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An Improved Estimate of Trapped Ion Mode Energy Loss from Tokamak Reactors

J. F. Clarke

MASTER

OAK RIDGE NATIONAL LABORATORY

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Fusion Energy Division

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ENERGY LOSS FROM TOKAMAK REACTORS

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ABSTRACT

An analytic model describing the energy and particle transport in a tokamak reactor dominated by trapped ion mode (TIM) instability is developed. The solution of the transport equations leads to an energy confinement time which can be from 10 to 10^3 times larger than the "standard" τ_{TIM} used in reactor design studies.

1. INTRODUCTION

The empirical scaling law, $n\tau \propto n^2 a^2$, seems to describe the energy confinement of present-day tokamak experiments [1] operating with collisionalities ranging from $\nu^* = 1$ to $\nu^* = 50$. The coefficient of this scaling law is such that a tokamak should be able to achieve ignition in a plasma with a minor radius of roughly 1 m. However, the collisionality of such an ignition tokamak would be on the order of 10^{-4} . In this collisionality range, trapped particle modes [2] are expected to appear and to make a contribution to particle and energy transport. The driving force of both the trapped electron and trapped ion modes is the differential drift velocity between trapped and passing particles produced by density and temperature gradients. As a result, the diffusion coefficients resulting from the existence of these modes turn out to be proportional to local plasma gradients. For trapped ion modes, the particle diffusion coefficient and heat conductivity are

$$D = \chi \propto \frac{r^{5/2} T^2}{R^{5/2} \nu_{ei} B^2 n^2} \left(\frac{dn}{dr} \right)^2 \quad (1)$$

In estimating the $n\tau$ of a tokamak reactor subject to trapped ion modes, it has been customary to replace the density gradient in these diffusion coefficients by $\langle n \rangle / a$ where $\langle n \rangle$ = average electron density and to replace the temperature T by $\langle T \rangle$, the average plasma temperature. The energy containment time is then estimated from $\tau = a^2 / 4D$. The conventional value of $n\tau$ resulting from this approximation [2-3] is

$$\langle n \rangle \tau_{TIM} = 8 \times 10^{-21} a^4 \left(\frac{B}{50} \right)^2 \frac{\langle n \rangle^2}{\langle T \rangle^{7/2}} Z A^{5/2} \left(\frac{1 + \sigma^2}{2} \right) \quad (2)$$

where the units are kiloelectron volts, centimeters, and kilogauss; σ is the elongation of an elliptical plasma of radius a ; and the sharp brackets $\langle \rangle$ denote a volume average.

This expression has become a standard in comparing various tokamak fusion reactor designs. Ignition in a tokamak with a minor radius of 1 m typically requires an $n\tau$ a factor of ten larger than that allowed by Eq. (2).

In large reactor plasmas, the techniques [4] available for fueling lead quite naturally to small density gradients in the plasma core. This will serve to reduce the level of trapped ion mode (TIM) turbulence [5] to below that obtained from the simple estimate. It is the purpose of this paper to show that the approximations leading to Eq. (2) are valid only for simple linear systems. For nonlinear diffusion processes, such as trapped ion turbulence, one must solve the coupled set of nonlinear differential equations describing particle and energy transport with particle and energy sources in order to correctly determine the energy containment time. We show that a more correct estimate of the $n\tau$ resulting from TIM turbulence can exceed that given by Eq. (2) by orders of magnitude. As a result, we conclude that transport due to trapped ion modes will allow ignition in small tokamaks and, in fact, can prove beneficial in providing a convenient method for controlling the reactivity.

2. THE TRAPPED ION MODE DIFFUSION MODEL

In order to obtain an analytically tractable set of equations, let us define a simplified model which describes the particle and energy transport of a tokamak reactor. We take the particle and energy transport

equations to be given by

$$\frac{1}{r} \frac{\partial}{\partial r} r D \frac{\partial n}{\partial r} = -S' \quad (3)$$

$$\frac{1}{r} \frac{\partial}{\partial r} r \left[n_x \frac{\partial T}{\partial r} + \frac{3}{2} T D \frac{\partial n}{\partial r} \right] = -P' \quad (4)$$

where we have added the electron and ion energy transport equations and used $T = T_i = T_e$; S' and P' are the particle and energy source rates per unit volume.

We specify the particle source function, $S' = S'_e x^4$, and the input power per unit volume, $P' = P'_0 [2(1 - H)x^2 + H]$ where $x = r/a$. Of course, the particle source rate per unit volume S' is not arbitrary but is given by $n_g n_e \langle \sigma v \rangle_e$ where n_g and n_e are the neutral and electron density and $\langle \sigma v \rangle_e$ is the ionization rate. Because we determine n_e from the solution of Eqs (3) and (4), choosing S proportional to x^4 constitutes an implicit definition of the neutral particle density within the plasma. This model for S allows the exploration of various fueling conditions ranging from edge gas injection to pellet injection deep in the plasma.

Similarly, choosing the power per unit volume in the form given above constitutes an implicit definition of the profile of externally injected power added to the naturally occurring fusion alpha heating. By varying H from 0 to 2, the total power can be kept constant

$(2\pi^2 a^2 R_0 P'_0)$ while one goes from a condition approximating edge heating by neutral beams to a centrally peaked heating characteristic of fusion alpha particles. This particular choice of source functions simplifies the solution of the equations and is sufficiently general to allow a general estimate of the nonlinear effects flowing from the nature of

the trapped particle diffusion coefficients under a variety of heating and fueling conditions.

Normalizing all variables to their peak values results in

$$\chi = D = \left[\frac{C T_0^{7/2}}{Z_{\text{eff}} B^2 n_0 a^2 A^{5/2}} \right] \frac{x^{5/2} \tilde{T}^{7/2}}{\tilde{n}^3} \left(\frac{\partial \tilde{n}}{\partial x} \right)^2 \quad (5)$$

for the TIM heat conductivity and particle diffusion where $\tilde{T} = T/T_0$ and $\tilde{n} = n/n_0$, and C is a constant [3]. In order to simplify the analysis while preserving the essential nonlinear features of the problem, we approximate the exponents of the normalized radius and temperature by the next lowest integer. This approximation is reasonable since, as we shall see, the product $(xT)^{5/2}$ is within 10% of $(xT)^2$ over the region of interest. Had we chosen the next larger integer, the results would differ from those presented below by only a few percent over the range of interest. This simplification results in our model trapped ion particle diffusion coefficient

$$D = \left[\frac{a^2}{4\tau_0} \right] \frac{x^2 \tilde{T}^3}{\tilde{n}^3} \left(\frac{\partial \tilde{n}}{\partial x} \right)^2 \quad (6)$$

where τ_0 is the conventional TIM energy containment time [Eq. (2)] evaluated at the peak temperature and density.

Notice that the thermal conductivity and particle diffusion due to trapped particle turbulence vanish at the origin. In order to keep the temperature bounded in the vicinity of the origin, we add a small heat conductivity given by $C_0 \tilde{T}/\tilde{n}$ to the trapped ion heat conductivity. The particular form of this residual heat conductivity is chosen for analytic convenience. However, the parametric dependence chosen for this additional heat conductivity is motivated both by the existence of the $m = 1$

sawtooth oscillations [6] at the center of experimental tokamak plasmas, which should lead to a heat conductivity of $\chi \propto \tilde{T}^{5/2}/\tilde{n}$, and by the heat conductivity derived from the empirical scaling mentioned above, $\chi \sim 1/\tilde{n}$. As shown below, we choose the coefficient C_0 such that this additional heat conductivity is only dominant in the vicinity of the origin, and its particular functional form should have little effect on the nature of our results. Thus, our model heat conductivity is given by $\chi = C_0 \tilde{T}/\tilde{n} + D$.

3. THE $\tilde{T}$ PROFILE

Given the above approximations we can solve Eq. (3) directly for the density gradient in terms of the plasma temperature. Since we choose the particle source rate to minimize this density gradient in the region of large $\tilde{T}$ (see Fig. 1), we neglect the convective energy transport term in Eq. (4). This requires $\frac{\partial \ln \tilde{n}}{\partial \ln \tilde{T}} < \frac{2}{3}$, which is the case for our solutions. Then substituting the density gradient into Eq. (4), we can solve for the temperature gradient. Since we wish the residual thermal conductivity to dominate only for a small region ($\chi < \chi_0$) near the origin, we choose

$$C_0 = \left(\frac{2\tau_0 S_e}{3n_0} \right)^{2/3} x_0^4 \quad (7)$$

where x_0 is the radius at which the residual heat conductivity $C_0 \tilde{T}/\tilde{n}$ equals the trapped ion thermal conductivity D . Typically we choose $x_0 = 0.1$. Given this choice of C_0 , we can solve for the temperature

$$\tilde{T} = \left\{ 1 - \frac{2P_0'\tau_0}{x_0^2 n_0 T_0} \left(\frac{3n_0}{2S_e'\tau_0} \right)^{2/3} \left[H \tan^{-1} (x/x_0)^2 + x_0^2 (1 - H) \ln (1 + x^4/x_0^4)^{1/2} \right] \right\}^{1/2} \quad (8)$$

Applying the boundary condition that $\tilde{T} = 0$ at $x = 1$ provides a relationship between the particle and energy inputs and the conventional energy confinement time which can be used to determine T_0 . Taking $\lim x_0 \rightarrow 0$ for the functions inside the brackets in Eq. 8 produces the relationship

$$\left(\frac{S_e'\tau_0}{n_0} \right)^{2/3} = \frac{\pi H}{x_0^2} \frac{P_0'\tau_0}{T_0 n_0} \quad (9)$$

between the particle and energy sources. For x much greater than x_0 and $H \neq 0$, the temperature equation reduces to

$$\tilde{T} = \sqrt{\frac{\pi}{2}} \frac{x_0}{x} \sqrt{1 - x^2} \quad (10)$$

which is sufficiently accurate to obtain the solution to the density equation. Notice that the approximate temperature profile is independent of H as long as the heating profile is not strongly peaked toward the plasma edge ($H = 0$).

4. THE $\tilde{n}$ PROFILE

We now use the latter form of the temperature to determine the particle density. From Eq. (3) we can solve for the logarithm of the plasma density

$$\ln \tilde{n} = - \frac{\sqrt{\pi}}{2^{3/2}} \left[\frac{2\tau_0 S_e'}{3n_0} \right]^{1/3} \left(\frac{\sin^{-1} x - x \sqrt{1-x^2}}{x_0} \right) \quad (11)$$

Evaluating this expression at the surface of the plasma determines the particle source function in terms of the ratio of central plasma density to surface density

$$\left(\frac{2\tau_0 S_e'}{3n_0} \right) = 8x_0^3 \left(\frac{2}{\pi} \right)^{9/2} \ln^3 [1/\tilde{n}(a)] \quad (12)$$

In combination with Eq. (9), this relationship fixes the central temperature in terms of the heat and particle sources and the TIM coefficient.

Figures 1 and 2 show the self-consistent profiles of temperature and density obtained for three values of $\tilde{n}(a)$. Figure 2 shows the collisionality parameter $\tilde{\nu}^*$ consistent with these solutions. For the reference parameters of the ORNL ignition tokamak, TNS [4], the ν_0^* is about 10^{-4} , showing that all of the plasma could be subject to trapped ion instability.

5. THE ENERGY CONFINEMENT TIME

Let us now turn to the determination of the energy confinement time obtained from the self-consistent temperature and density profiles. The definition of energy confinement time is

$$\tau \equiv \frac{3}{2} \frac{\langle nT \rangle}{\langle P' \rangle} \quad (13)$$

In terms of our normalized variables, this can be written as

$$\frac{\langle n \rangle \tau}{\langle n \rangle \tau_0} = \frac{3}{2} \frac{\langle \tilde{n} \tilde{T} \rangle}{\langle P_0' \tau_0 / n_0 T_0 \rangle} \quad (14)$$

In reactor design studies, the TIM energy confinement time given in Eq. (2) has the parametric dependence of τ_0 , but is usually evaluated

at the average temperature and density rather than at the peak values T_0 and n_0 . In terms of our self-consistent profiles, the ratio of the self-consistent energy confinement time to the conventional value τ_{TIM} is found by using Eqs (9) and (12) to evaluate the denominator of Eq. (14),

$$\frac{\langle n \rangle \tau}{\langle n \rangle \tau_{TIM}} = \frac{3}{4} \left(\frac{\pi}{2} \right)^6 \left\{ \frac{\langle \tilde{n} \tilde{T} \rangle}{\langle \tilde{n} \rangle \langle \tilde{T} \rangle} \right\} \frac{H}{\ln^2 [1/\tilde{n}(a)]} \quad (15)$$

Depending on the magnitude of the particle source function or, equivalently, on the density ratio $\tilde{n}(a) = n(a)/n(0)$, one sees that the self-consistent trapped particle mode $n\tau$ can greatly exceed that obtained from the simple linear estimate.

Figure 1 shows three self-consistent density profiles, and Table I gives the volume averages of $\tilde{n}$ and $\tilde{T}$ obtained from these profiles. The enhancement of $n\tau$ over the estimate in Eq. (2) is due to the fact, evident in Fig. 1, that the product $T^{7/2} \frac{\partial n}{\partial x}$ is never large or equal to the simple estimate $\{\langle T \rangle^{7/2} \langle n \rangle\}/a$. The profile for $\tilde{n}(a) = 0.9$ would be typical of that resulting from gas blanket operation with some pellet injection and results in $\langle n \rangle \tau / \langle n \rangle \tau_{TIM} = 2151$. The profile for $\tilde{n}(a) = 0.18$ is typical of operation with an efficient particle-removing divertor and pellet fueling and gives $\langle n \rangle \tau / \langle n \rangle \tau_{TIM} = 9.4$. The intermediate case $\tilde{n}(a) = 0.5$ gives $\langle n \rangle \tau / \langle n \rangle \tau_{TIM} = 53$. Notice that the more peaked heating profiles ($H = 2$) give the maximum $n\tau$, but the dependence is only linear. Note also that the approximations used to arrive at Eq. (15) are not good for values of H much less than one. More exact calculations show that $n\tau$ varies less than a factor of three as H goes from 0 to 2.

6. REACTOR CONTROL BY PARTICLE INJECTION

The $n\tau$ obtained from the flatter self-consistent profile with $\tilde{n}(a) = 0.9$ is so large that other particle transport mechanisms such as that responsible for the presently observed empirical scaling should dominate the particle transport rather than (as is commonly assumed) the TIM turbulence. In the face of these milder loss mechanisms, it should be possible to control the particle source distribution within the plasma so as to maintain density gradients similar to those calculated here and thus to eliminate completely the effect of trapped ion turbulence on particle or energy transport.

On the other hand, the TIM transport has a temperature dependence which allows a stable operating temperature for an ignited fusion reactor [5]. Thus, one can use the profile of injected fuel as a control mechanism if the trapped ion mode is the dominant loss mechanism. Varying the particle source profile will adjust the energy loss rate which in turn adjusts the temperature and consequently the reactor power output.

Neglecting radiation losses, one can write the Lawson criterion as

$$\langle n \rangle \tau = \langle \tilde{T} \rangle \frac{\langle \tilde{n} \rangle \langle \tilde{n} \tilde{T} \rangle}{T_0 \langle \tilde{n}^2 \tilde{T}^2 \rangle} \left\{ \frac{6T^2}{E_\alpha \langle \sigma_F v \rangle} \right\} \quad (16)$$

Now taking advantage of the fact that the bracket has a roughly constant value of 1.8×10^{15} over the range of interest from 8 to 35 keV, one can find the reactor operating point from Eqs (15) and (16) to be

$$\langle T \rangle = 1.9 \times 10^{-14} \left\{ a^4 \left(\frac{B}{50} \right)^2 Z A^{5/2} \right. \\ \left. \frac{(1 + \sigma^2)}{2} \frac{\langle \tilde{n}^2 \tilde{T}^2 \rangle}{\langle \tilde{T} \rangle^2} \frac{H}{\ln^2 [1/\tilde{n}(a)]} \right\}^{2/5} n_0^{4/5} \quad (17)$$

Thus, the reactor temperature can be varied as $\langle T \rangle \propto [n_0 / \ln 1/\tilde{n}(a)]^{4/5}$, and the power output which is given by

$$P = \frac{3}{2} V \frac{E_F}{E_\alpha} \left\{ \frac{E_\alpha \langle \sigma_F v \rangle}{6T^2} \right\} \frac{\langle \tilde{n}^2 \tilde{T}^2 \rangle}{\langle \tilde{T} \rangle^2} \langle T \rangle^2 n_0^2 \quad (18)$$

varies as

$$P \propto n_0^{18/5} / \ln^{8/5} [1/\tilde{n}(a)] \quad (19)$$

For example, changing $\tilde{n}(a)$ from 0.18 to 0.5 at constant density increases the reactor power by about a factor of four. Table I shows the values of the peak and average values of all plasma parameters obtained from the TIM transport model derived in this paper using the reference parameters of TNS.

CONCLUSION

The model derived in this paper allows an analytic solution to the particle and energy transport in a TIM-dominated tokamak reactor. The results show that an ignition condition can be obtained in a small plasma with a radius of roughly 1 m with reasonable edge fueling. The heating profiles and particle source function in the model are chosen for convenience, but are compatible with a wide variety of conditions of alpha heating and gas or low velocity pellet fueling. Thus, the conclusion should have general validity.

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TABLE I
TYPICAL PROPERTIES OF SELF-CONSISTENT TIM PROFILES
IN A TOKAMAK REACTOR

	$\tilde{n}(a)$		
	0.9	0.5	0.18
$\langle \tilde{n} \rangle$	0.97	0.85	0.70
$\langle \tilde{T} \rangle$	0.14	0.14	0.14
$\langle \tilde{n} \tilde{T} \rangle$	0.15	0.14	0.12
$\langle \tilde{n}^2 \tilde{T}^2 \rangle$	0.035	0.033	0.031
$\langle T \rangle$	60	28	18.3
T_0	473	197	127
n_0	0.6×10^{14}	1.5×10^{14}	2.2×10^{14}
$\langle n \rangle \tau / \langle n \rangle \tau_{TIM}$	2151	53	9.4
$P(MW)$	3400	3400	3400

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FIGURE CAPTIONS

Fig. 1. Self-consistent normalized $\tilde{T}$ and $\tilde{n}$ profiles for a plasma dominated by trapped ion energy and particle transport. The points indicate the average value of each curve. Notice that the product $T \, dn/dx$ is small.

Fig. 2. The collisionality parameter ν^* consistent with $\tilde{n}(a) = 0.5$ and $\tilde{n}(a) = 0.18$. The collisionality is normalized to $\nu_0^* = 3.3 \times 10^{-18} q(A/3)^{5/2} (a n_0/T_0)^2$ which is typically 3×10^{-5} for reactor conditions.

Fig. 1.

Fig. 2.

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