T_e(0) \sim P_{\rm total}^{1/2}$ ($P_{\rm total} = 3MW$ is taken as reference point).

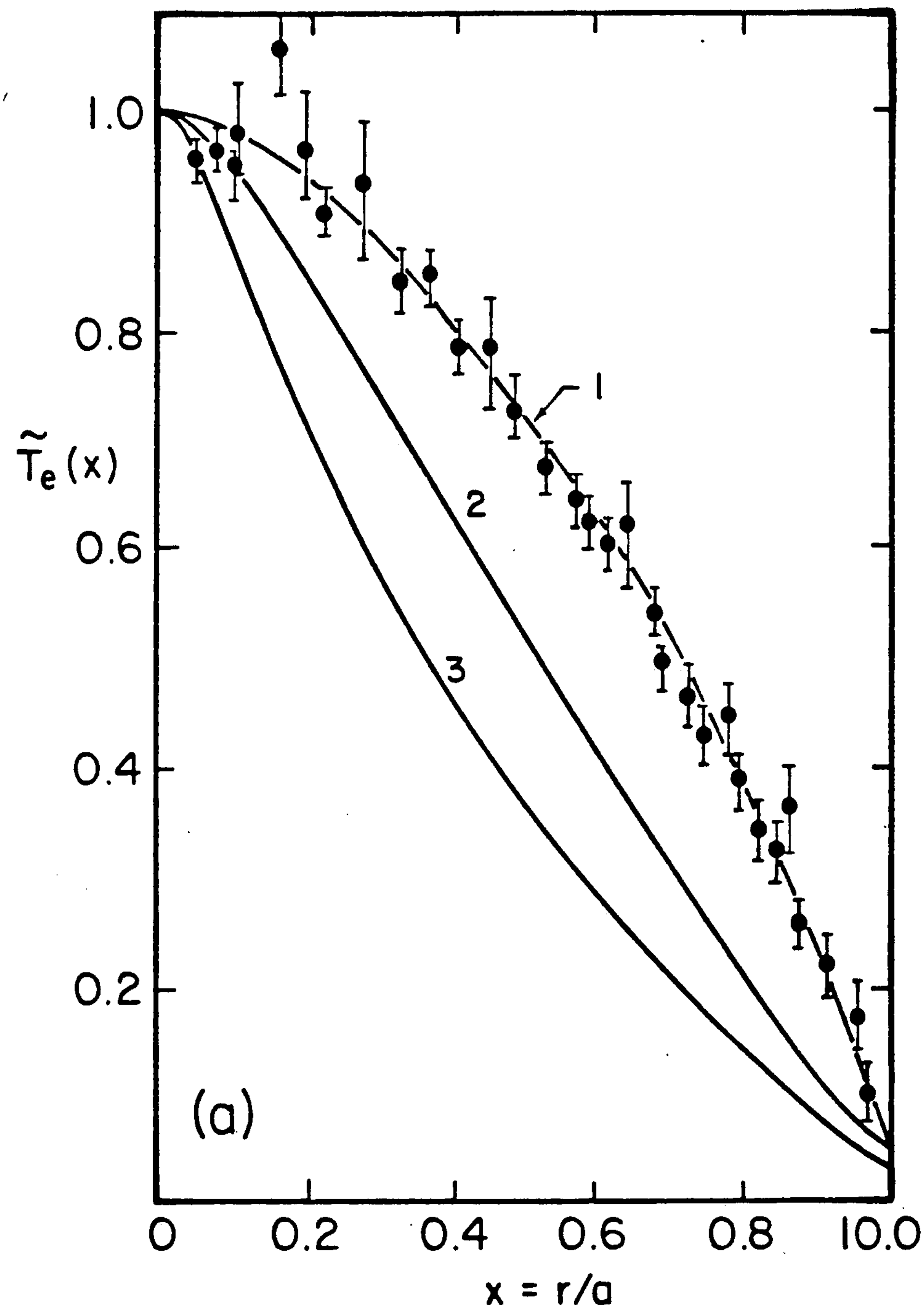


Fig. 1(a)

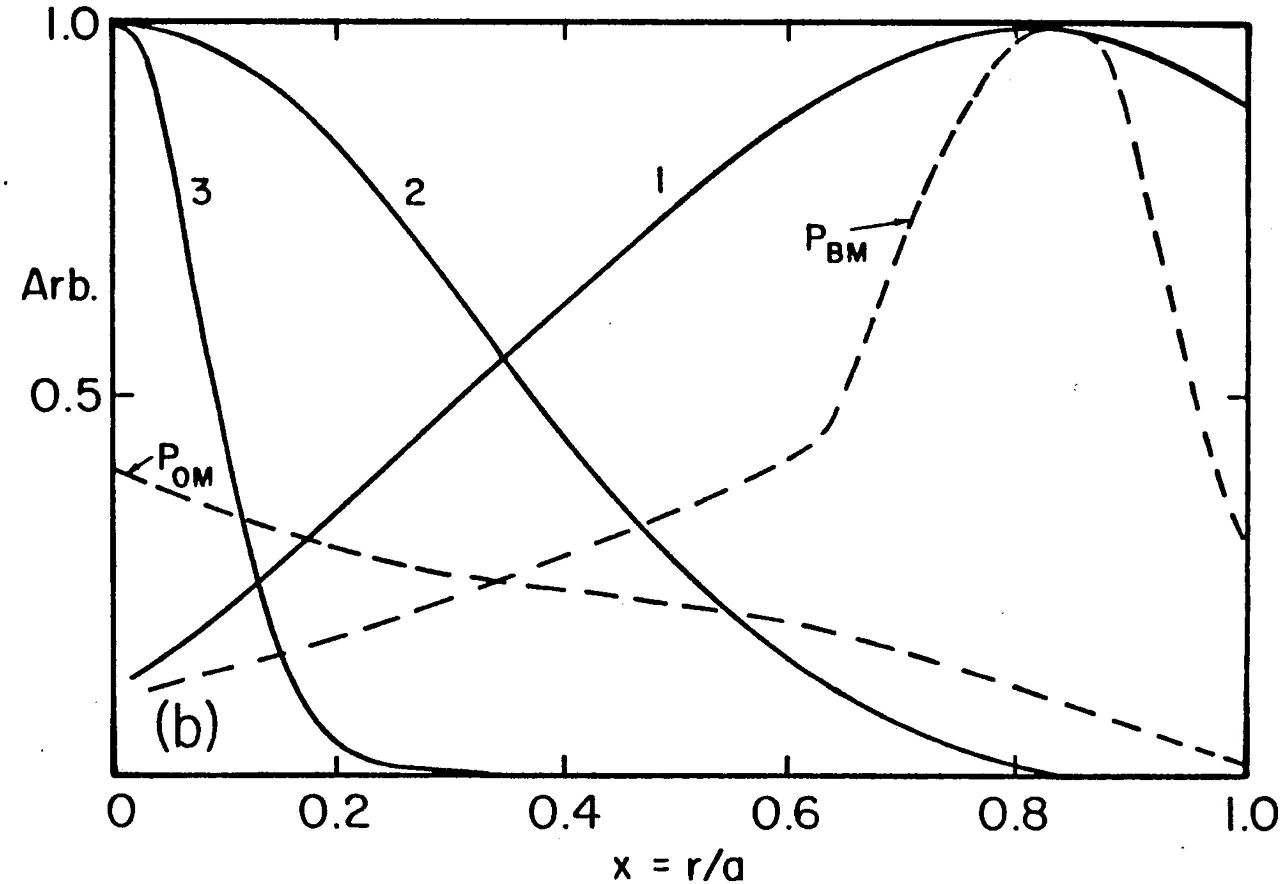
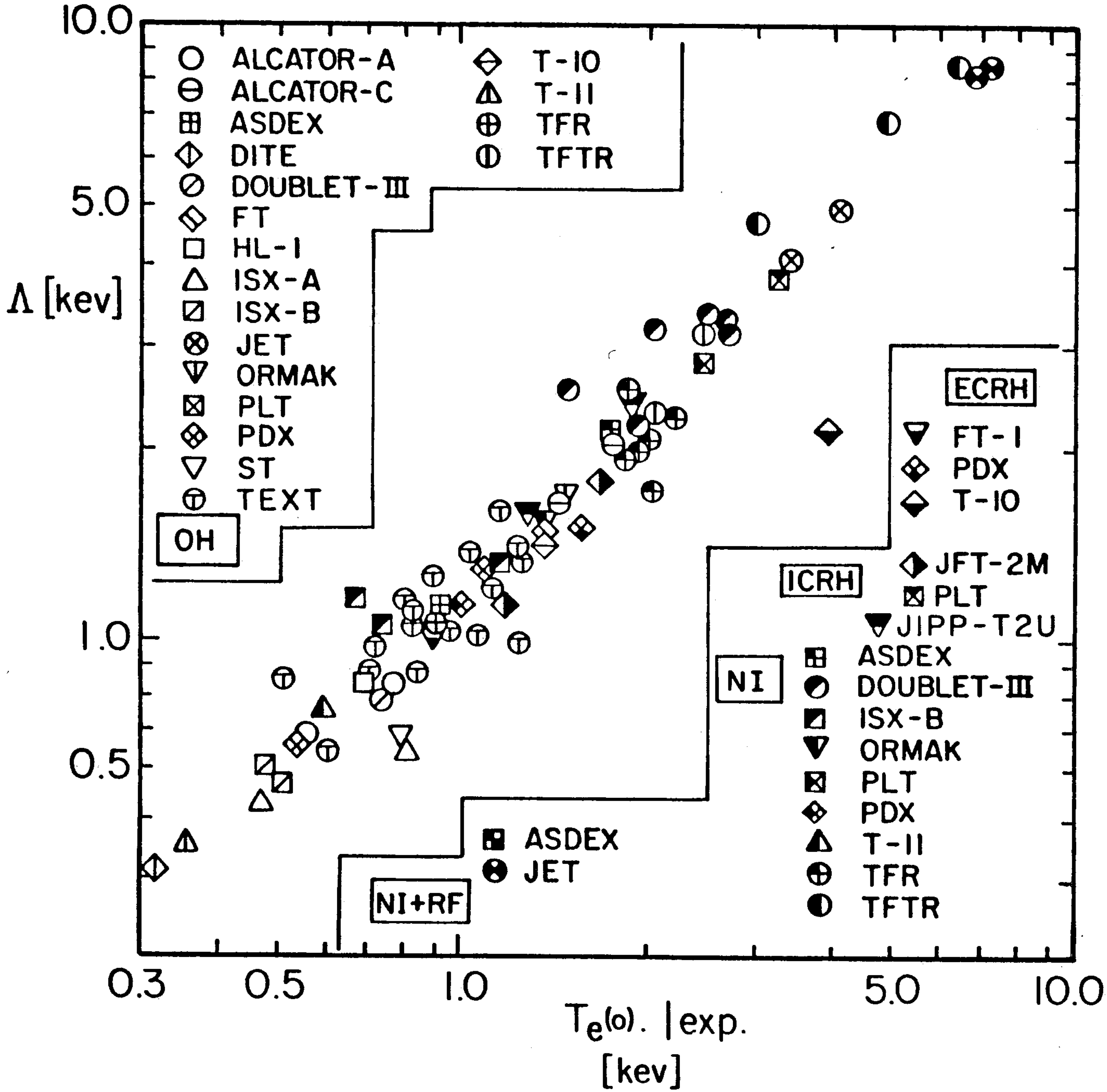


Fig. 1 (b)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Λ [kev]
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ALCATOR-C
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DITE
DOUBLET-III
FT
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ISX-A
ISX-B
JET
ORMAK
PLT
PDX
ST
TEXT
T-10
T-11
TFR
TFTR
OH
ECRH
FT-1
PDX
T-10
ICRH
JFT-2M
PLT
JIPP-T2U
NI
ASDEX
DOUBLET-III
ISX-B
ORMAK
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Fig. 2

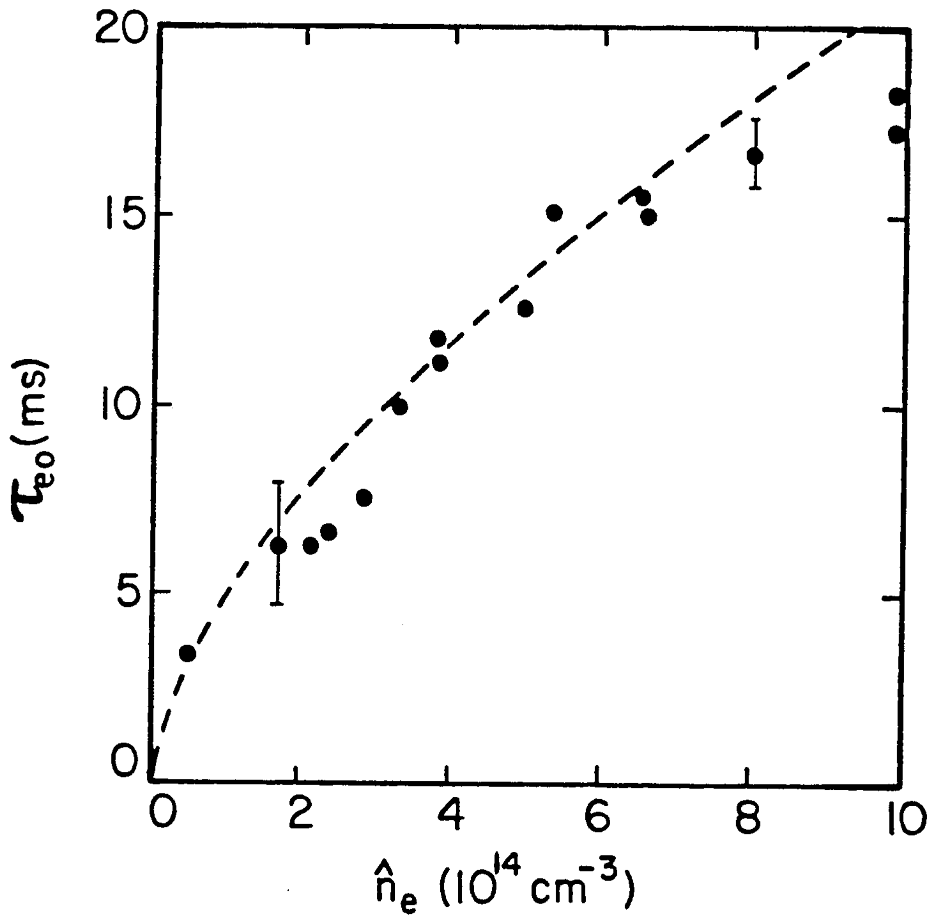

Fig. 3

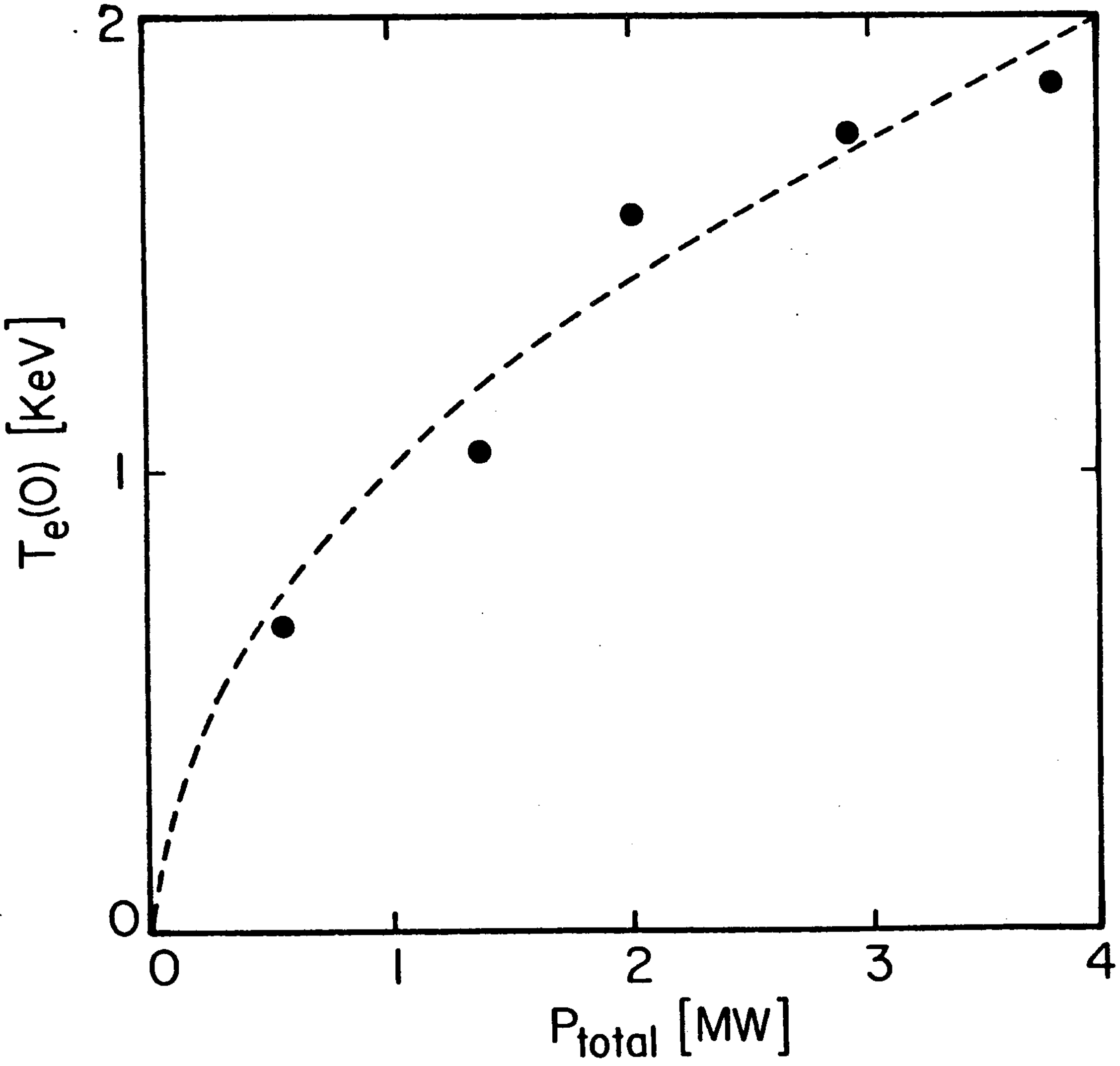

Fig. 4