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ON IGNITION CONDITIONS

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INFLUENCE OF ANOMALOUS THERMAL LOSSES
ON IGNITION CONDITIONS

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ABSTRACT

In the process of achieving ignition conditions, it is likely that microinstabilities, which lead to anomalous thermal transport of the fusing nuclei, will be present. When such phenomena are taken into account, an appropriate formulation of ignition criteria becomes necessary. In particular, a new type of plasma density limit is identified.

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I. INTRODUCTION

One of the most important parameters characterizing an ignition experiment is the plasma density that can be maintained through the heating cycle. It is clearly advantageous to arrive at the conditions, where heating due to the fusion reaction charged products begins to be significant (i.e., near the minimal ignition temperature, T_{mI}), with the maximum possible density. The minimal ignition temperature can be defined to be the central plasma temperature where fusion heating balances the local losses due to bremsstrahlung and collisional (neoclassical) ion thermal conduction. Under typical conditions $T_{mI} \approx 5$ keV for a deuterium-tritium plasma. However, it is shown in this paper that the presence of microinstabilities, which lead to anomalous thermal transport of the fusing nuclei, introduce a critical density limit that is reached when $T \approx T_{mI}$. Thus, we show another facet of the observation made in previous studies that the anomalous ion thermal conductivity is connected, in practice, with the problem of particle transport. In addition, we give prescriptions for either preventing the excitation of the relevant modes or avoiding their effects by programming the plasma particle density as a function of time in such a way that the density limit is not exceeded when approaching T_{mI} .

In order to illustrate clearly the main points, we analyze a simplified form of the basic homogeneous (zero-dimensional) thermal energy balance equation. The principal heating and loss terms under steady-state ignition conditions are included. The density limit and T_{mI} are evaluated in Sec. II, assuming that the ion losses are collisional (neoclassical). In Sec. III, an analysis of the enhanced ion losses due to ion temperature gradient instabilities is given. Finally, in Sec. IV, the principal consequences of the effects studied are discussed and suggestions are given for optimal ignition scenarios.

II. BASIC ENERGY BALANCE CONSIDERATIONS

Under steady-state conditions, the simplest form of the basic homogeneous (zero-dimensional) thermal energy balance equation is given by

$$P_h = P_{Le} + P_{Li} + P_R - P_\alpha \quad , \quad (1)$$

where P_h is the external heating power density, P_{Le} and P_{Li} represent the losses due to electron and ion thermal conductivity, P_R represents the radiation losses, and P_α is the α -particle heating term. In ohmically heated plasmas, P_{Le} can be expressed as

$$P_{Le} = nT_e \chi_e^A / a^2 \quad (2)$$

with the well-known anomalous diffusivity scaling as $\chi_e^A \propto 1/n$ and with a denoting the radius of the plasma column. Except for the class of high density discharges discussed in the next section of this paper, the ion losses in ohmic cases have usually been treated with a collisional (neoclassical) model. This prescription leads to

$$P_{Li} = nT_i \chi_i^{NEO} / a^2 \quad (3)$$

with

$$\chi_i^{NEO} = v_{ii} n_i^2 (v_{*i}, a/R) \quad ,$$

where ρ_i is the ion gyroradius, ν_{*i} is the collisionality parameter for trapped ions, f is the usual neoclassical function,¹ and R is the plasma column major radius. The variation of aB (B being the toroidal magnetic field) is relatively small when comparing different ignition devices (in view of the fact that in order to confine the 3.5 MeV α -particles produced by D-T reactions, a minimum value of the plasma current I_p has to be produced). Given the requirements of macroscopic plasma stability, $I_p \propto aB$ in practice, and Eq. (3) can be simplified to the form

$$P_{Li} = \gamma_L^i n^2 T_i^{1/2} \quad (4)$$

with γ_L^i absorbing all the other relatively weak parametric dependences.

If we neglect synchrotron emission and take, for simplicity, $T_e = T_i = T$, then Eqs. (2) and (4) can be combined with Eq. (1) to give

$$P_0(T) \equiv P_h - \gamma_L^e \frac{T_e}{a^2} = n^2 T^{1/2} (\gamma_B + \gamma_L^i) - \gamma_\alpha n^2 F(T) \quad , \quad (5)$$

where $\gamma_L^e \equiv n \chi_e^A$, $\gamma_B n^2 T^{1/2}$ represents the loss by bremsstrahlung emission, and $\gamma_\alpha n^2 F(T)$ is the α -particle heating term. Equation (5) shows clearly that there is a density limit which is a function of temperature. This can be written as

$$n^2 = \frac{G(T)}{1 - S(T)} \quad , \quad (6)$$

where

$$G(T) \equiv P_0(T) / [(\gamma_B + \gamma_L^i) T^{1/2}]$$

and

$$S(T) \equiv [\gamma_\alpha / \gamma_B + \gamma_L^i] F(T) / T^{1/2} .$$

We define the minimal ignition temperature, T_{mI} , to correspond to the balance between the α -particle heating and the bremsstrahlung and collisional ion losses. Thus,

$$S(\hat{T} = 1) = 1 \quad (7)$$

with $\hat{T} \equiv T/T_{mI}$. Note that the minimal density is reached at the temperature, $\hat{T}_c$, where

$$\frac{dS}{d\hat{T}} = - \frac{1}{G} \frac{dG}{d\hat{T}} (1-S) \quad (8)$$

From the definition for $G(T)$ following Eq. (6) we obtain

$$\frac{1}{G} \frac{dG}{d\hat{T}} = - \frac{1}{2\hat{T}} \left(1 - \frac{2\hat{T}}{P_o} \frac{dP_o}{d\hat{T}} \right) \quad (9)$$

In the simple limit where dP_h/dT and the anomalous electron loss can be neglected, $(dG/d\hat{T})/G = -1/(2\hat{T})$, and Eq. (8) reduces to

$$\frac{dS}{d\hat{T}} = \frac{1-S}{2\hat{T}} \quad (10)$$

Since $F(T) \propto \langle \sigma_F v \rangle$, where σ_F is the D-T fusion cross section, we may take $F(T) \propto T^3$ to approximate the real temperature dependence on T of $\langle \sigma_F v \rangle$ for temperatures lower than 6 keV. Then $S = \hat{T}^{5/2}$ and Eq. (10) yields

$$\hat{T}_c = (1/6)^{2/5} \approx 0.488 \quad (11)$$

Hence, for $T_{m1} = 5$ keV, $T_c = 2.5$ keV indicating that the effect of α -particle heating becomes important at relatively low temperatures. Notice that, in general, $\hat{T}_c^{5/2} = (1/6) (1+2H)/(1+H/3)$ for $H \equiv - (2T/P_0) dP_0/dT$.

It is significant to note that when approaching ignition conditions, the optimal value of the density to be chosen for a given heating cycle is close to the density limit. In fact, determining the optimal density involves a compromise between the need to maximize the rate of α -particle heating and the need to reach temperature regimes where this can occur within a sufficiently short time. The latter requirement must be compatible with the design characteristics (e.g., duration of the confining magnetic field) of the experimental device considered.

We observe also that in experiments with low values of n_{TE} , τ_E being the total energy replacement time, the maximum density has been found to be related to the plasma current density when only ohmic heating is present. To illustrate this point, note that if only bremsstrahlung losses are considered, then in the central part of the plasma column,

$$\eta_{||} J_{||}^2 \approx \gamma_E n^2 T \quad (12)$$

where $\eta_{||}$ is the longitudinal resistivity. Consequently,

$$n_M = J_{||} \left(\frac{\eta_{||}}{T} \right)^{1/2} \approx \frac{B\chi}{R} \left(\frac{\eta_{||}}{T} \right)^{1/2} \quad (13)$$

where χ is the magnetic twist parameter (so-called rotational transform) and, to the extent that $\eta_{||}$ is classical, $T \propto (T/\eta_{||})^{1/2}$. It should also be

remembered that the average current density $\langle J_{\parallel} \rangle$ in a toroidal configuration is constrained by stability considerations to be proportional to B/R . These points were initially noted in 1969 during the design phase of the Alcator experiment.² Subsequently, from a direct experimental perspective, Murakami³ found, after analyzing a relatively large number of plasma discharges in different devices, that there was evidence for a density limit of the form

$$\bar{n}_{Mu} \propto \frac{B}{R}, \quad (14)$$

where $\bar{n}$ is the line average density. It can be argued that this is consistent with an energy balance equation of the form

$$P_h = P_{Le}, \quad (15)$$

where the energy loss by anomalous electron thermal conductivity prevails, and with the conjecture that

$$P_{Le} \approx P_R|_{\text{edge}}, \quad (16)$$

where the radiation loss at the edge of the plasma column can be written as

$$P_R|_{\text{edge}} = \gamma_R n^2 g(T)|_{\text{edge}}$$

and

$$\bar{n} \propto \langle J_{\parallel} \rangle \times \frac{\langle n_{\parallel} J_{\parallel} \rangle^2 \bar{n}^2}{\langle J_{\parallel} \rangle^2 \times n^2 g(T)|_{\text{edge}}}^{1/2} \quad (17)$$

Hence, to the extent that the term multiplying $\langle J_{\parallel} \rangle$ on the right side of Eq. (18) can be considered nearly constant in practice, we have

$$\bar{n}_M \propto \frac{B}{R} \lambda(r=a) \quad , \quad (18)$$

where $\lambda(r=a) = 1/q_a$ and q is the so-called safety factor of the plasma column. In fact, this is the type of density limit that is frequently considered when assessing the regimes of operation of currently proposed ignition experiments.

Recent detailed studies by Perkins and Hulse⁴ have shown that in the case where the heavy impurities are predominant, the relevant density limit scales as B/R , i.e., in conformity with the Murakami limit of Eq. (14). However, for light impurities (i.e., low Z_{eff} plasmas), it does not appear from the numerical results displayed in Ref. 4, that the temperature dependence of the density limit can really be ignored.

It is important to note that experiments have now attained record high values of $n\tau_E$ by increasing n and, therefore, depressing the anomalous electron thermal conductivity. In these cases the collisional ion thermal conductivity in the main body of the plasma column becomes important, and the temperature-dependent density limit discussed in the present section should be considered instead of the Murakami limit. In the next section we will show that the density limits can be further dramatically modified if anomalous thermal losses of the fusing nuclei are present.

III. CONSEQUENCES OF ANOMALOUS THERMAL TRANSPORT OF FUSING NUCLEI

In the preceding section the ignition requirements and associated density limits were calculated assuming that the thermal confinement

properties of the fusing nuclei followed the predictions of the usual collisional transport theory. However, there is increasing experimental evidence that even under operating conditions characteristic of many present experimental facilities, the ion confinement can be anomalous. For example, recent direct ion temperature profile measurements on the neutral-beam-heated Doublet-III experiment have shown that the associated ion thermal diffusivity, $\chi_i(r)$ is not only larger than the usual neoclassical estimate but also exhibits very different profiles.⁵ These observations are consistent with those made in many high density, ohmically heated discharges involving refueling by gas puffing. Here the saturation of the thermal confinement time, τ_E , was seen to occur at considerably lower densities than predicted by the neoclassical theory.^{6,7} This trend has, in fact, been attributed to the excitation of relevant microinstabilities.⁸ Hence, it becomes important to reassess ignition criteria with anomalous ion thermal transport properties taken into account.

The class of instabilities most likely to be responsible for the anomalous effects of interest here are the ion-temperature-gradient drift modes.⁹ These can be excited when the profile parameter, $\eta_i \equiv d \ln T_i / d \ln n_i$, is sufficiently large [i.e., $\eta_i > (\eta_i)_c \approx 1$ to 2]. This criterion can be readily satisfied in beam-heated plasmas when the ion temperature profiles are relatively steep compared to the density profiles. For high density ohmic discharges refueled by gas injection (puffing), the characteristic (relatively flat) density profiles can also lead to sufficiently large values of η_i . Systematic studies involving nonlinear analytic derivations¹⁰⁻¹² and gyrokinetic particle simulations¹³ have indicated that the basic scaling of the thermal diffusivity for η_i -driven modes follows:

$$x_i \propto T^{3/2} / (B^2 L \tau^{1/2}) \quad (19)$$

with $\tau = T_e/T_i$, and L being a characteristic equilibrium scale length. Although the specific magnitude of x_i and the particular prescription for L are different in the various evaluations made,¹⁰⁻¹² the dominant scaling trends all conform to Eq. (19). Applications of this type of prescription for x_i in recently proposed anomalous thermal transport models¹⁴⁻¹⁶ have yielded results which correlate well with experimental trends in numerous experiments.

As just noted, the basic scaling given by Eq. (19) can be obtained from a variety of different theoretical derivations. At the simplest level, the usual ambient-gradient or mixing-length approximations for the saturated mode amplitude lead to the familiar $\gamma/k_\perp^2$ type estimate for x_i with γ and $k_\perp$ being respectively the growth rate and wave number of the instability. Recall that

$$\gamma \approx (k_\parallel^2 c_s^2 \omega_{*i}^T)^{1/3},$$

where $k_\parallel$ is the parallel wave number, c_s is the ion sound speed, and

$$\omega_{*i}^T = k_\perp c_{Ti} / eBL_{Ti}$$

with L_{Ti} being the ion-temperature-gradient scale length. Taking $k_\parallel \approx 1/qR$ and noting that the worst allowable $k_\perp$ (corresponding to the largest x_i) can be estimated from $\omega_{*i}^T = k_\parallel c_s$ then leads to the following upper-bound approximation

$$x_i \approx \frac{\rho_i^2 v_{Ti} qR}{\tau^{1/2} L_{Ti}^2} \approx \frac{T_i^{3/2}}{B^2 \tau^{1/2} L} \quad (20)$$

with ρ_i being the ion gyroradius and the scale length L corresponding here to qR/L_{Ti}^2 . However, the actual magnitude is very likely to be considerably smaller due to radial localization of the mode¹⁷ and also to toroidal modifications. With regard to the latter point, first recall that the toroidal η_i -modes^{11,18,19} are characterized by

$$\gamma = (\omega_{*i}^T \omega_{Di})^{1/2}$$

with

$$\omega_{Di} = \omega_{*i}^T L_{Ti}/R$$

Simply taking $k_{\perp}\rho_i$ to be a fixed quantity $\lesssim 1$ then leads to the estimate $L = (L_{Ti} R)^{1/2}$ in Eq. (20). Nevertheless, the basic scaling given by Eq. (19) persists. An analysis of the scaling for x_i resulting from the excitation of the η_i -driven "ubiquitous" mode²⁰ is yet to be carried out.

In view of the preceding considerations, it is appropriate to introduce anomalous effects associated with η_i -modes into the steady-state ignition energy balance equation of the previous section, i.e., Eq. (5). Specifically, referring to Eqs. (1) and (3) we write

$$P_{Li} = \frac{nT_i}{a^2} (x_i^{NEO} + x_i^{cm}) \quad , \quad (21)$$

where x_i^{cm} is related to the excitation of collective modes. In particular, referring to the general properties of the ion-temperature-gradient-driven drift modes, we argue that since the relevant frequency and growth rate are density independent, x_i^{cm} is also density independent. The relevant energy balance equation is, instead of Eq. (5),

$$P_o(T) = n^2 [\gamma_{BL} T^{1/2} - \gamma_a F(T)] + \gamma_{cm}^i n \frac{T^{5/2}}{L}, \quad (22)$$

where

$$\gamma_{BL} \equiv \gamma_B + \gamma_L^i,$$

and

$$\gamma_{cm}^i \equiv \chi_i^{cm} B^2 L T^{3/2}$$

following Eq. (20). Note that this equation, unlike Eq. (5), yields a density limit at $T \sim T_{mI}$.

In order to illustrate graphically the effect we are considering, we rewrite Eq. (22) with dimensionless variables. In particular, we define $\hat{T} \equiv T/T_{mI}$. We then obtain

$$\hat{P}_o(\hat{T}) = \hat{n}^2 [\hat{T}^{1/2} - \hat{F}(\hat{T})] + A \hat{n} \hat{T}^{5/2}, \quad (23)$$

where

$$\hat{P}_o(\hat{T}) \equiv \frac{P_o(T)}{P_o(T = T_{mI})},$$

$$\hat{n} \equiv n \frac{\gamma_{BL} T_{mI}^{1/2}}{P_o(T = T_{mI})},$$

$$A \equiv \frac{\gamma_{cm}^i T_{mI}^{5/2}}{L \gamma_{BL}^{1/2} T_{mI}^{1/4} [P_o(T = T_{mI})]^{1/2}},$$

and

$$\hat{F}(\hat{T}) \equiv \frac{\gamma_a}{\gamma_{BL}} T_{mI}^{1/2} F(T)$$

It is reasonable to assume that $\hat{P}_O(\hat{T})$ is a monotonically decreasing function of T even though, in principle, it is possible to control an externally applied heating system in such a way as to obtain the most desirable dependence of P_O or T .

In order to illustrate qualitatively how the presence of an anomalous ion thermal conductivity affects the density limit discussed in Sec. II, we have represented in Figs. 1, 2, and 3 the solution of Eq. (23) assuming $P_O(T) = \text{constant}$; i.e., $\hat{P}_O(\hat{T}) = 1$. In addition, we have modeled $\hat{F}(\hat{T})$ simply by $\hat{T}^3$. For the situation considered in Fig. 1 there is no anomaly, and, consequently, the density limit is reached below the minimal ignition temperature. For $T \geq T_{mI}$ there is no density limit in this case. In Fig. 2, we represent the intermediate case where $A = 1$; i.e., where the anomalous ion thermal conductivity is comparable to the collisional one. Here we see that the density limit is lower than in the previous case and occurs at $T_c = T_{mI}$. In Fig. 3 we represent the more likely case where the anomalous thermal conductivity is well in excess of the collisional one.

From the results presented we can conclude that the presence of an anomalous ion thermal conductivity introduces a density limit that is monotonically decreasing with temperature up to $T_c > T_{mI}$. Beyond a certain temperature, as illustrated in Figs. 2 and 3, it is instead necessary to operate above a certain density because in such a case, α -particle heating should offset the losses caused by the anomalous ion thermal conductivity. We note, in particular, that the minimum density limit is approximately

$$n_m^{(1)} = \frac{P_o (T_{mI}) L}{\gamma_{cm}^i T_{mI}^{5/2}} \quad (24)$$

and we observe that if this density is maintained past the temperature T_{mI} , the density of fusion power produced scales as

$$P_F \propto \frac{[P_o (T_{mI}) L]^2}{T_{mI}^5 (\gamma_{cm}^i)^2} F(T) \quad (25)$$

This indicates a drastic dependence on T_{mI} and on the magnetic field as $\gamma_{cm}^i \propto (aB)^{-2}$. The consequence is that for a given experiment the conditions under which ignition can be achieved are considerably eased (i) if T_{mI} is decreased or (ii) a stronger magnetic field is applied, or (iii) if the effective heating power, P_o , is increased. One possible way to decrease T_{mI} is to use a spin-polarized D-T plasma,²¹ assuming that the state of plasma polarization may be maintained through the heating phase. More generally, the possibility of increasing P_o depends on the intrinsic machine design parameters in the case of ohmic heating and on the characteristics of the auxiliary heating systems in the case of injected heating.

Finally, we note that a series of numerical simulations of the plasma heating cycle in compact ignition experiments and with different values of anomalous ion thermal conductivity have recently been carried out.²² The results from these 1 1/2-D transport code calculations indicate agreement with our conclusions.

IV. DISCUSSION

In the present paper we have pointed out how the ignition criteria for a magnetically confined plasma have to be reformulated when taking into account

the possibility that collective modes, producing enhanced thermal transport of the fusing nuclei, can be excited. In this case a new type of density limit is found, and the intrinsic relation between particle transport and ion thermal energy transport is illustrated.

In order to avoid the unfavorable consequences of the instabilities, we would suggest that techniques like injection of pellets of moderate size in the central part of the plasma be adopted before the plasma temperature reaches the minimal ignition temperature. This corresponds to the regime where α -particle heating can compensate both the bremsstrahlung and the collisional ion thermal conductivity losses. In this case the desirable situation is to keep the profile parameter, n_1 , sufficiently small, so that the relevant modes are not excited. If only weakly excited, the relevant modes occur at very short wavelengths and tend not to produce large thermal transport. On the other hand, the peak plasma density should not be so high as to choke the heating process. After the critical temperature, T_c , is attained, a less discriminating refueling system can be adopted to offset directly the enhanced ion thermal energy transport by a relatively strong rate of α -particle heating.

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FIGURE CAPTIONS

Fig. 1. Density limit obtained by solving Eq. (23) in the limit where $P_0(T)$ is constant and the anomaly factor A is zero. This case is relevant to the considerations made in Sec. II.

Fig. 2. Density limits for the case where the anomaly factor A , introduced in Eq. (23), is unity. In this case the allowable values of the density are below the curve i) for $T < T_c = T_{m1}$ and above the curve ii) $T > \alpha_2 T_{m1} > T_c$. The latter case is justified by the need to have a strong rate of α -particle heating in order to offset the losses due to the anomalous ion thermal conductivity. The region corresponding to the dark heavy line corresponds to one where there are no density limits.

Fig. 3. Density limits for the case where the anomaly factor A appearing in Eq. (23) equals 2. Notice that, by comparison with the case of Fig. 2, the region where there are no density limits is absent. The comment made on the curves i) and ii) apply to the present case. The transition temperature, which separates the upper density limit from the lower, is evident.

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Fig. 1

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Fig. 2

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Fig. 3

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