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OPTIMIZATION OF STEADY-STATE BEAM-DRIVEN TOKAMAK REACTORS

By

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OPTIMIZATION OF STEADY-STATE BEAM-DRIVEN TOKAMAK REACTORS

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ABSTRACT

Recent developments in neutral beam technology prompt us to reconsider the prospects for steady-state tokamak reactors. A mathematical reactor model is developed which includes the physics of beam-driven currents and reactor power balance, as well as reactor and beam system costs. This model is used to find the plasma temperatures which minimize the reactor cost per unit of net electrical output. The optimum plasma temperatures are nearly independent of β and are roughly twice as high as the optimum temperatures for ignited reactors. If beams of neutral deuterium atoms with near-optimum energies of 1-2 MeV are used to drive the current in an INTOR-sized reactor, then the optimum temperatures are typically $T_e \approx 12-15$ keV and $T_i \approx 17-21$ keV for a wide range of model parameters. Net electrical output rises rapidly with increasing deuterium beam energy for $E_b \lesssim 400$ keV, but rises only slowly above $E_b \sim 1$ MeV. We estimate that beam-driven steady-state reactors could be economically competitive with pulsed-ignition reactors if cyclic-loading problems limit the toroidal magnetic field strength of pulsed reactors to $\lesssim 85\%$ of that allowed in steady-state reactors.

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I. INTRODUCTION

The prospects for development of neutral beams, deuterium, and oxygen with energies of ~ 1 MeV/amu have been brightened by progress in neutral beam technology [1,2,3]. The superior current drive efficiency of such beams would permit the construction of neutral-beam-driven steady-state tokamak reactors which could be competitive with pulsed tokamak reactors.

In pulsed tokamak reactors, the cyclic character of thermal and mechanical loads creates problems which may require expensive solutions [4]. The desire to eliminate these problems has led to the search for an efficient, continuously operable plasma-current driver that could replace the inductive current driver which now precludes steady-state operation. Other benefits of driven systems are the additional opportunities for control of plasma temperature, pressure, and current profiles, and the possible improvement in the beta limit [5] caused by the suprathermal ion population generated by some current drive methods.

Ohkawa first suggested the use of fast ions for steady-state current drive [6], and since then many other methods have been proposed [7]. Beam-driven current experiments on the Culham Levitron and DITE produced results which were consistent with theoretical expectations [8,9]. In unpublished studies, Jassby [10], Rutherford [11], and one of the authors [C. E. Singer] concluded that current-driven reactors based on ~ 200 keV D^0 would have a $Q = P_{\text{fusion}}/P_{\text{beam}}$ of only ~ 3 . This would not produce an economically viable "pure fusion" electrical power reactor.

We have re-examined this problem because the current drive efficiency obtainable with ~ 2 MeV D^0 is much higher than that with 200 keV D^0 . Neutral beams of moderate Z atoms ($Z \lesssim 10$) with energies of ~ 1 MeV/amu also appear feasible [3], so we have investigated their potential for current drive as

well. We do not analyze the beam systems in isolation because the optimum plasma temperatures will differ from the optimum for pulsed operation and may depend on the choice of beam system. An analysis of the integrated system is necessary in order to simultaneously find the best beam and plasma parameters. The goal of this exploratory work is to provide general direction for follow-up studies and, ultimately, for reactor design and the neutral beam development program.

The multiplicity of important parameters in a beam-driven steady-state reactor, as well as the complication of opposing positive and negative effects which some parameters produce in different elements of the system, encouraged us to develop a simplified mathematical model of the reactor. The inaccuracies inherent in such a model are minimized by calibrating the model against results from more complete numerical simulations of the plasma-transport and beam-injection processes. The resulting model provides guidance and a framework for understanding the results of what would otherwise be a bewildering multiparameter search for the optimum reactor configuration.

The detailed physical processes of beam-driven currents are reviewed and analytically approximated in Section II. The results are summarized in a single expression and, in Section III, we find that the optimum beam for the standard INTOR configuration is deuterium with $E_b \sim 2$ MeV. In Section IV, a model for the physical processes and economic features of a reactor is developed. This model is used in Section V to find that the optimum plasma temperatures for an INTOR-sized reactor are $T_e \sim 14$ keV and $T_i \sim 18$ keV. These optimum temperatures are relatively insensitive to changes in many of the model parameters. Finally, conclusions and suggestions for lines of further research are presented.

II. PHYSICS OF BEAM-DRIVEN CURRENT

We develop an analytic expression for the total plasma current driven in a tokamak by an injected fast ion distribution. The expression is calibrated with the results of a Monte-Carlo beam-injection and tokamak-orbit code, IO, which calculates beam deposition and subsequent collisional evolution of the fast-ion trajectories in an axisymmetric tokamak with elliptical cross section. The analytic expression for the beam-driven current is incorporated into a reactor model in Section IV.

We begin with a discussion of the importance of beam penetration and its role in the choice of beam particle energy. The analytic solution of the Fokker-Planck equation is used to obtain an expression for the circulating fast-ion current. The current induced in the plasma is discussed, and an analytic approximation to it is given. All these results are collected, discussed, and simplified into the final zero-dimensional (0-d) expression. The effects of a possible beam-driven Alfvén-wave instability are briefly described.

II.A. Beam penetration

The beam deposition profile plays an important role in determining the current profile and the beam driven current efficiency, η_{drive} . Since the current profile must not be hollow and since η_{drive} is greater in high temperature regions, it is essential that the beam penetrate to the center of the plasma column. This is especially true if we consider the bootstrap current, which has a hollow radial profile and requires a very centrally peaked seed current to produce an acceptable total current profile. The strong dependence of penetration on the beam particle energy, E_b , relates penetration requirements to constraints on E_b . The optimum beam energy

derived below in Section II.B is generally higher than that needed for good penetration, so the important constraint in the present section is obtained by limiting the beam shinethrough.

A quantitative measure of the shinethrough fraction, f_s , is

$$f_s = e^{-\tau}, \quad (1)$$

where τ , the number of ionization mean free paths traversed through the plasma, is given by

$$\tau = 2 \int_0^L n_e(r) \sigma_{\text{eff}}(r) ds. \quad (2)$$

Throughout this work we approximate the plasma geometry with a model based on concentric elliptic toroidal flux surfaces. The major radius of the magnetic axis is denoted by R_0 , the half-width and half-height of the plasma column are denoted by a and b , respectively. In a cylindrical coordinate-frame which is coaxial with the torus centerline, the flux surfaces are defined by

$$(R - R_0)^2 + (Z a/b)^2 = r^2,$$

where the minor radius, r , varies from 0 to a . In this cylindrical coordinate system, the path of integration in Eq. (2) is given by $Z = 0$ and

$$r(s) = |[R_{\text{tang}}^2 + s^2]^{1/2} - R_0|,$$

where R_{tang} is the minimum major radius along the beam path. The optimum

choice is typically $R_{\text{tang}} = R_0 - a/4$. The upper limit of integration in Eq. (2) is

$$L = [(R_0 + a)^2 - R_{\text{tang}}^2]^{1/2}.$$

A handy approximation to Eq. (2) is

$$\begin{aligned} \tau &\approx 2 \bar{n}_e \sigma_{\text{eff}} L \\ &= 2 \bar{n}_e \sigma_{\text{eff}} R_0 [(1 + a/R_0)^2 - (1 - a/4R_0)^2]^{1/2}, \end{aligned}$$

where $\bar{n}_e$ is the line-average density. Adopting a typical aspect ratio of $R_0/a \approx 4$, this is

$$\tau \approx 1.7 \bar{n}_e \sigma_{\text{eff}} R_0. \quad (3)$$

The total effective ionization cross section for D^0 ,

$$\sigma_{\text{eff}} = [n_e \langle \sigma v \rangle_e / v_0 + \sigma_p \sum_i n_i + \sum_I \sigma_{I1} n_I] / n_e,$$

is the sum of electron impact ionization [12], proton impact ionization (Ref. 12) for $E_b < 400$ keV/amu, Ref. 13 for $E_b > 400$ keV/amu), and impurity impact ionization and charge transfer [14,15] contributions (v_0 is the speed of the neutral deuterium). At the energies considered here, charge transfer with hydrogenic ions is negligible. We show σ_{eff} in Fig. 1 for $T_e = 20$ keV and $Z_{\text{eff}} = \sum_i n_i Z_i^2 / n_e = 2$, with Argon as the only impurity. It is apparent that raising the beam energy reduces σ_{eff} and τ , thus increasing the shinethrough.

In a survey study [16] with the IO code, we found that the current driven by D^0 is greatest when $2 \lesssim \tau \lesssim 4$. At higher values of τ , the beam deposition in the plasma core falls due to poor penetration, and, for lower values, the shine-through loss becomes significant.

The penetration of non-hydrogenic beam species occurs only partially in the neutral state, and the complexities of the subsequent orbit penetration do not lend themselves to simple analysis. However, the results of the IO survey can be easily summarized: the optimum energies for penetration in a current drive configuration are 1-2 MeV/amu for INTOR-size tokamaks with $n_e \sim 10^{14} \text{ cm}^{-3}$ [16].

II.B. Fast ion current

After neutral beam injection, collisions with the background plasma cause the circulating fast ions to slow down and change their value of $\xi = v_{\parallel}/v$. We obtain the average current of fast ions from an analytic solution to the Fokker-Planck equation. (According to IO results, the absence of toroidal orbit effects here causes only $\sim 10\%$ errors for $R_{\text{tang}} \approx R_0$.) The power efficiency for the circulating current, η_{circ} , is defined, and the beam energies which optimize η_{circ} are found.

Gaffey's [17] steady-state solution to the fast-ion Fokker-Planck equation which includes the average drag on electrons and ions as well as pitch-angle diffusion is

$$f(v, \xi) = \frac{S \tau_{se}}{v^3 + v_c^3} \sum_{n=C}^{\infty} \frac{2n+1}{2} P_n(\xi_0) P_n(\xi) \left[\frac{v^3}{v_0^3} \frac{v_0^3 + v_c^3}{v^3 + v_c^3} \right]^{n(n+1)Z_2/6} U(v_0 - v),$$

where the ions are created at a delta-function source of strength S (in $\text{cm}^{-3} \text{sec}^{-1}$) at $v = v_0$, $\xi = \xi_0$. In the above, $U(v_0 - v)$ is the Heaviside step

function, P_n are Legendre polynomials, and

$$\begin{aligned} \tau_{se} &= 3A_b m_p v_e^3 / (16 \sqrt{\pi} e^4 Z_b^2 n_e \log \Lambda_e) \\ &\approx (0.37 \text{ sec}) \frac{A_b}{Z_b^2} \left(\frac{T_e}{10 \text{ keV}} \right)^{3/2} \frac{10^{14} \text{ cm}^{-3}}{n_e}; \end{aligned} \quad (4)$$

$$\begin{aligned} v_c &= \left[\frac{3\sqrt{\pi}}{4} \frac{m_e}{m_p} \frac{\Sigma_i \log \Lambda_i}{\log \Lambda_e} \frac{n_i Z_i^2}{n_e A_i} \right]^{1/3} v_e \\ &\approx (4.5 \times 10^8 \text{ cm/sec}) (T_e/10 \text{ keV})^{1/2}; \end{aligned}$$

$$E_c \equiv A_b m_p v_c^2 / 2 \approx (100 \text{ keV}) A_b (T_e/10 \text{ keV}); \quad (5)$$

$$\begin{aligned} Z_2 = 2\tau_{se}/\tau_p &= (\Sigma_i \log \Lambda_i n_i^2 Z_i / A_b) / (\Sigma_i \log \Lambda_i n_i^2 Z_i / A_i) \\ &\approx 12 Z_{eff} / 5 A_b \quad [\text{when } (n_D + n_T)/n_e \gtrsim 0.9]; \end{aligned}$$

$$\begin{aligned} \log \Lambda_e &= \log \{ (T_e / h e) (m_e^2 \pi / n_e)^{1/2} \} - 1/2 \\ &= 17.0 + \log \{ (T_e / 10 \text{ keV}) (10^{14} \text{ cm}^{-3} / n_e)^{1/2} \}; \end{aligned}$$

for $E_b \leq (100 \text{ keV}) A_b (Z_b Z_i)^2$,

$$\begin{aligned} \log \Lambda_i &= \log \{ E_b A_i / (e^3 Z_b Z_i (A_b + A_i)) (T_e / \pi n_e)^{1/2} \} \\ &= 25.4 + \log \{ (E_b / \text{MeV}) A_i / (Z_b Z_i (A_b + A_i)) [(T_e / 10 \text{ KeV}) (10^{14} \text{ cm}^{-3} / n_e)]^{1/2} \}; \end{aligned}$$

for $E_b > (100 \text{ keV}) A_b (Z_b Z_i)^2$,

$$\log \Lambda_i = \log \left\{ \left(\frac{2 A_b m_p v_o}{h e (A_i + A_b)} \right) \left(T_e n_e / n_i \right)^{1/2} \right\} - 1/2$$

$$= 23.7 + \log [A_i / (A_i + A_b)] \left[(T_e / 10 \text{ keV}) (A_b E_b / \text{MeV}) (10^{14} \text{ cm}^{-3} / n_e) \right]^{1/2},$$

where m_p is the proton mass, $v_e = (2T_e/m_e)^{1/2}$, Z_b and A_b are the nuclear charge and mass numbers of the injected beam ions. In calculating the simplified expressions above, the Coulomb logarithms [18] were taken to be $\log \Lambda_e = 17$, and $\log \Lambda_i = 24$, and we assumed $n_D/n_T = 1$.

The toroidal current density of circulating fast ions in a torus of major radius R_0 is (neglecting toroidal orbit effects)

$$j_{\text{circ}} = n_b \langle v_{\parallel} \rangle e Z_b / (2\pi R_0)$$

$$= \frac{e Z_b}{2\pi R_0} \int_0^{\infty} v^3 dv \int_{-1}^1 f(v, \xi) \xi d\xi, \quad (6)$$

where the orthogonality of the P_n has reduced the sum over n to the $n = 1$ term. We define

$$I(x, y) = \frac{1}{x} \left(\frac{1+x^3}{x^3} \right)^y \int_0^x \left(\frac{u^3}{1+u^3} \right)^{y+1} du,$$

(which is a hypergeometric function) and then Eq. (6) reduces to

$$j_{\text{circ}} = \frac{e Z_b S \tau_{se}}{2\pi R_0} v_o \xi_o I(x, y), \quad (7)$$

where the scaled beam velocity parameter is

$$x = v_o / v_c = \sqrt{E_b / E_c},$$

and the pitch-angle scattering parameter is

$$y = Z_2/3 \approx 4 Z_{\text{eff}}/5A_b. \quad (8)$$

The power efficiency, or the circulating current produced per unit of injected power, is

$$\begin{aligned} \eta_{\text{circ}} &= i_{\text{circ}} / (SA_b m_p v_o^2 / 2) \\ &= \frac{e Z_b \tau_{se} \xi_o}{\pi R_o A_b m_p v_c} J(x, y), \end{aligned} \quad (9)$$

where $J(x, y) = I(x, y)/x$, and we note that all the dependence on E_b is now in $J(x, y)$. If $x_{\text{max}}(y)$ is defined as the x which maximizes $J(x, y)$ for a given y , then the beam energy which maximizes η_{circ} is

$$\begin{aligned} E_{\text{opt}} &= x_{\text{max}}^2(y) E_c \\ &\approx x_{\text{max}}^2(y) (100 \text{ keV}) A_b (T_e / 10 \text{ keV}). \end{aligned} \quad (10)$$

In Figure 2, we show $J(x, y)$ for several values of y . For D^0 injection into a $Z_{\text{eff}} = 2$ plasma, $y \approx 0.8$; for injection of impurities with $A_b \gg 1$ into $Z_{\text{eff}} \sim 1$ plasmas, $y \approx 0$. Some values for $x_{\text{max}}^2(y)$ are given in Table I together with the range of x^2 for which $J(x, y) \geq 0.90 J(x_{\text{max}}, y)$. This optimal operating region has a considerable dynamic range: typically a factor of 5-

6. Note that $J(x_{\max}, y)$ is reduced by raising y (i.e., higher Z_{eff} and lower A_b).

We have devised an analytic approximation to $J(x, y)$:

$$J(x, y) \sim x^2 / [4 + 3y + x^2 (x + 1.39 + 0.61 y^{0.7})] . \quad (11)$$

For all $0 < y < 4$, this approximation matches $J(x_{\max}, y)$ to within 0.1% and the greatest deviation from $J(x, y)$ is $\sim 6\%$ for $x^2 \sim 0.8$.

II.C. Total driven current

The directed flow of the fast ions tends to drag the electrons with it. This tendency is opposed by electron collisions with the background ions. In the classical description; i.e., when tokamak orbits of the electrons are not considered, this collisional electron current will cancel the ion current if $Z_b = Z_{\text{eff}}$. When trapped-electron orbits are included, the electron current is reduced and Z_b need not differ from Z_{eff} in order to drive a net current.

In the limit where $v_o^2 \ll v_e^2$, the total classical current in the plasma is [6,19]

$$j_{\text{classical}} = j_{\text{circ}} (1 - Z_b / Z_{\text{eff}}) .$$

This is modified by the neoclassical electron-orbit correction to

$$j_{\text{bd}} = j_{\text{circ}} [1 - Z_b / Z_{\text{eff}} (1 - G(Z_{\text{eff}}, \epsilon))], \quad (12)$$

where the trapped-electron correction, $G(Z_{\text{eff}}, \epsilon)$ depends on the inverse aspect ratio $\epsilon = r/R_0$, and Z_{eff} . (We are concerned with only the banana

regime of collisionality, where G is independent of collisionality. The finite aspect ratio calculation [20] of $G(Z_{\text{eff}}, \epsilon)$ is approximated by

$$G(Z_{\text{eff}}, \epsilon) \approx (1.55 + 0.85/Z_{\text{eff}})\sqrt{\epsilon} - (0.20 + 1.55/Z_{\text{eff}})\epsilon. \quad (13)$$

This approximation is in error by less than 1% for $1 < Z_{\text{eff}} < 3$ and $0 < \epsilon <$

0.2. The related collisional current factor

$$F(Z_b, Z_{\text{eff}}, \epsilon) \equiv \frac{j_{bd}}{Z_b j_{\text{circ}}} = \frac{1}{Z_b} - \frac{1}{Z_{\text{eff}}} [1 - G(Z_{\text{eff}}, \epsilon)], \quad (14)$$

is shown in Fig. 3.

When $Z_b = 1$, the current is carried mainly by the fast ions, and the electron current reduces the beam driven current. This reduction can be minimized by using $Z_{\text{eff}} > 1$ and by taking advantage of the neoclassical electron term away from the magnetic axis.

When $Z_b \gg Z_{\text{eff}}$, the current is carried mainly by the electrons and is directed against the fast ion current. Thus, counter-injection of the beam is required. The trapped-electron correction now reduces the total driven current, and the value of $|F|$ is maximized near the magnetic axis where $G(Z_{\text{eff}}, \epsilon)$ is small.

The dependence of F on Z_b is simply additive; therefore, we have shown F for different Z_{eff} and Z_b in Fig. 3 by using a set of shifted vertical axes. Note that when $Z_b/Z_{\text{eff}} \approx 1-2$, the magnitude of F (and, hence, the driven current) is low due to cancellation of terms.

At present, there is uncertainty over the existence of the neoclassical correction to the plasma resistivity in tokamaks [21]; in its absence, we would use the classical result, or $G = 0$ (see the $\epsilon = 0$ axis of Fig. 3). The

removal of the trapped-electron correction lowers $|F|$ for deuterium and raises $|F|$ for impurity beams; as a result, the overall beam-driven current efficiency of impurity beams is comparable to or superior to that of deuterium beams if trapped-electron effects do not occur.

II.D. Global beam-driven current efficiency

By combining Eqs. (7) and (14), we arrive at the final expression for the local beam-driven current density:

$$j_{bd}(r) = \left(\frac{e}{\pi m_p R_o} \right) \left(\frac{H(r) P_b}{V_p} \right) \left(\frac{Z_b^2 \tau_{se}}{A_b V_c} \right) \xi_o J(x, y) F(Z_b, Z_{eff}, \epsilon) . \quad (15)$$

We have used $S = 2H(r)P_b / (V_p A_b m_p v_o^2)$ where V_p is the plasma volume, and the neutral beam deposition function, $H(r)$, is defined in Ref. [22]. Equation (15) is integrated to obtain the global power efficiency

$$\eta_{drive} = I_{bd} / P_b = \frac{2e}{\pi m_p R_o a^2} \int_0^a H(r) \xi_o \left(\frac{Z_b^2 \tau_{se}}{A_b V_c} \right) J(x, y) F(Z_b, Z_{eff}, \epsilon) r dr .$$

This is put into a more comprehensible form by approximating ξ_o as R_{tang}/R_o (negative R_{tang} indicates counter-injection, which is used when $Z_b > Z_{eff}$) and using Eq. (4) to produce

$$\eta_{drive} \approx 0.10 \left(\frac{\text{amps}}{\text{watt}} \right) \left(\frac{5m}{R_o} \right) \times \int_0^a H(r) \left(\frac{R_{tang}}{R_o} \right) \left(\frac{T_e}{10 \text{keV}} \right) \left(\frac{10^{14} \text{cm}^{-3}}{n_e} \right) \left(\frac{J(x, y)}{0.2} \right) F(Z_b, Z_{eff}, \epsilon) (2r dr/a)^2 . \quad (16)$$

Before converting Eq. (16) to a 0-d expression, we discuss the radial profiles which are involved. In order to maximize η_{drive} , the deposition

profile, $H(r)$, should evidently be tailored to emphasize regions of high T_e and $|F|$, and low n_e , while the beam energy should be chosen to maximize $J(x,y)$. In reactor simulations [23], the temperature profile is markedly more centrally peaked than the density profile, so centrally peaked deposition is advantageous; this is especially true for beams with $Z_b \gg Z_{eff}$ where $|F|$ is strongly peaked due to the neoclassical trapped electron correction.

Although maximizing η_{drive} is desirable, there are other important profile effects to be considered. For instance, high- β MHD stability evidently requires a broad current profile [24-27]; this can be easily achieved by flattening $H(r)$, albeit at some cost to η_{drive} . A broad current profile is particularly difficult to achieve with impurity beams if the neoclassical trapped-electron correction is valid because then $|F|$ is small except near the magnetic axis. (There would, however, be little reduction in η_{drive} for deuterium if the beam energy were $\lesssim 500$ keV because then $J(x,y) \propto x^2 \propto T_e^{-1}$, and the integrand of Eq. (16) is nearly independent of T_e .) On the other hand, if the neoclassical bootstrap current [28,29] is generated, it would require a centrally peaked beam-driven seed current to produce a broad current profile.

Having acknowledged the complexities of Eq. (16), we now approximate it by

$$\eta_{drive} \approx A_{bd} \left(\frac{5m}{R_0} \right) (1-f_s) \left(\frac{R_{tang}}{R_0} \right) \left(\frac{T_e}{10 \text{ keV}} \right) \left(\frac{10^{14} \text{ cm}^{-3}}{n_e} \right) \frac{J(x',y)}{0.2} F(Z_b, Z_{eff}, \epsilon = a/2R_0), \quad (17)$$

where

$$x' = (E_b / B_{bd} E_c)^{1/2} = [E_b / (B_{bd} 10 A_b T_e)]^{1/2}. \quad (18)$$

The calibration coefficients, A_{bd} and B_{bd} , are found by setting Eq. (17) equal to the η_{drive} calculated by IO for some reference plasma and beam species. In this way, the main parameter dependencies are explicit and the subtleties involved in the integral of Eq. (16) are retained implicitly in A_{bd} and B_{bd} . The factor of $(1 - f_g)$ is motivated by the observation that the fraction, f_g , of the beam which passes completely through the plasma will drive no current. The volume averaged plasma density and density weighted volume averaged temperature are to be used in Eqs. (17 and 18); this produces compatibility with the 0-d reactor model developed below.

II.E. Beam-driven Alfvén-wave microinstabilities

The non-thermal distribution of injected fast ions may drive microinstabilities in reactor plasmas. In particular, if $v_{||}$ of a fast ion is greater than the Alfvén speed

$$v_A = B_0 / (4\pi m_p \sum_i n_i A_i)^{1/2} ,$$

then the ion can couple its energy to the wave. The beam power required for steady-state reactors is ~ 100 MW; this is sufficient to drive Alfvén-wave microinstabilities according to linear stability estimates [30]. Nonlinear analyses have not yet been done to determine what the consequences would be, but it is possible that the fast ions would be very rapidly slowed to $v \sim v_A$. This would greatly reduce the beam-driven current efficiency, and it would be disadvantageous to use an injection energy greater than

$$E_{AL} = m_p A_b v_A^2 / 2 , \quad (19)$$

which is typically ~ 400 keV for deuterium in a reactor.

III. BEAM OPTIMIZATION

Before enlarging the scope of the discussion to cover the entire reactor system, we analyze Eq. (17) and use the standard INTOR parameter set [31] in a sample optimization of beam particle charge, mass, and energy.

III.A. Beam particle charge

The beam particle charge, Z_b , appears in η_{drive} only as an argument of $F(Z_b, Z_{\text{eff}}, \epsilon)$, which is shown in Fig. 3. (In principle, A_{bd} depends implicitly on Z_b ; we ignore this dependence since A_{bd} is approximately the same for different beam species [16].) The optimum charge is not obvious, however, as Z_b interacts in a complex way with Z_{eff} and the trapped-electron correction, $G(Z_{\text{eff}}, \epsilon)$. To avoid prejudicing the choice of beam species, we assume that Z_{eff} can be controlled through the action of divertors and impurity injection.

If the neoclassical trapped-electron correction is valid, then $|F| \approx 1$ over a sizable volume only if $Z_b = 1$. It is apparently best to choose $Z_{\text{eff}} \gg 1$, but the improvement in $|F|$ as Z_{eff} is raised and opposed by a reduction in $J(x', y)$. We have found that a compromise value of $Z_{\text{eff}} \sim 2$ is best.

If the electrons respond classically, sizable values of $|F|$ can be had with either $Z_b = 1$ and $Z_{\text{eff}} \gtrsim 3$, or $Z_b \gtrsim 8$ and $Z_{\text{eff}} \lesssim 1.5$. In the latter case, $F(Z_b, Z_{\text{eff}}, \epsilon) < 0$, so counter-injection is required ($R_{\text{tang}} < 0$). To make a complete comparison of deuterium versus impurity beams, we must take into account the details of beam penetration, the lower degree of pitch angle scattering of impurity ions, and the dilution of reacting fuel by impurities. All we can say here is that impurities with $Z_b \gtrsim 8$ and deuterium have comparable current drive efficiencies when the neoclassical trapped-

electron correction is absent from Eq. (14).

III.B. Beam particle mass

The dependence of η_{drive} on the particle mass, A_b , is in $J(x',y)$ where $y \propto Z_{\text{eff}}/A_b$. Since the maximum value of $J(x',y)$ increases as y decreases, E_{inj} is improved by raising A_b . Of course, once the beam species has been chosen there is little choice of what A_b is, because for all species of interest, except $Z_b = 1$, we have $A_b \approx 2Z_b$. For $Z_b \gg 1$, the dependence on A_b is negligible since $y \sim 0$. For $Z_b = 1$, however, it makes a significant difference in y , and hence $J(x',y)$, whether $A_b = 1, 2$, or 3 . Hydrogen injection is ruled out by the unfavorable y and dilution of fuel ions. From the perspective of our theoretical considerations tritium is preferred over deuterium. However, this may be outweighed by the practical disadvantage of increased tritium contamination, so we have used deuterium in our optimization.

III.C. Beam particle energy

The dependence of η_{drive} on the beam energy, E_b , is contained in the shinethrough fraction, f_s , and in $J(x',y)$ where $x' = [E_b/B_{\text{bd}}E_c]^{1/2}$. In general, η_{drive} is maximized when $f_s \lesssim 0.1$, and this amounts to a lower bound on the beam energy. The value of $J(x',y)$ is near its maximum for a broad range of E_b (see Table I).

We illustrate these issues with an example based on INTOR. Taking $f_s < 0.1$, we change Eq. (3) into

$$-\log(0.1) < \tau \approx 1.7 \bar{n}_e \sigma_{\text{eff}} R_0.$$

For $\bar{n}_e = 1.3 \times 10^{14}$ [20], we find $\sigma_{\text{eff}} > 2 \times 10^{-17} \text{ cm}^2$ or $E_b < 4.5 \text{ MeV}$. For $Z_{\text{eff}} = 2$, Eqs. (10 and 18) become $E_{\text{opt}} \approx 5.3 B_{\text{bd}} E_c \approx 106 B_{\text{bd}} T_e$. With an INTOR temperature of 10 keV and $B_{\text{bd}} \approx 2$, $J(x', y)$ is optimized with $E_b \sim 2 \text{ MeV}$ and is within ten percent of its maximum in the range $1 \text{ MeV} < E_b < 5 \text{ MeV}$. In this example, nearly the entire range of beam energies which optimize $J(x', y)$ has acceptable shinethrough.

IV. PHYSICS AND ECONOMICS OF REACTORS AND BEAM SYSTEMS

We optimize the reactor plasma temperatures in the next section. As a preliminary, a set of equations is developed in this section to quantify the several relationships between the injected beam, the β limit on the plasma pressure, the plasma power-balance, the fusion power, and the net electrical power produced by the reactor.

We must also identify the measure of performance which is to be optimized. The traditional choice is the ratio of fusion to input power, $Q = P_{\text{fusion}}/P_{\text{beam}}$. However, this is the wrong choice because Q is optimized with very high temperatures, and thus, for any given β , at relatively low density and low fusion power. Since a large part of the cost of the reactor is independent of the fusion power, the cost per unit of electricity produced is relatively high if Q is optimized. The appropriate performance standard for an electrical power reactor is the cost of the produced electricity. Accordingly, we have developed a cost model for beam-driven tokamak reactors.

IV.A. Reactor physics

Several major interdependent features of reactor physics are discussed and reduced to a simple mathematical model. MHD stability considerations determine the required plasma current. At the heart of the model, the plasma

power balance expresses a complex relationship between the beam power, fusion power, the energy confinement, and the equilibrium plasma density and temperature. Finally, the net electrical power depends on both the fusion power and the recirculated power needed for the beams.

IV.A.1. Beta limit and plasma current

When the plasma β becomes too large, MHD oscillations are expected to produce increased transport or loss of equilibrium [24-27]. Since the fusion power and, hence, the net electrical power scale roughly as β^2 , the β limit is a prominent reactor design parameter. There being no definitive determination of the maximum value of β , we treat it as the major unknown variable and express our results as a function of it.

We include only the thermal ion and electron pressures in β . Since the densities and temperatures appearing in our 0-d model are volume averaged and density-weighted volume averaged, respectively, we have

$$\beta = (n_e T_e + (n_D + n_T + n_{He} + n_W) T_i) (8\pi/B_0^2), \quad (20)$$

where B_0 is the nominal toroidal field strength. The pressure due to suprathermal alpha-particles and beam-ions may improve the MHD stability and raise the critical β of the thermal populations [5]. By comparing pulsed and driven reactors at the same thermal β , we assume that the suprathermal ions have no effect on the β limit for the thermal plasma.

The β limit for ballooning modes increases with increasing plasma current. To avoid disruptions, the surface safety factor, $q(a)$, must be $\gtrsim 2$, and this effectively places an upper limit on the current. The INTOR design value is near this limit so we use $I_p = 6.4$ MA [31], in all our calculations.

Perhaps, some of this current could be the neoclassical bootstrap current or current driven passively by synchrotron radiation [32]. The inclusion of these effects would further complicate our study so we do not consider them here. Instead, we require that the entire plasma current be driven by the beams, and the beam power is given by

$$P_b = I_p / \eta_{\text{drive}}. \quad (21)$$

IV.A.2. Power balance

Since there is no a priori reason for the optimum electron and ion temperatures to be equal, we use a two-fluid (electrons and ions) power balance model. The model is restricted to 0-d in order to limit the computational effort involved in the multi-parameter optimization procedure. The electron thermal transport and ion-electron coupling employ simple scalings normalized to BALDUR [33] transport simulations; the normalization constants, denoted by asterisks, are given in Table II, and the calibration procedure is discussed in a later section.

The electron and ion power balance equations are:

$$0 = -U_e / \tau_{Ee} + P_b (1 - h_b) + P_\alpha (1 - h_\alpha) - Q_{ie}, \quad (22)$$

$$0 = -U_i / \tau_{Ei} + P_b h_b + P_\alpha h_\alpha + Q_{ie}, \quad (23)$$

where the terms not previously defined are the thermal electron and ion energies, $U_e = 3V_p n_e T_e / 2$ and $U_i = 3V_p n_i T_i / 2$; the plasma volume, $V_p = 2\pi^2 a b R_0$; the electron and ion energy confinement times, τ_{Ee} and τ_{Ei} ; the ion heating fraction for beams and alpha particles, h_b and h_α ; and the ion-electron temperature equilibration power, Q_{ie} .

To maintain thermal equilibrium in the strongly heated plasmas which we will consider, controllable loss channels must be added to the usual anomalous electron and neoclassical ion transport. These losses could be due to impurity radiation, ripple transport, or enhanced transport from unstable MHD oscillations. Accordingly, we define τ_{Ee} by

$$1/\tau_{Ee} = 1/\tau_{INTOR} + 1/\tau_{Eec}$$

where $\tau_{INTOR} = \tau_{INTOR}^* (n_e/n_e^*)$ exhibits the INTOR scaling of the anomalous electron thermal transport and τ_{Eec} is adjusted to achieve electron power balance. The neoclassical ion thermal transport is negligible in the present context, so we use

$$\tau_{Ei} = \tau_{Eic},$$

which is adjusted to achieve ion power balance.

We have given special attention to the 0-d approximation of the fusion power since it is a large term in the power balance. The total power in a torus with concentric elliptical flux surfaces is

$$P_\alpha = E_\alpha V_p \int_0^1 n_p(u) n_T(u) \langle \sigma v(T_1(u)) \rangle 2u \, du,$$

where E_α is 3.5 MeV and $u = r/a$ is the scaled radius. If the density and temperature profile shapes are fixed, this can be expressed as

$$P_\alpha = E_\alpha V_p n_p n_{T\alpha} f(T_1), \quad (24)$$

where only the scaling with T_i is not transparent. For $5 \text{ keV} \lesssim T_i \lesssim 60 \text{ keV}$, $f_\alpha(T_i)$ can be fit very well by

$$f_\alpha(T_i) = A_\alpha \exp\{-B_\alpha [\log(T_\alpha/T)]^2\}, \quad (25)$$

where A_α , B_α , and T_α are weakly dependent on the density and temperature profiles.

A fast ion divides its initial energy among the background electrons and ions as it thermalizes. The fraction going to ions is [34]:

$$h(x) = (2/x^2) \{ [\arctan^{-1}((2x-1)/\sqrt{3}) + \pi/6]/\sqrt{3} - (1/6) \log \frac{(1+x^2)}{(1-x+x^2)} \}, \quad (26)$$

where $x^2 = E_0/E_c$, the ratio of the initial and critical energies (the critical energy is defined in Section II). The fractions for beams, $h_b = h(x_b = [E_b/E_c]^{1/2})$, and alphas, $h_\alpha = h(x_\alpha = [E_\alpha/E_c]^{1/2})$, are generally not equal since both the initial energy and the critical energy are usually different.

The electron-ion equilibration power, Q_{ie} , is the most difficult term to model in 0-d since it depends on the relatively small difference between the temperatures, and the sign is frequently different in the core and outer regions of the plasma. Its main function in our model is to restrain the difference between the electron and ion temperatures. We use the simple local scaling of Q_{ie} [35] normalized to a BALDUR simulation:

$$Q_{ie} = Q_{ie}^* (n_e/n_e^*)^2 (T_e^*/T_e)^{1.5} [(T_e - T_i)/(T_e^* - T_i^*)]. \quad (27)$$

IV.A. Calibration of 0-d reactor model

We used a BALDUR simulation of the standard INTOR configuration to

calibrate the 0-d model described in Section IV.B. The simulation code and transport model we used are described in Ref. [23]. Tangential injection of 70 MW of 2 MeV deuterium was used to heat a D-T plasma which was seeded with tungsten to produce $Z_{\text{eff}} \approx 2$. The simulation pulse length of 10 sec was long enough to permit the densities and temperatures to relax to their steady-state profiles.

The calibration density and temperatures, together with the normalization constants are shown in Table II. The global electron energy confinement time is used for τ_{INTOR}^* , and the maximum value of the radially integrated electron-ion equilibration power is used for Q_{ie}^* . The density profile and scaled ion temperature profiles were used to calculate the integrated fusion power as a function of T_i ; the results were well approximated by the expression in Eq. (25) with the fitting parameters given in Table II. In calculating A_{bd} , we used $f_s = 0.031$ from the IO calibration, and $F(Z_b = 1, Z_{\text{eff}} = 2, \epsilon = 0.5a/R_0) = 0.77$.

IV.B. Reactor economics

In the following section, we optimize the capital cost per unit of net electrical output capacity of several reactor configurations. This is only one component needed to compute the actual busbar cost of electricity; the others are the interest rate, availability, operating costs, and plant lifetime. The interest rate is independent of the choice of beam or reactor type (steady state vs. pulsed) and the availability of a steady-state tokamak is presumably at least as high as that of a pulsed tokamak. The changes in the busbar cost due to the differences in plant lifetimes and operating costs in the various cases are probably small. For these reasons, we feel that the capital cost rankings will be close to the busbar cost rankings.

Since our ultimate goal is to compare pulsed and steady-state fusion reactors only to each other (and not to, say, a fission reactor), we need only the relative costs of tokamaks and beams and not absolute costs. The beam costs are quite uncertain so there is little to be gained by discussing the relatively smaller uncertainty in tokamak costs.

IV.B.1. Net electrical power

The electrical product of a reactor is generated from the thermal energy collected from the first wall and blanket. Taking into account the multiplication of the fusion neutron power by exothermic nuclear reactions in the blanket, the total thermal power is

$$P_{th} = P_b + P_\alpha + 4M_b P_\alpha, \quad (28)$$

where M_b is the neutron power multiplication factor in the blanket. The gross electrical power is

$$P_{gross} = \eta_{th} P_{th},$$

where the electricity needed to power the equipment in all the generating systems is presumed to be proportional to P_{th} and has been accounted for in calculating the thermal electrical efficiency, η_{th} .

We have divided the remaining local electrical consumption into two categories. The largest is the power needed to run the beam systems. The other category is the power for the normal-conducting magnetic field coils, cooling, cryogenic refrigeration, and vacuum pumping systems, all of which we presume will scale with the size of the tokamak. The net electrical power

capacity is thus

$$P_e = \eta_{th} P_{th} - P_b / \eta_b - P_E V_E, \quad (29)$$

where the overall electrical to beam power efficiency, η_b , the coefficient P_E , and the tokamak engineered volume, V_E , are discussed below.

The average electrical output capacity differs from the capacity defined above by the availability factor. Replacement of the first wall and blanket is the major expected reason for downtime in a fusion reactor. Other factors will further reduce the availability but they are difficult to quantify and are ignored here. Wall and blanket life will be determined by the fluence limit which is estimated [36] to be within a factor of two of

$$I_w = 10 \text{ MW yr m}^{-2}.$$

(Thus, a wall will last 2.5 years with an incident neutron flux of 4 MW m^{-2} .) If it takes a time t_{down} (see Table III) to effect the wall-blanket replacement, then the availability will be

$$f_{\text{up}} = t_{\text{on}} / (t_{\text{on}} + t_{\text{down}}),$$

where

$$t_{\text{on}} = I_w A_w / 4 P_\alpha,$$

is the wall-blanket lifetime and A_w is the first wall area. The average net electrical capacity is finally

$$P_{\text{av}} = f_{\text{up}} P_e. \quad (30)$$

IV.B.2. Capital costs

The costs discussed in this report are not intended as estimates of total installed costs; rather, they are the relatively quantifiable costs. Our estimates include the direct cost of equipment and buildings, design and contingency allowances, and spares. We do not include the cost of administration buildings, overhead, or indirect costs in the belief that these would be proportional to the included costs, and therefore would not affect the relative ordering of our results. All costs have been adjusted to 1981 dollars in order to permit a more meaningful comparison of cost estimates made in different years. The dollar sign is omitted from all costs to emphasize our intention that the reactor costs estimated here be used only for internal comparisons and not for comparisons with real cost estimates of electrical generating plants.

The reactor costs have been sorted into three groups which reflect their different parameter scalings. The cost of the tokamak, its power supplies and building, should scale roughly as the engineered volume [37] V_E , of the tokamak. We define the engineered volume to be the volume enclosed by the toroidal field coils minus the plasma volume. The costs of the neutral beam system and its power supplies should scale with the delivered beam power. The costs of the electrical generating equipment (turbines, boilers, pumps, etc.) and its buildings should scale with the total thermal power of the tokamak. The result is then

$$C = C_E V_E + C_b P_b + C_{th} P_{th} , \quad (31)$$

where the coefficients are determined in the next section.

The final economic measure of a reactor is the cost per unit of net

electrical capacity:

$$C_W = C/P_{av} . \quad (32)$$

This, and not C, is referred to below, for brevity, as the cost of a reactor system.

IV.B.3. Reactor characteristics

For the electrical and cost coefficients defined above, we use the average of values extracted from the NUWMAK [38] and STARFIRE [39] design reports. The final adopted values are shown in Table III. The two reports are generally in close agreement; the only exception is P_E , which corresponds to a relatively small power drain in both studies.

IV.B.4. Neutral beam characteristics

We consider only deuterium beams in the reactor optimization study; impurity beam characteristics would be similar to those of the FQ deuterium beam system described below. Deuterium has been chosen because the most efficient of the proposed high energy neutral beam systems is based on deuterium, and the current-drive efficiency of deuterium is comparable to or greater than that of impurities [16].

Since prototypes for the neutral beams discussed here have not yet been built, the costs and electrical efficiencies are necessarily of questionable reliability. Our feeling is that the characteristics of the system which ultimately emerges as the best will be in the range considered here even though they may differ from the specific values assigned in Table III.

Since positive ions cannot be neutralized efficiently at the desired particle energies, all these neutral beams employ negative-ion sources. Likewise, all the designs use photodetachment neutralizers instead of gas neutralizers because of the premium on high efficiency. The only major design differences are in the choice of accelerator; their voltage limitations and power efficiencies determine the system characteristics with which we are concerned.

The first system considered is a "conventional" one, using electrostatic acceleration (denoted by EA) [40]. It is generally thought that the accelerator voltage will be limited by high-voltage breakdown to $E_{\max} \sim 400$ kV. At these energies a simple gas neutralizer could be used with only a modest reduction in efficiency.

The system proposed by TRW [2] employs an electrostatic accelerator with electrostatic-quadrupole (ESQ) focusing to suppress electron losses, and thereby raise $E_{\max}$ to ~ 2 MeV.

The third system uses a radio-frequency quadrupole (RFQ) accelerator [3]. The RFQ is the only proven accelerator at these energies, but the efficiency is relatively low. The RFQ could also be coupled to a source of impurity negative ions and a photodetachment neutralizer to produce neutral beams of impurity atoms. The cost and efficiency of impurity beams would possibly be slightly more favorable than a D^0 RFQ beam [41].

V. REACTOR OPTIMIZATION

The reactor equations developed above are now incorporated into a reactor optimization procedure, which is used to determine the optimum beam energy and plasma temperatures for beam-driven steady-state tokamaks. The resulting cost is compared with the optimum cost for pulsed tokamak reactors. The

sensitivity of the results to variations in several of the model parameters is discussed, and we describe an attempt to produce optimum conditions in a numerical simulation using a 1-d transport code.

V.A. Optimization procedure

In essence, our optimization procedure finds the combination of plasma temperatures and beam energy which produce the lowest cost reactor; the constraints and qualifications to this statement are detailed below. We have chosen not to optimize the tokamak and plasma dimensions in order to keep the problem manageable. The effects of changing some parameters, such as R_{tang} and Z_{eff} , are too subtle to be reliably represented in our simple 0-d model so we have fixed these parameters at what we expect are nearly optimum values.

For each point in a two-dimensional grid of values of β_{th} and $T_{av} = (T_e + T_i)/2$, the cost is calculated for each combination of trial values for T_e (which runs from 0.5 to 1.5 times T_{av}) and E_b (which runs from 100 keV to E_{max}). The parameters of the lowest cost case for each (β_{th}, T_{av}) pair are saved. The steps of the computation of the cost are:

- (1) T_i is computed from the choices of T_{av} and T_e ;
- (2) n_e is computed from β , T_e , and T_i ;
- (3) we compute η_{drive} from Eq. (17), P_b from Eq. (21), P_α from Eq. (24), h_b and h_α , from Eq. (26), and Q_{ie} from Eq. (27);
- (4) Eq. (22) and (23) are solved for τ_{Eec} and τ_{Eic} ;
- (5) if either τ_{Eec} or τ_{Eic} are < 0 that parameter set is disallowed.
- (6) We compute P_{th} from Eq. (28), P_e from Eq. (29), P_{av} from Eq. (30), C from Eq. (31), and C_W from by Eq. (32).

V.B. Optimum pulsed-ignition reactor

The procedure described above has been applied to pulsed-ignition operation by raising the coefficient A_{bd} in Eq. (17) to a large value; the beam power which is then required is so small that it has little effect on either the power balance or the cost. The resulting optimized cost is shown as a function of β and T_{av} in Fig. 4. The optimum temperatures are $T_i = 8.6$ keV and $T_e = 8.3$ keV.

These results are easily understood. The optimum value of T_a is close to the temperature which maximizes $P_\alpha(T)/T^2$ (when $\beta \approx n$ T is fixed, $P_\alpha = V_p n_D n_T f_\alpha(T) \propto \beta^2 f_\alpha(T)/T^2$) which is

$$T_{opt} = T_\alpha \exp(-1/B_\alpha),$$

when f_α is given by Eq. (25). Raising T_i above T_{opt} and lowering T_e increases the fusion power and lowers the cost, but this improvement is severely restricted by electron-ion equilibration in the power balance equations. As the temperature difference grows, equilibration transfers energy from the ions to the electrons, and the equilibration power equals the alpha power to the ions when $T_i - T_e$ is only a few hundred eV.

V.C. Optimum steady-state reactor

By using the calibrated value of A_{bd} (Table II), the optimization procedure finds the combination of density, temperatures, and beam energy which produce the lowest-cost steady-state operation. The results for the ESQ beam are shown in Fig. 5 as a function of β and T_{av} . The results for all three beams and the case where the beam energy is limited to the "Alfvén energy" (see Eq. 19) are shown for $\beta = 6.7\%$ in Table III. (In order to

facilitate comparisons with the NUWMAK [38] and STARFIRE [39] design studies, we used their value of $\beta = 6.7\%$.)

The qualitative features of the beam-driven results (Fig. 5) can be easily related to the pulsed-ignition results (Fig. 4). Raising T_e and T_i above T_{opt} , and lowering n_e to maintain constant β , improves η_{drive} , and thus reduces both the cost of the beams and their power consumption; as T_{av} is raised above T_{opt} , the fusion power drops only slightly at first. Eventually, the effect of the reduction in fusion power overtakes the benefits of reduced beam power; this occurs (by definition) at the optimum temperature for steady-state operation. The more costly and inefficient the beam system is, the more important it is to raise η_{drive} , and the higher the optimum temperatures are.

The importance of optimizing the plasma for current drive can be seen by observing that the cost for steady-state operation at the optimum temperature for pulsed operation is much higher than the cost at the optimum temperature for steady-state operation.

V.D. Comparison of pulsed-ignition and steady-state reactors

Since beams add extra cost and consume electricity, the steady-state reactors can compete with pulsed-ignition reactors only if they have some advantage in operating conditions. One such advantage could be higher availability; the duty cycle is higher, of course, and the much lower frequency of cyclic stresses suggests greater reliability. In addition, the presence of the injected fast-ion population may permit stable operation at higher β in the steady state case [5]. Most importantly, there is presently some concern that cyclic-stress-induced fatigue may not permit pulsed INTOR operation at the design magnetic field strength [4]. This could seriously reduce the fusion power, which is proportional to $\beta^2 B_0^4$.

An important property of the model we have employed is that the optimum cost for the pulsed-ignition case is a function of βB_0^2 ; that is, changes in B_0 can be offset exactly by corresponding changes in β . The same is true of the steady-state reactors in the EA, ESQ, or RFQ cases. As a result, it is a simple matter to determine reactor costs even if the magnetic field which is ultimately achievable in practice is lower than the value used here.

We have chosen to compare pulsed-ignition and steady-state operation by calculating the reduction in B_0 assigned to the pulsed case which produces a cost equal to that of steady state case (see Fig. 6). If for example, steady-state reactors could operate with ESQ beams, then they would be competitive with pulsed-ignition reactors operating with a fifteen percent reduction in B_0 and the same β .

V.E. Sensitivity of results to variations in model parameters

As noted above, a number of the reactor model parameters must be extrapolated beyond the range of present experience and, hence, are only rough estimates. It is important, therefore, to determine the sensitivity of the results to variations in these parameters. This has been done by changing one or two parameters at a time, and comparing the new optimum temperatures and cost with those in Table III. We fixed β at 6.7%; the dependence of the results on β has already been shown in Figs. 5 and 6.

The parameter sensitivities which are of the most concern are those which are the least understood and most uncertain. Two examples are the electron-energy confinement time, τ_{INTOR}^* , and the ion-electron equilibration power, Q_{ie}^* . There is little understanding of electron energy transport in present tokamaks, and the correct scaling to larger devices is uncertain. The equilibration power depends on details of the temperature profiles, and these

cannot be reliably predicted.

In view of these uncertainties, it is comforting to find that both the cost, C_W , and the optimum temperatures are insensitive to changes in either τ_{INTOR}^* or Q_{ie}^* . When τ_{INTOR}^* is changed by factors of 2 and 0.5, there is no change in the results. This occurs because τ_{INTOR}^* sets only an upper limit to τ_{Ee} , and in strongly heated reactors τ_{Ee} is much smaller than this limit. The equilibration power, Q_{ie}^* , does play a direct role in determining the results by limiting the difference between the electron and ion temperatures. When Q_{ie}^* is changed by factors of 2 and 0.5, the optimum cost varies by ± 5 -10%, while T_i varies by $\pm 5\%$, and T_e varies by ± 10 -15%.

Changes in the beam-driven current, I_p , should have a strong influence on the results because when I_p vanishes the reactor model reduces to the pulsed ignition case, and the costs and optimum temperatures change dramatically (compare Figs. 4 and 5). The value assigned to I_p is uncertain because the maximum desirable current in INTOR could be up to fifty percent larger than the design value; in addition, the bootstrap current could conceivably be substantial and that would reduce the beam-driven current requirement. When I_p is changed by factors of 1.5 and 0.5, the cost of the ESQ case varies by $\pm 25\%$. The other cases have more recirculating power and their cost varies by +200% and -50%. The optimum temperatures vary by ± 15 -25%.

Since the plasma current and beam-driven current efficiency, η_{drive} , appear in the reactor model equations in only Eq. (21) and in the combination I_p/η_{drive} , changes in I_p are exactly equivalent to inverse changes in η_{drive} . The above variations in I_p correspond to changes in η_{drive} by factors of 0.67 and 2. Discrepancies between the experimentally inferred and theoretically predicted beam-driven current in DITE [9] are generally $\lesssim \pm 30\%$.

The last prominent category of uncertain model parameters is the neutral beam characteristics. In one sensitivity survey, we fixed E_b at 2 MeV and considered all pairs of values from the sets $C_b = (1.3, 2.6, 5.2) \text{ watt}^{-1}$, and $\eta_b = (0.8, 0.6, 0.4)$. Using the ESQ case as a reference point, we find that C_W varies by $\sim 10\%$ to $+80\%$, T_e varies by -5% to $+15\%$, and T_i varies by -5% to $+25\%$.

The reactor net electrical output per unit cost, C_W^{-1} , is shown in Fig. 7 as a function of deuterium neutral beam energy, E_b , for $\eta_b = 0.8, 0.6$, and 0.4 . Reactor output rises rapidly with E_b near $E_b \sim 400 \text{ keV}$, but rises only slowly for $E_b > 1 \text{ MeV}$ and has essentially saturated at $E_b = 2 \text{ MeV}$. The price paid for inefficiency in either η_b or η_{drive} (low E_b produces low η_{drive}) is quite apparent. It is clearly very desirable to develop neutral beams with $\eta_b \geq 0.6$, and $E_b \geq 1 \text{ MeV}$.

V.F. 1-d transport simulation of beam-driven reactor

We used the 1-d tokamak transport code BALDUR to simulate a beam-driven reactor with a near-optimum plasma. Optimum conditions of $n_e = 1.5 \times 10^{14}$, $T_e = 14 \text{ keV}$, $T_i = 17 \text{ keV}$, and $E_b = 2 \text{ MeV}$ were chosen as the target values for the BALDUR simulation. The ion temperature was successfully controlled by the ripple-enhanced ion thermal conductivity. We attempted to control the electron temperature by adding tungsten impurity to radiate away electron thermal energy. As impurities were added the radiation from the cooler outer regions grew to the point of causing local thermal collapse before the radiation from the plasma interior was sufficient to reduce the central electron temperature. The simulation results were not optimum (compare Table II with the ESQ case in Table III); the cost figure for this simulation case is 5.61 watt^{-1} or $\sim 20\%$ higher than the optimum for this $\beta = 6.7\%$ case.

VI. CONCLUSIONS

It is now possible to envision neutral-beam-driven, steady-state reactors with acceptable levels of recirculated power. This result relies on the development of efficient, high-energy (> 1 MeV) neutral deuterium beams which would have much higher current-drive efficiency than the 200 keV D^0 beams used in previous studies [10,11]. More specifically, we find that steady-state reactors could be economically competitive with pulsed-ignition reactors if cyclic-loading problems force the toroidal magnetic-field strength of pulsed reactors to be $\sim 15\%$ lower than that of steady-state reactors (see Fig. 6).

We have found that the optimum plasma in a steady-state reactor differs considerably from the optimum for pulsed-ignition operation (compare Figs. 4 and 5). The optimum temperatures for the beam-driven reactors are higher, typically $T_e \sim 12-15$ keV and $T_i \sim 17-21$ keV, because the beam-driven current efficiency is proportional to $T_e/n_e \propto T^2/\beta$. In addition, it is actually advantageous to "seed" the plasma with impurities to improve the current-drive efficiency (high-Z impurities, which minimize fuel dilution, are preferred).

The large beam power (~ 100 MW) in a steady-state reactor has a strong impact on the confinement required for reactors. On the one hand, a driven reactor would tolerate lower bulk plasma confinement than an ignited reactor. On the other hand, a driven reactor would require a higher power handling capability in the burn-control mechanism; in addition, independent control of both the electron and ion temperatures is desirable in order to maintain each temperature near its optimum. The needed burn control processes have yet to be demonstrated experimentally, but ISX-B results [42] suggest that the expected high- β MHD instabilities may provide effective control if they manifest themselves through increased transport rather than catastrophic failure of equilibrium.

The effectiveness of the neutral beam systems can be judged by their overall current drive efficiency; i.e., the product of their current drive power efficiency, η_{drive} , and their electrical power to beam power conversion efficiency, η_b . (For the beams considered, the beam cost is inversely correlated with the overall efficiency, and thus not an independent consideration.) As can be seen in Fig. 7, near-optimum reactor output is possible with beams of 1 MeV neutral deuterium if $\eta_b \geq 0.6$. If either the beam particle energy or η_b is reduced significantly below these values, then the reactor output is seriously reduced.

Several key issues requiring further research have been identified. The most basic need is a broader experimental foundation for estimates of beam-driven current efficiencies; this includes resolution of the doubts concerning the existence of the trapped-electron correction to the driven electron current. The existence of deleterious beam-driven phenomena (e.g., the Alfvén-wave microinstabilities) must also be determined. The neoclassical bootstrap current would be very beneficial; a discovery of the conditions necessary to produce the bootstrap current could be invaluable. Finally, the study of the effects of fast-ion pressure and fast-ion driven currents on MHD equilibria and stability [42,43] should be extended to higher values of β .

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TABLE I

Maxima of $J(x,y)$

$$x = \sqrt{E_o/E_c} ; \quad y = Z_2/3 \approx 0.8 Z_{eff}/A_b$$

y	x_{max}^2	$J(x_{max},y)$	(a) x_-^2	(a) x_+^2
0	3.6	0.228	1.7	9
0.4	4.6	0.220	2.1	11
0.6	4.9	0.192	2.2	12
0.8	5.3	0.184	2.3	13
1.0	5.7	0.178	2.5	14
1.2	6.0	0.172	2.8	15
1.6	6.6	0.162	3.0	17
2.0	7.2	0.154	3.2	18
2.4	7.8	0.147	3.4	20

(a) $J(x,y) \geq 0.9 J(x_{max},y)$ for $x_- \leq x \leq x_+$

TABLE II

O-D Physics Model Parameters

n_e^*	$1.5 \times 10^{14} \text{ cm}^{-3}$	A_α	$7.0 \times 10^{-16} \text{ cm}^2$
T_e^*	17.7 keV	B_α	0.512
T_i^*	16.3 keV	T_α	54 keV
Q_{ie}^*	17 MW	A_{bd}	0.11 A/watt
τ_{INTOR}^*	1.6 sec	B_{bd}	2.0

TABLE III
Reactor Parameters

BEAM SYSTEM			EA	ESQ	RFQ	AL
1. Machine	Units					
V_E	m^3	3600				
A_w	m^2	356				
C_V	cm^{-3}	0.28				
2. Plasma						
R_o	m	5.3				
a	m	1.2				
b/a		1.6				
n_e	cm^{-3}		1.6×10^{14}	1.6×10^{14}	1.4×10^{14}	1.4×10^{14}
β_e	%	6.7				
I_P	MA	6.4				
T_i	keV		22	18	23	24
T_e	keV		12	14	16	13
Z_{eff}		2.1				
$f_{D,T} = n_{D,T}/n_e$		0.44				
$f_{He} = n_{He}/n_e$		0.05				
$f_W = n_W/n_e$		0.00019				
P_b	MW		220	110	85	190
E_b	MeV		0.4	2.0	2.5	0.43
R_{tang}	m	5.0				
E_{max}	MeV		0.4	2.0	--	^a E_{AL}
η_b			0.6	0.8	0.33	0.8
C_b	watt ⁻¹		2.6	2.6	5.5	2.6
Q			7	12	14	7
3. Power Output						
$5P_\alpha$	MW		1450	1350	1150	1300
M_b		1.2				
P_{gross}	MW _e		670	590	500	590
η		0.35				
C_{th}	watt ⁻¹	0.11				
P_{EVE}	MW _e	72				
P_E	watt _e cm ⁻³	0.02				
P_{av}	MW _e		190	330	150	250
C_W	watt _e ⁻¹		9.6	4.6	10.6	6.7
4. Reactor Components						
I_w	MW yr m ⁻²	10				
t_{down}	yr	0.5				
B_o	T	5.5				

^aSee Eq.19.

FIGURE CAPTIONS

- Fig. 1 The effective ionization cross section, σ_{eff} , for neutral deuterium in a hydrogen plasma with $Z_{\text{eff}} = 2$ due to Argon impurities.
- Fig. 2. The efficiency function $J(x,y)$ for selected values of y .
- Fig. 3. The collisional current factor $F(Z_b, Z_{\text{eff}}, \epsilon)$ for selected values of Z_b and Z_{eff} . The vertical axes are labeled with the element symbol.
- Fig. 4. Contours of constant cost, C_W (in watt_e^{-1}), for an INTOR-size pulsed tokamak reactor as a function of β and T_{av} . For each choice of β and T_{av} , we varied T_e to minimize the cost.
- Fig. 5. Contours of constant cost, C_W (in watt_e^{-1}), for an INTOR-size beam-driven steady-state tokamak reactor as a function of β and T_{av} . For each choice of β and T_{av} , we varied T_e , and E_b to minimize the cost.
- Fig. 6. The lowest cost, C_W , and the ratio of toroidal fields (pulsed/steady-state) for which the cost of pulsed and steady-state reactors are equal are shown as a function of β (for $B_0 = 5.5\text{T}$). The lowest cost is found by varying T_e , T_i , and E_b while holding β constant.
- Fig. 7. The reactor output per unit cost, C_W^{-1} , as a function of deuterium beam energy, E_b , for $\eta_b = 0.8, 0.6$, and 0.4 .

#81P0219

Fig. 1

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Fig. 2

Fig. 3

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Fig. 4

Fig. 5

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Fig. 6

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Fig. 7