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Oak Ridge TNS Program: Low Density Ignition Scenarios Using Injection Heating

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OAK RIDGE TNS PROGRAM: LOW DENSITY IGNITION

SCENARIOS USING INJECTION HEATING

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ABSTRACT

Plasma heating and ignition by neutral injection have been studied using a Monte Carlo neutral injection computer code coupled to a single fluid, one-dimensional (1-D) transport code and a two-dimensional (2-D) flux conserving equilibrium code. It is shown that, by taking advantage of central α heating, profile effects, and flux surface shifts in elongated plasmas, it is possible to ignite a modeled, prototypical reactor plasma using 45-30 MW of 100-150 keV (D^+) neutral beams. To do this, the plasma is started at full bore but low density. The density is then increased by peripheral fueling so that the central core begins to ignite before the time when the neutral beams no longer penetrate to this region. The fusion α particles take over the heating requirements in the core region. Because of the decreasing beam line efficiency with increasing energy, it is found that a nearly constant extracted power of about 85-95 MW is needed for ignition in the range studied. There is thus little economic difference in this energy range. However, higher energy beams around 150 keV imply fewer injectors and perhaps lower impurity production rates during heating to ignition.

1. INTRODUCTION

Neutral beam injection is the most proven method of auxiliary heating of tokamak plasmas and is thus the primary heating method in TFTR and in many prototypical reactor designs. A cursory treatment of this problem indicates that high energy (~ 300 keV) beams are required, forcing the use of negative ion sources. However, a careful treatment which includes α heating, profile effects, and flux surface shifts, allows ignition to be achieved with beam energies well below 200 keV.

To see why there might be undue pessimism over injection energy requirements, a brief historical overview is in order. Because the mechanism for neutral beam attenuation is the ionization of the high energy neutrals as they pass through the plasma, the beam attenuation depends on the integral of the plasma density over the path length of the beam in the plasma. The earliest treatments, such as the one by Sweetman [1], placed this factor in an exponent and assumed that effective plasma heating would require a penetration distance of one e-folding length to reach the magnetic axis.

Later, Rome, Callen, and Clarke [2] included the effect of the fast ion drift orbits. They found that, because the volume of the drift orbits shrinks to zero near the magnetic axis, fewer deposited beam ions are required to attain a given fast ion density. The effect was to increase the acceptable penetration distance to about 4 e-folding lengths for tangential injection.

Finally, Rome, Peng, and Holmes [3] suggested that in a reacting plasma, it would not be necessary to continue heating the central region with the beam after significant α heating has begun, and this

leads to another decrease in the required penetration distance. However, for success, this scenario depends critically on the manner in which the plasma density and temperature are built up.

It is important to realize that a fusion reactor does not have to be (or may not be able to be) started with its full size and density. In particular, the initial plasma conditions may be chosen to allow good beam penetration and centrally peaked heating with moderate energy beams [100-150 keV (D^+)]. Good penetration can be obtained by starting with a low density or small radius plasma, or a combination of the two. With enough neutral beam power, the center of the plasma can be heated to the point where α particle heating exceeds the initial beam heating before the beam is shielded out due to increasing plasma size and/or density. Thus, it may not be necessary to continue to heat the central region and the beams need only be used to heat the new plasma particles introduced by the edge fueling.

In this work, we study the full bore, low density startup scenario proposed in Ref.[3]. Accordingly, we are interested in the interplay among the fueling rate, the beam energy, and the beam current, taking into account the effect of 2-D noncircular plasma flux surfaces and the geometry and energy species of the injected beams.

To explore these scenarios we have employed a one fluid transport model coupled to a 2-D MHD equilibrium solver and a Monte Carlo beam deposition code. It was assumed that the beam particles and α particles are confined and instantaneously deliver all of their energy to the plasma in the vicinity of their birth flux surfaces. Fueling is accomplished by a crude model of pellet injection. The fueling source

is flat in the outer 25 cm and then decays exponentially inward with an e-folding length of 5 cm.

The result of these studies is that for reasonable models of injection, fueling, and transport, it is possible to ignite a prototypical reactor plasma (see Table I) with about 85-95 MW of extracted (unneutralized) beam power in the energy range of 100-150 keV (D^+). This corresponds to 45-30 MW of neutralized beam power into the plasma.

TABLE I

Parameters for Elongated, D-shaped Reactor Employed in This Study

Parameter	R_o	a	Elongation	B (axis)	q(o)	q(a)	I_p
Value	5.0	1.25	1.65	4.3	1.0	4.0	4.0
Units	m	m		T			MA

2. PHYSICS MODELS AND EQUATIONS

For the neutral beam heating scenarios studied here we have modeled the plasma as a sequence of 2-D toroidal MHD equilibria evolving in time according to the requirements of magnetic flux conservation and of 1-D particle and energy balances. Because the resistive skin time of the plasma is large compared to the time scale of the neutral heating scenario [4,5], we consider the case where flux is conserved and the q-profile is unchanged. Since our emphasis is on injection and equilibrium evolution, we use only a single fluid approximation to the particle and energy balances.

The instantaneous equilibrium state of an axisymmetric toroidal plasma is described by the Grad-Shafranov equation:

$$\Delta^*\psi \equiv R \frac{\partial}{\partial R} \left(\frac{1}{R} \frac{\partial \psi}{\partial R} \right) + \frac{\partial^2 \psi}{\partial Z^2} = -4\pi R^2 p'(\psi) - FF'(\psi), \quad (1)$$

in which

$$\psi(R, Z) = RA_\phi \quad (2)$$

is the poloidal magnetic flux function, p is the pressure,

$$F(\psi) = RB_\phi \quad (3)$$

is the toroidal magnetic flux function, and ' denotes differentiation with respect to ψ . In terms of ψ and F the magnetic field is given by

$$\vec{B} = \frac{1}{R} \vec{\nabla} \psi \times \hat{\phi} + \frac{F}{R} \hat{\phi}, \quad (4)$$

where $\hat{\phi}$ is the unit vector in the toroidal direction.

Magnetic flux conservation requires that the safety factor

$$q(\psi) = \frac{1}{4\pi^2} FV' \langle R^{-2} \rangle \quad (5)$$

remain a fixed function of ψ as the plasma evolves in time. The $\langle \rangle$ notation in Eq. (5) denotes an average [6] over surfaces of constant

ψ , while $V(\psi)$ denotes the volume contained within a flux surface ψ . Flux conservation also requires the range of ψ to remain constant in time. That is, $\psi_{\text{edge}} - \psi_0 = \text{constant}$, where ψ_{edge} is the value at the plasma surface and ψ_0 the value at the magnetic axis.

The single fluid particle and energy balance equations can be averaged over each flux surface to obtain the following 1-D equations in which ψ is the independent variable:

$$\frac{\partial n}{\partial t} = \frac{n}{V'} (V' \langle \frac{\partial \psi}{\partial t} \rangle)' + \frac{1}{V'} (n' V' \langle D R^2 B_p^2 \rangle)' + \langle S_{\text{inj}} \rangle + \langle S_{\text{fuel}} \rangle, \quad (6)$$

$$\begin{aligned} \frac{\partial p}{\partial t} = & \frac{5}{3} \frac{p}{V'} (V' \langle \frac{\partial \psi}{\partial t} \rangle)' + \frac{1}{V'^{5/3}} (p V'^{5/3} \frac{n'}{n} \langle D R^2 B_p^2 \rangle)' \\ & + \frac{2}{3} p (\frac{n'}{n} \langle D R^2 B_p^2 \rangle)' + \frac{2}{3} \frac{1}{V'} ((\frac{p}{n})' n V' \langle \chi R^2 B_p^2 \rangle)' \\ & + \frac{2}{3} (\langle \eta J_\phi^2 \rangle + \langle H_\alpha \rangle + \langle H_{\text{inj}} \rangle). \end{aligned} \quad (7)$$

The time derivatives on the left-hand side of Eqs. (6) and (7) are evaluated at constant ψ , and the time derivatives of ψ on the right-hand side of Eqs. (6) and (7) are taken at fixed (R, Z) . In Eqs. (6) and (7), $n(\psi, t)$ is the particle density, $p(\psi, t)$ is the pressure, D is the particle diffusion coefficient, χ is the heat conduction coefficient, B_p^2 is the square of the poloidal $\vec{B}$ field, η is the plasma resistivity, S_{inj} is the injection particle source, S_{fuel} is the fueling particle source, ηJ_ϕ^2 is the ohmic heating source, H_α is the nuclear alpha heating source, and H_{inj} is the injection heating source. The ideal gas law $p = nkT$ is assumed.

Equations (1), (5), (6), and (7) comprise our present approach to the time evolution of MHD equilibria subject to particle and temperature diffusion and to various particle and energy source terms. We now describe the simple models of the transport coefficients and source terms.

The plasma resistivity is chosen to be the Spitzer resistivity. The heat conduction coefficient χ is taken to consist of anomalous electron plus neoclassical ion contributions

$$\chi = (\chi_e + \chi_i)/2 . \quad (8)$$

The anomalous electron term is given by

$$\begin{aligned} \chi_e &= 5 \times 10^{17}/n_e \frac{\text{cm}^2}{\text{s}} \\ &= 1 \times 10^{18}/n \frac{\text{cm}^2}{\text{s}} \end{aligned} \quad (9)$$

where the electron density is in particles/cm³. The formulation for the neoclassical ion conductivity is taken from Hinton and Hazeltine [7]. The particle diffusion coefficient is taken to be

$$\begin{aligned} D &= \chi_e/5 = 1 \times 10^{17}/n_e \frac{\text{cm}^2}{\text{s}} \\ &= 2 \times 10^{17}/n \frac{\text{cm}^2}{\text{s}} . \end{aligned} \quad (10)$$

These diffusion coefficients are in reasonable agreement with experimental observations [8] of low density ohmically heated tokamak discharges for reasonable choices of $p(\psi)$, $n(\psi)$, and $q(\psi)$. These simple diffusion coefficients are considered adequate for our purpose since the physical mechanisms for heat conduction and particle transport in tokamaks are yet to be fully understood [9,10].

We have included D-T burning to provide the α particle heating which is needed to ignite the plasma. To do this, the D and T particle densities are each assumed to be equal to one-quarter the total particle (electrons and ions) density. The depletion of deuterium and tritium and the creation of α particles by the nuclear reactions were ignored because of the small burn-up rate ($<1\%$) and short time scale (~ 10 s). Reaction rate formulas for D-T burning were taken from the fitting formulas of Fowler, Caughlan, and Zimmerman [11]. The particle fueling source is assumed to be uniform to a chosen depth in the plasma beyond which it drops exponentially at a chosen rate. The neutral injection energy and particle source terms are calculated using the modified Monte Carlo injection package, NFREYA [12]. A deuterium beam is used.

3. THE FULL RADIUS, LOW DENSITY SCENARIO

Our primary goal in this study is to determine the applicability of the scenario presented in Ref. [3]. The basics of the scenario are simple: a full bore, low density plasma is established by conventional means in an elongated, D-shaped prototypical reactor having major radius $R_0 = 5$ m, minor radius $a = 1.25$ m, and elongation factor of 1.65 (Table I). For this scenario, we began all cases with an initial

electron density of $8 \times 10^{13} \text{ cm}^{-3}$ at the magnetic axis and a low value of $\bar{\beta}_T = 0.01\%$ where $\bar{\beta}_T$ is the volume averaged toroidal beta. About 30-70 MW of neutral injection power at energies of 100-150 keV (D^+) is used to heat the plasma. One-half and one-third energy beam components are included. Simultaneously the density is increased near the edge by an assumed fuel source such as might arise from pellets and/or gas puffing. We assume a net particle fueling profile which is constant to a depth of 25 cm in the plasma beyond which it drops off exponentially with a scale length of 5 cm. Initially the beams easily penetrate the plasma, and the heating is strongly peaked near the magnetic axis. The rate of density increase is controlled to avoid an inverted density profile (caused by the edge fueling). With enough injection power the center of the plasma begins to ignite before the beam ceases to penetrate. As β is increased the magnetic axis shifts outward, somewhat easing the required penetration (see Fig. 1).

Consider the base case having $E_B = 120 \text{ keV}$ and an extracted current $I_B = 1167 \text{ A}$ (Fig. 2). This current consists of full, one-half and one-third energy components in the ratio 90:9:1. Initially the injection heating is highly peaked in the center at a value of about 0.85 w/cm^3 . The chosen fueling rate increases the central density to a value of $1.2 \times 10^{14} \text{ cm}^{-3}$ in about 3.5 s. This increase begins to shield out the beam, and the corresponding injection power at the center drops to about 0.6 w/cm^3 . However, at 3.5 s the center is sufficiently hot so that the α heating is now 0.7 w/cm^3 . This is large enough to have effectively taken over from the beams. The overall central heating at 3.5 s is greater than the initial value of 0.85 w/cm^3 . This central heating is adequate to maintain the high plasma

temperature in the core because the confinement is proportional to n and has improved.

By 5.8 s, $\bar{\beta}_T$ has increased to 7.5%, and the outward shift of the magnetic axis exceeds 0.3 m. Due to this outward shift the central neutral heating is virtually unchanged in spite of an increase in the central density to $1.6 \times 10^{14} \text{ cm}^{-3}$. As shown in Fig. 3, the plasma has ignited. Figure 3 shows the time evolution of $\bar{\beta}_T$, the average electron density, the average temperature, and the injection and nuclear alpha heating powers for the base case with three different beam shut-off times. If the injection is turned off at $t = 5.8$ s and $\bar{\beta}_T = 7.5\%$, the plasma $\bar{\beta}_T$, temperature, and nuclear heating all continue to increase. This remains true if the beam is turned off at $t = 5.1$ s and $\bar{\beta}_T = 5.4\%$. However, if the injection is turned off at $t = 4.3$ s and $\bar{\beta}_T = 4.0\%$, the plasma $\bar{\beta}_T$, temperature, and nuclear heating all decrease, and ignition is not achieved. The results indicate that, for this case, ignition is achieved when the overall nuclear heating exceeds the overall injection heating.

The effects of injection power and neutral beam energy upon the above scenario were studied by varying the neutral beam power and energy from that used in the base case. All other parameters remained unaltered. The results of this survey are shown in Figs. 4-7. In all cases ignition was tested by turning off the neutral injection and following the subsequent evolution of the plasma. In the ignited cases the beam was turned off when the value of $\bar{\beta}_T$ reached 7%. In the nonignited cases, values of $\bar{\beta}_T$ generally failed to rise above 3%, and the plasma rapidly cooled after injection was turned off.

Figure 4 plots the cases considered as functions of the beam energy and neutral power reaching the plasma. The base case has 120-keV beam energy and 60 MW neutral power. The energies surveyed include 100, 120, and 150 keV, while the neutral power varies from about 25-100 MW. The neutral power required to heat the plasma to ignition increases with decreasing beam energy due to the poorer penetration of low energy neutrals into the plasma. In Fig. 5, the same results are plotted as functions of the extracted (preneutralization) power. Here the trend is somewhat reversed; i.e., slightly more extracted power is required to ignite the plasma using 150-keV neutrals than is required for 120- and 100-keV neutrals. This reversal is due to the decrease in the neutralization efficiency of the extracted ion beam with increasing beam energy. Due to the competing effects of neutralization efficiency and plasma penetration by the beam, the extracted power required for ignition varies by only about 10% (85-95 MW) for neutral energies between 100-150 keV.

Scaling studies of beam heated tokamaks by Scott and Sheffield [13] show that up to ignition the product

$$P_B(E_B)^\delta \approx \text{constant} \quad (11)$$

where P_B is the neutralized beam power, E_B is the beam energy of the full energy species, and δ is a profile related constant of order 1. This is satisfied quite well for our barely ignited cases with $\delta = 1.0$ which corresponds to a rather broad density profile model of $n \sim [1-(r/a)^2]^{1/2}$ and is not too unlike our profiles which are broad and only slightly inverted.

Figures 6 and 7 show the times required to heat the ignited cases to $\bar{\beta}_T = 7\%$ as functions of the neutral power and extracted power, respectively. Except for powers very close to borderline ignition, for which the heating time increases rapidly, the heating time is a slowly decreasing function of power with values between 4 and 10 s for the cases surveyed.

Because $\tau \sim n$, we raised the fueling rate by about 15% to see if the increase in n and τ would offset the decrease in heating resulting from less beam penetration. With all other parameters fixed to the base case the time required to reach 7% $\bar{\beta}_T$ was 5.5 s, a slight decrease. However, the increased fueling resulted in a sizeable inversion of the density profile. Decreasing the fueling rate by 15% achieved $\bar{\beta}_T = 7.3\%$ in $t = 6.7$ s, compared with $\bar{\beta}_T = 7.5\%$ in $t = 5.8$ s for the base case. It therefore appears favorable to fuel as rapidly as possible without creating a sizeable density inversion.

As an alternative to the low density scenario we considered the base case with the initial density doubled. The time required to heat the higher density plasma to 7% $\bar{\beta}_T$ was 6.9 s, so that the additional reduction in penetration more than offset the better confinement provided by the higher density.

The favorable geometric effects of the flux surface shifts were considered by running the base case with transport only. The equilibrium data were not reevaluated during the run, so that the initial equilibrium was used throughout to provide the geometric factors. The resulting case ignited with $\bar{\beta}_T$ reaching 7.7% in 7.0 s is compared with 5.8 s to reach $\bar{\beta}_T = 7.5\%$ in the base case. At these

times the volume averaged density and temperature values ($\bar{n} = 0.90 \times 10^{14} \text{ cm}^{-3}$ and $\bar{T} = 15.3 \text{ keV}$) contrasted with those of the base case ($\bar{n} = 1.14 \times 10^{14} \text{ cm}^{-3}$ and $\bar{T} = 11.5 \text{ keV}$). Furthermore, the density profiles in the nonshifted case are rounded (Fig. 8) rather than the squarish shape with slight inversion of the base case. The lower density aids the ignition of the nonshifted case by permitting the beam easier penetration into the plasma. Even so the final central neutral heating of $<0.4 \text{ w/cm}^3$ for a rounded density profile of $1.4 \times 10^{14} \text{ cm}^{-3}$ central value is less than that of the base case — 0.55 w/cm^3 for a squarish slightly inverted density profile having a central value of $1.6 \times 10^{14} \text{ cm}^{-3}$. We attribute the difference to the favorable effect of the shift of the magnetic axis upon the beam penetration. The differences in the density profiles arise from the absence of the $\langle \frac{\partial \psi}{\partial t} \rangle$ term from the right-hand side of Eq. (6) when the changes in equilibrium are ignored. This "adiabatic" term and the equivalent term in Eq. (7) account for the effects of the shifting flux surfaces on the plasma pressure and density. Ignoring flux surface shifts can, therefore, lead to substantial differences in the evolution of the plasma.

Finally as a test of the effects of energy confinement upon the proposed scenario, a pessimistic model in which χ was doubled, i.e., $\chi_e = 10^{10}/n_e \frac{\text{cm}^2}{\text{s}}$ and $\chi_i = 2 \times \chi_{\text{neoclassical}}$, was considered. Varying only the neutral beam energy and current, the plasma ignition results are summarized in Figs. 9-12, which are analogous to Figs. 4-7 discussed above. Basically the above discussion applies except that the required beam power is approximately doubled — the neutral power required for ignition decreases from 100-70 MW as the beam energy

increases from 100-150 keV, and the associated extracted power increases from 200-225 MW. Once again neutralization efficiency offsets beam penetration to result in a nearly constant extracted beam power requirement in the energy range studied.

4. CONCLUSIONS

This study has shown that beam energies well below 200 keV (D^+) are sufficient to ignite a prototypical reactor plasma using a full bore, low density scenario. To make the scenario succeed, it is necessary to provide enough beam power to locally ignite the central region before the beam is shielded out by the increasing plasma density. Local ignition occurs approximately when the local α heating exceeds the local beam heating. Global ignition occurs when the total α power exceeds the total beam power.

For transport which scales as $\tau_E \sim n$, the product of the beam energy and neutral beam power is constant for the barely ignited cases. This is due to the increase of beam penetration with energy. In contrast, because the neutralization efficiency of the beam decreases with increasing energy, the extracted (preneutralization) beam power is almost constant at ~ 85 -95 MW.

Since the cost of injection systems seems to be proportional to the extracted power (and hence the power supply), there is not too much economic difference among the 100-, 120-, and 150-keV systems based on the present calculations. But, if direct recovery of the unneutralized beam ions is available, it would be better to use the higher energy system since it would require fewer beam lines and have greater overall efficiency.

The results of this study agree with similar work performed by Houlberg, Mense, and Attenberger [14], although our scenario is different from theirs in that we increase the density after initiation of the discharge while they expand the plasma radius. It is an open question which one is easier to implement in practice.

Finally, our one fluid transport model may be pessimistic because of the assumption that $T_i = T_e$. This disallows the possible decoupling of the well-confined ion energy from the poorly confined electron energy. If the fluids are decoupled and $T_i > T_e$, increased fusion rates occur at lower values of β_T . It may also be pessimistic in not including any temperature dependence in the anomalous electron heat conductivity [15]. However, our model may be optimistic because it did not include several loss mechanisms such as radiation and instabilities. More definitive results can be obtained by coupling more complete multifluid transport models to the existing MHD equilibrium and Monte Carlo beam deposition packages.

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FIGURE CAPTIONS

- FIG. 1. Equilibrium contours are shown at several times for a typical case [$E_B = 120$ keV (D^+)]. The outward shift of the plasma for increasing β eases the neutral beam penetration requirements.
- FIG. 2. Electron density, temperature, injection heating, and nuclear alpha heating profiles are shown at several times for the base case [$E_B = 120$ keV (D^+), $I_B = 1167$ A].
- FIG. 3. Time evolution of $\bar{\beta}_T$, the average electron density, the average temperature, and the injection and nuclear alpha heating powers for the base case [$E_B = 120$ keV (D^+), $I_B = 1167$ A] for three different beam turn-off times.
- FIG. 4. Survey of plasma ignition for various neutral beam powers at energies of 100, 120, and 150 keV (D^+).
- FIG. 5. Survey of plasma ignition for various extracted beam powers at energies of 100, 120, and 150 keV (D^+).
- FIG. 6. Time required to heat ignited cases to $\bar{\beta}_T = 7\%$ versus neutral power for 100, 120, and 150 keV (D^+) beam energies.
- FIG. 7. Time required to heat ignited cases to $\bar{\beta}_T = 7\%$ versus extracted power for 100, 120, and 150 keV (D^+) beam energies.
- FIG. 8. Electron density, temperature, injection heating, and nuclear alpha heating profiles are shown at several times for the base case [$E_B = 120$ keV (D^+), $I_B = 1167$ A]. The equilibrium and all geometric quantities have been fixed at their initial values throughout the run so that the evolution is entirely through 1-D transport.
- FIG. 9. Survey of plasma ignition for various neutral beam powers at energies of 100, 120, and 150 keV (D^+) with plasma heat conductivity doubled from the base case.
- FIG. 10. Survey of plasma ignition for various extracted beam powers at energies of 100, 120, and 150 keV (D^+) with plasma heat conductivity doubled from the base case.
- FIG. 11. Time required to heat ignited cases to $\bar{\beta}_T = 7\%$ versus neutral power for 100, 120, and 150 keV (D^+) beam energies with plasma heat conductivity doubled from the base case.
- FIG. 12. Time required to heat ignited cases to $\bar{\beta}_T = 7\%$ versus extracted power for 100, 120, and 150 keV (D^+) beam energies with plasma heat conductivity doubled from the base case.

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Fig. 12

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