JUNE 1975

MATT-1145

REACTOR ASPECTS OF
COUNTERSTREAMING-ION
TOKAMAK PLASMAS

BY

D. L. JASSBY

PLASMA PHYSICS
LABORATORY

MASTER

PRINCETON UNIVERSITY
PRINCETON, NEW JERSEY

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REACTOR ASPECTS OF COUNTERSTREAMING-ION TOKAMAK PLASMAS

D. L. JASSBY

Plasma Physics Laboratory, Princeton University
Princeton, N. J. 08540 U.S.A.

ABSTRACT

Toroidal DT plasmas in which the D and T ions make up two distinct, quasi-thermal velocity distributions, oppositely displaced in velocity along the magnetic axis, are discussed. Such counterstreaming distributions can be set up by introducing all ions by tangential injection of neutral beams, and by removing ions from the plasma shortly after they have decelerated to an energy $\sim 2T_e$ by Coulomb drag on the plasma electrons. A simple physical model for counterstreaming-ion operation is postulated, which allows one to deduce the ion velocity distributions and required energy and particle confinement times that are in good agreement with the results of previous Fokker-Planck calculations. The variations of fusion reactivity, power gain, and power density with injection energy and electron temperature are presented. We consider the practical problems of implementing counterstreaming operation in a tokamak, such as charge-exchange losses, the prompt removal of cold ions, and the effect of impurities.

Several properties of the counterstreaming system that make it advantageous for near-term reactor applications are (1) optimal performance at relatively low injection energy

(40-60 keV), because of head-on nuclear collisions,
(2) small electron energy confinement time for $Q \sim 1$
(e.g., $n_e \tau_E \approx 2 \times 10^{12} \text{ cm}^{-3}$ at $T_e \approx 3.5 \text{ keV}$), (3) re-
fueling solely by the injected beams, and (4) maintenance
of the plasma current by the beams. For realistic tokamak
parameters, neutron wall loading is limited to 0.5 MW/m^2
or less. Plasma parameters appropriate to radiation test
reactors and fusion-fission hybrid reactors are discussed.

1. INTRODUCTION

In thermal DT plasmas at moderate ion temperatures ($T_i \lesssim 15$ keV), head-on collisions between D and T ions are the most effective in producing fusion energy. In this paper we consider a toroidal plasma in which the D and T ions are grouped in two distinct, quasi-thermal velocity distributions, oppositely displaced in velocity along the magnetic axis, as shown in Fig. 1. The relative displacement velocity v_r is near the optimum value for generating fusion power, which occurs principally by nearly head-on collisions between D and T ions of the counterstreaming distributions. Such velocity distributions can be set up in a toroidal plasma container by introducing all ions by means of neutral beams injected parallel and anti-parallel to the toroidal magnetic field. The energetic ions slow down by Coulomb drag on the plasma electrons, and if decelerated ions are removed from the plasma in a time that is short compared with the fast-ion slowing-down time, virtually all plasma ions will be found in the two oppositely directed "beams". The electron temperature T_e is maintained entirely by power flow from the fast ions. This system of space-charge neutralized, counterstreaming-ion distributions has been called a colliding-beam toroidal plasma, or CBT [1]. The Fokker-Planck analysis of Ref. [1] has shown that collisions among ions of the same beam cause significant diffusion of ion energy during their slowing-down history following injection, so that the ions form displaced Maxwellian-like distributions, as in Fig. 1. The beam "temperature" is sufficiently large so that the counterstreaming system is apparently stable to both electrostatic and electromagnetic velocity-space modes.

The CBT system possesses certain advantages over thermonuclear and target-plasma (TCT) reactors. For the same fusion reactivity, injection voltages can be 3 to 4 times smaller than in the TCT, because head-on collisions increase the relative collision velocity. Somewhat larger Q -values can be attained, because of the reduction in cold-ion drag. For $T_e = 2-6$ keV, the electron confinement parameter $n_e \tau_E$ is reduced by a factor of 2-4 in the CBT, to the range $0.5-5 \times 10^{12} \text{ cm}^{-3} \text{ s}$. Other properties advantageous for reactor applications are that the plasma is fueled entirely by the beams, so that pellet injection technology is not required, and the plasma current can be maintained indefinitely by means of the injected beams.

Because $Q \sim 1$ operation is attainable at modest values of T_e , $n_e \tau_E$, and beam voltage, it seems appropriate to explore the practical use of CBT plasmas. While the ideal fusion gain in the CBT is limited to perhaps 5-10, there are several potential applications of fusion reactors for which $Q \sim 1$ is quite adequate. These applications include fission-related activities such as fissile breeding and actinide burning [2], the large-scale production of tritium, and large-area neutron bombardment testing of wall materials and blankets. For such applications, fusion power density P_f is of primary importance [3]. It is with an eye to these applications that this paper has evolved. (Another study will discuss means of maximizing Q in a counterstreaming reactor, in order to enhance its viability as a stand-alone electrical power producer.)

In Sec. 2 we develop a simple physical model that allows us to conveniently derive the ion velocity distributions and the energy and particle confinement times in CBT plasmas. These distributions are used to calculate Q and P_f over a range of injection voltages and T_e . While beam-driven, counterstreaming operation can in principle be applied to most toroidal configurations, in this study we shall direct our attention particularly to tokamak plasmas, since tokamak operation has proved to be the most successful to date. Section 3 discusses practical problems in the implementation of a counterstreaming-ion tokamak plasma, including the removal of decelerated ions, and the effects of charge exchange and impurities. The attainable neutron wall loading and illustrative reactor parameters for several applications are discussed in Sec. 4.

2. POWER BALANCE AND FUSION REACTIVITY

2.1 ION VELOCITY DISTRIBUTIONS

In this paper we consider only steady-state operation, although the CBT could also be operated in a pulsed mode [1]; the latter mode seems somewhat unsuitable for reactor applications, however. Consider a steady-state counterstreaming-ion plasma into which energetic neutrals are injected parallel and anti-parallel to the magnetic axis. The neutrals are trapped by charge-exchange with ions travelling in the same direction, by impact ionization on oppositely travelling ions, and by electron ionization. The relative importance of the various trapping processes is discussed in Sec. 3.1.

The energetic ions are assumed to be confined until they decelerate to an energy $W \sim 2T_e$ by Coulomb drag on the electrons. Ions with energy $W < 2T_e$ are assumed to be lost at a rate $\tau^{-1} \sim \tau_E^{-1}$, where τ_E is the electron energy confinement time. It turns out that for $T_e \gtrsim 10$ keV, the slowing-down time of the fast ions is many times longer than τ_e . Methods of removing cold ions from the plasma are discussed in Sec. 3.2. By means of a magnetic divertor, the inward flow rate of cold neutrals and ions from the wall can be kept small compared to the large rate of particle injection by the beams.

Let us consider the nature of the ion velocity distributions. The equilibration time among ions in the same beam is many times shorter than the slowing-down time on electrons. Interaction between ions in opposite beams is particularly important in causing angular scattering, and therefore a degree of isotropization. Thus one expects the ions in each beam to form a "thermal" distribution, with a mean velocity in the laboratory frame somewhat less than the injection velocity. Ref. [1] reports calculations of the ion distribution functions $f(v)$ with a multi-species, two-dimensional Fokker-Planck code that numerically solves the Fokker-Planck equation over a uniform mesh in velocity space. More extensive results are given in [4]. These Fokker-Planck results demonstrate that $f(v)$ does indeed resemble a displaced Maxwellian, with $T_{\perp}$ slightly larger than $T_{\parallel}$. The determination of most plasma parameters, as well as the reactor performance, is dependent on the accurate evaluation of $f(v)$. The numerical Fokker-Planck solutions are very expensive, however, and are not practical for extensive investigation of a wide range of parameters such as T_e , W_0 , n_D/n_T . In

the following, we describe a simple physical model that provides distribution functions that agree remarkably well with the known Fokker-Planck results for $Z = 1$.

Consider a test ion, injected at energy W_0 , that slows down by Coulomb drag on plasma electrons, but has no interaction with other ions. The equation of motion of this ion is

$$\frac{dW}{dt} = - \frac{C n_e \ln \Lambda_e}{A_i T_e^{3/2}} (W - \frac{3}{2} T_e) \quad (1)$$

Here $C = 1.06 \times 10^{-13}$ (T_e, W in keV), and A_i is the atomic mass of the ion. The average energy of this ion during its slowing-down time to an energy $2T_e$ is

$$W_s = \frac{W_0 - 2T_e + \frac{3}{2} T_e \log \left(\frac{W_0 - \frac{3}{2} T_e}{\frac{1}{2} T_e} \right)}{\log \left(\frac{W_0 - \frac{3}{2} T_e}{\frac{1}{2} T_e} \right)} \quad (2)$$

In the case of continuous injection of noninteracting ions, each of which is removed immediately after deceleration to $W = 2T_e$, the average energy of the distribution is also W_s . In practice, when $n_b \sim n_e$, where n_b is the fast-ion density, the ion-ion collision time is very short compared to the slowing-down time t_s on electrons, unless T_e is very small. If τ_{12} is the energy exchange time between two ions with energies W_0 and $2T_e$, respectively, we have

$$\frac{\tau_{12}}{t_s} \approx \frac{1}{60} \left(\frac{W_0}{T_e} \right)^{3/2} \frac{1}{A_i^{1/2}} \frac{1}{\ln \left(\frac{W_0}{2T_e} \right)} \sim \frac{1}{10} + \frac{1}{3} \quad (3)$$

Ions with a smaller energy differential will interact even more rapidly, and some ions will be accelerated to energies well above the injection energy. One may reasonably propose that the ions

form a displaced Maxwellian of temperature T_b , centered at $v_p = (2W_s/M)^{1/2}$, as illustrated in Fig. 2. The normalized ion distribution function is

$$f(\vec{v}) = \left(\frac{M}{2\pi T_b} \right)^{3/2} \exp \left[- \frac{M}{2T_b} (\vec{v} - v_p \hat{\phi})^2 \right] \quad (4)$$

where $\hat{\phi}$ is the unit vector along the magnetic axis. The average energy of this distribution is

$$\bar{W} = \frac{1}{2} M v_p^2 + \frac{3}{2} T_b \quad (5)$$

The determination of T_b requires a further condition. Because of rapid ion-ion thermalization, the entire kinetic energy of the beam tends to be randomized. However, the directed motion is maintained by continuous ion injection at $W_o > W_s$ and by removal at $W \lesssim 2T_e < W_s$. The injected energy is not entirely available for randomization among the ions, but also must be used to maintain the electron random motion (i.e., T_e). We postulate an equipartition of energy between the directed ion motion and the plasma thermal motion. That is,

$$\frac{1}{2} M v_p^2 = \frac{3}{2} T_e + \frac{3}{2} T_b \quad (6)$$

The equipartition postulate is justified by demonstrating that Eq. (4), with W_s and T_b determined from Eqs. (2) and (6) respectively, predicts plasma performance in agreement with the results of the previous Fokker-Planck calculations. (This postulate is also expected to hold when $Z_{eff} > 1$, but with $T_{b\perp} > T_{b\parallel}$.)

Note that this model does not require interaction between ions of opposite beams. (In practice, angular scattering by ions of the opposing beam is important for the isotropization about v_p .) In most of this work, the deuterons and tritons are injected with the same energies; since W_s is independent of mass, the D and T distributions have the same temperature.

Figure 3 shows the dependence of v_p , T_b , and $\bar{W}$ on T_e , using $W_0 = 60$ keV and Eqs. (2), (5), and (6). These results, particularly the temperature dependence of $\bar{W}$, are in good agreement with the Fokker-Planck results [1,4]. In the latter calculations, "thermal" ions, that is, ions decelerated below $2T_e$, have a lifetime τ_E . While the proportion f_T of "thermal" ions may be as large as 30% for $T_e \gtrsim 12$ keV, the vast majority of these ions retain their direction of motion along the toroidal magnetic field.

For practical values of W_0 and T_e , the system of velocity distributions in the CBT is apparently stable to electrostatic and electromagnetic velocity-space modes [1,5]. If v_0 is the injection velocity, then $(v_0/v_{\text{Alfvén}})^2 \approx \beta_t$, independent of v_0 . Since $\beta_t \ll 1$ in tokamak operation, Alfvén instabilities seem unlikely to occur. The stability of finite-medium modes has not yet been investigated, but presumably a counterstreaming-ion plasma would be as susceptible as most toroidal plasmas to drift instabilities.

2.2 POWER BALANCE AND ENERGY CONFINEMENT

The electron temperature is maintained entirely by power flow from the fast ions, so that in practice T_e and n_e are determined by τ_E and the ion source strength S . T_e is found from the

steady-state power balance relation

$$\frac{3}{2} \frac{n_e T_e}{\tau_E} = \frac{C n_e \ln \Lambda_e}{T_e^{3/2}} \sum_{i,2} \frac{n_i (\bar{W}_i - 3/2 T_e)}{A_i} \quad (7)$$

Here n_i is the density and $\bar{W}_i$ the average ion energy in each beam, as given by Eq. (5). In this work, we consider only the case $Z=1$, so that $n_e = n_1 + n_2$. The ion source strength is

$$S_i = \frac{n_i}{\tau_{hi}} \quad (8)$$

where τ_{hi} is the ion lifetime.

In the following we consider only cases where $n_1 = n_2$ and $W_{o1} = W_{o2}$. (It should be noted, however, that for controlling plasma momentum and current, these conditions may have to be altered.) Defining $\tau_h^{-1} = (\tau_{h1}^{-1} + \tau_{h2}^{-1})/2$, τ_h/τ_E is given by

$$\frac{3}{2} \frac{n_e T_e}{\tau_E} = \frac{n_e (W_o - \alpha T_e)}{\tau_h} \quad (9)$$

where $\alpha T_e \sim 3/2 T_e$ is the average energy of the ions leaving the plasma. Both τ_E and τ_h are inversely proportional to n_e , except for the weakly varying $\ln \Lambda_e$ factor. Figure 4 shows τ_h/τ_E as a function of T_e and W_o for $\alpha = 2$. For most conditions, the ion lifetime must be many times τ_E . This requirement seems compatible with observations on beam-injected tokamak plasmas [6], where in the absence of charge-exchange losses, the energetic ions remain in the plasma for their slowing-down time ~ 10 ms, while τ_E is in the range 2-4 ms. In Ref. [7] it was shown that microinstabilities, such as drift instabilities, are relatively ineffective in causing

diffusion of superthermal ions, so that one would expect fast ions to be lost prematurely only by faulty magnetic surfaces, or by charge-exchange.

In analyzing CBT reactor performance, it is often most convenient to choose n_e and T_e , and to calculate the corresponding value of $n_e \tau_E$ from Eq. (7). Figure 5 shows $n_e \tau_E$ as a function of T_e and W_0 . Evidently, the predicted $n_e \tau_E$ for $W_0 = 60$ keV is in excellent agreement with the values obtained from the numerical Fokker-Planck solutions, thus lending support to the model of Sect. 2.1. Also shown in Fig. 5 is $n_e \tau_E$ for a TCT plasma with $W_0 = 200$ keV, the same relative collision energy as in the 60-keV CBT case. In the TCT, $n_e \tau_E$ may have a wide range of values for a given T_e , depending on the choice of n_b/n_e ; the curve in Fig. 5 is calculated for the case of maximum P_f [3]. In the CBT, on the other hand, $n_b/n_e \approx 1$, so that there is a unique value of $n_e \tau_E$ for each T_e (at constant W_0). At low T_e , this value can be many times smaller than in the TCT, where many more particles must be heated by a given population of energetic ions. But at large T_e , the fast-ion slowing-down time is much longer in the CBT, so that $n_e \tau_E$ values tend to be comparable.

A useful quantity in discussing beam-maintained plasmas is $\bar{\Gamma}$, defined as

$$\bar{\Gamma} = \frac{\text{superthermal-ion energy density}}{\text{thermal-plasma energy density}} = \frac{n_b (1-f_T) \bar{W}}{\frac{3}{2} n_e (1+f_T) T_e} \quad (10)$$

In typical operation of a target-plasma reactor, $\bar{\Gamma} \lesssim 1$ [3]; in the CBT, $\bar{\Gamma} \approx \frac{2}{3} (\bar{W}/T_e) \gg 1$, and $\bar{\Gamma}$ is essentially the ratio of the ion pressure to the electron pressure.

2.3 FUSION REACTIVITY

In this section we calculate the fusion reactivity of the system of counterstreaming D and T ions. The reactivity is

$$\langle \sigma v \rangle_c = \frac{\iint d\vec{v}_1 d\vec{v}_2 \sigma(v_{12}) v_{12} f_1(\vec{v}_1) f_2(\vec{v}_2)}{\iint d\vec{v}_1 d\vec{v}_2 f_1(\vec{v}_1) f_2(\vec{v}_2)} \quad (11)$$

where $v_{12} = |\vec{v}_1 - \vec{v}_2|$, f_1 and f_2 are given by Eq. (4), and $\sigma(v_{12})$ is the fusion cross-section for D and T nuclei with relative collision velocity v_{12} . We consider only the case of equal injection energies for D and T. Then as discussed in Sect. 2.1, each distribution has the same T_b , while $v_{pT} = -\sqrt{\frac{2}{3}} v_{pD}$. The six-fold integrals in Eq. (11) can be reduced to single integrals by transforming the variables of integration to $\vec{v}_{12} = \vec{v}_1 - \vec{v}_2$ and $\vec{V} = (M_1 \vec{v}_1 + M_2 \vec{v}_2)/(M_1 + M_2)$, where M_1 and M_2 are the masses of the two ion species. The result is

$$\langle \sigma v \rangle_c = \frac{\int_0^{v_{\max}} dv_{12} \sigma(v_{12}) v_{12}^2 \left(e^{C_1 v_{12}} - e^{-C_1 v_{12}} \right) e^{-\mu v_{12}^2 / 2T_b}}{\int_0^{v_{\max}} dv_{12} v_{12} \left(e^{C_1 v_{12}} - e^{-C_1 v_{12}} \right) e^{-\mu v_{12}^2 / 2T_b}} \quad (12)$$

Here $\mu = M_1 M_2 / (M_1 + M_2)$, $C_1 = (1 + \sqrt{\frac{2}{3}}) \mu v_{pD} / T_b$, $v_{\max} \gg |v_{pD}| + |v_{pT}|$.

The integrals in Eq. (12) have been evaluated numerically, using $\sigma(v)$ from Ref. [8]. Figure 6 shows $\langle \sigma v \rangle_c$ as a function of $v_r = |v_{pD}| + |v_{pT}|$, the relative displacement of the peaks of $f_1(\vec{v}_1)$ and $f_2(\vec{v}_2)$ [Fig. 1]. Several points are worth noting:

- (1) For $v_r = 0$, $\langle \sigma v \rangle_c$ is the usual fusion reactivity of an isotropic Maxwellian DT plasma.

(2) As $T_b \rightarrow 0$, the peak value of $\langle \sigma v \rangle_c$ approaches the maximum value for a beam-target system. This value is $1.67 \times 10^{-15} \text{ cm}^3/\text{s}$ for $W = 128 \text{ keV}$.

(3) For any T_b , $\langle \sigma v \rangle_c$ can be larger than the maximum reactivity of a thermonuclear plasma, viz. $8.8 \times 10^{-16} \text{ cm}^3/\text{s}$ at $T_i = 60 \text{ keV}$.

Thus in the counterstreaming mode, a low temperature distribution of ions may have a very large fusion reactivity. However, the distribution must be maintained by continuous beam injection, so that the system has only a moderate Q , as shown in the next section.

2.4 FUSION POWER GAIN

In steady-state operation, the fusion power gain is $Q = P_f/P_b$, where P_f is the fusion power density and P_b is the injected power density:

$$P_f = n_D n_T \langle \sigma v \rangle_c E_f = \frac{1}{4} n_e^2 \langle \sigma v \rangle_c E_f, \quad (13)$$

where $E_f = 17,600 \text{ keV}$, and $n_D = n_T = \frac{1}{2} n_e$;

$$P_b = \frac{W_{OD} n_D}{\tau_{hD}} + \frac{W_{OT} n_T}{\tau_{hT}} = \frac{W_o n_e}{\tau_h}, \quad (14)$$

where τ_h is defined as in Eq. (9).

$$Q = \frac{1}{6} \langle \sigma v \rangle_c \frac{E_f}{W_o} \frac{n_e \tau_E (W_o - \alpha T_e)}{T_e} \quad (15)$$

$n_e \tau_E$ and $\langle \sigma v \rangle_c$ are found from Eqs. (7) and (12), respectively. Figure 7 compares the predicted temperature dependence of Q with previous Fokker-Planck results for $W_{OD} = 60$ keV and $W_{OT} = 40$ keV [1]. The agreement is everywhere better than 10%, again lending support to the model of Sect. 2.1. Note also the large increase in Q over the TCT reactor [9], in which the cold target ions exert appreciable drag on the relatively few beam ions.

Figure 8 shows the dependence of Q on injection energy and electron temperature. Evidently, Q is rather insensitive to W_0 ; for each T_e there is a broad maximum in the range $W_0 = 40$ -60 keV, corresponding to a relative collision energy of 130-200 keV at injection. For $W_0 \gg 3/2 T_e$, and ignoring the variation in W_s with T_e , Eqs. (7) and (15) predict that $Q \propto (n_e \tau_E)^{3/5}$ or $Q \propto T_e^{3/2}$. (The Q -curves are not plotted for $W_0 \leq 4 T_e$, since for these conditions, the calculation of W_s is somewhat sensitive to the "cutoff" $W = 2 T_e$, which is used in the derivation of Eq. (2).)

In this work we assume that alpha particles from fusion reactions are not confined. Retention of fusion alphas does not significantly increase Q , and finite-alpha pressure reduces the allowable beam-ion density, thereby lowering P_f .

2.5 FUSION POWER DENSITY

It is instructive to compare semi-quantitatively the permissible values of n_e and P_f in the counterstreaming and target-plasma reactors, for the same values of T_e and total plasma pressure p . In the TCT, with $T_e = T_i$, we have

$$p = 2(n_e)_{\text{TCT}} T_e (1 + \bar{\Gamma}) \approx 4(\bar{n}_e)_{\text{TCT}} T_e \quad (16)$$

where $\bar{\Gamma} \approx 1$ is near the optimal value for maximizing P_f [3]. In the CBT,

$$p = (n_e)_{\text{CBT}} (T_e + \frac{2}{3} \bar{W}) \approx \frac{2}{3} \bar{W} (n_e)_{\text{CBT}} \quad (17)$$

Here we have noted that superthermal ions produce an isotropic pressure, even if these ions have a streaming motion along B_t [10]. Combining Eqs. (16) and (17), we have

$$\frac{(n_e)_{\text{CBT}}}{(n_e)_{\text{TCT}}} \approx \frac{6T_e}{\bar{W}} \approx \frac{1}{2} \rightarrow 1 \quad (18)$$

(See, for example, Fig. 3c.) In the TCT, typically $n_D \approx \frac{1}{8} n_e$, while $\langle \sigma v \rangle_c \approx \langle \sigma v \rangle_{\text{TCT}}$. Therefore

$$\frac{(P_f)_{\text{CBT}}}{(P_f)_{\text{TCT}}} = \frac{1/4 (n_e)_{\text{CBT}}^2 \langle \sigma v \rangle_c}{(n_D n_T)_{\text{TCT}} \langle \sigma v \rangle_{\text{TCT}}} \approx \frac{1}{2} \rightarrow 2 \quad (19)$$

Thus both n_e and P_f are comparable in the two modes of operation.

The pressure limitation of a tokamak plasma is given by the MHD equilibrium condition on the poloidal beta:

$$\beta_p = \frac{8\pi \bar{p}}{B_p^2} = \frac{8\pi \bar{p} q^2 A^2}{B_t^2} \leq A \quad (20)$$

where $A = R/a$ is the aspect ratio, and q is the safety factor at the limiter. In this section we use the illustrative values

$B_t = 50$ kG, $q = 2.5$, $A = 5.0$, $\beta_p = 3.0$, so that $\bar{p} = 0.19$ J/cm³.

Then P_f is found from Eqs. (12) and (13), using $n_D = n_T = \frac{1}{2} n_e = \frac{1}{2} p / (T_e + \frac{2}{3} \bar{W})$. Figure 9 shows the dependence of P_f on W_0 and T_e . P_f is maximum at the smallest T_e and W_0 , since for constant pressure, larger n_e is permitted at smaller electron and beam energies. The minimum practical W_0 is determined by the difficulty of the neutral beams in penetrating the plasma (Sec. 3.1), while the minimum desirable T_e is determined by the minimum Q-value required (cf. Fig. 8).

Figure 10 shows Q vs P_f for constant W_0 and p. Each point on a constant- W_0 curve corresponds to a particular value of $n_e \tau_E$, and therefore a definite T_e given by Eq. (7). There is necessarily a trade-off between P_f and Q, the largest values of P_f being obtained at the smallest $n_e \tau_E$ and T_e . The large- $n_e \tau_E$ region is the most appropriate for stand-alone power producers, while the low- $n_e \tau_E$ region is the most attractive for maximizing neutron production. In the latter operating regime, the throughput of plasma energy is extremely large, and the most important confinement role of the tokamak is the confinement of the injected ions during their long slowing-down period (cf. Fig. 4).

3. IMPLEMENTATION OF COUNTERSTREAMING OPERATION IN A TOKAMAK

3.1 BEAM TRAPPING

A target plasma is required to initiate CBT operation. This target plasma can be a low-temperature tokamak plasma ($T_e \sim 1$ keV) of moderate density ($n_e \sim 2 \times 10^{13} \text{ cm}^{-3}$), set up in the usual way by gas filling and ohmic heating. Beam-injected neutrals are

trapped by charge-exchange and ionization by plasma ions and electrons. Once beam injection has begun, the divertor coils are energized, so that recycling of outwardly diffusing ions is minimized; essentially all further particle input to the plasma comes from the beams.

After injection for several slowing-down times, T_e reaches its final value, and the ion population is dominated by energetic ions. In the steady state, injected neutrals are trapped by the ions in the counterstreaming distributions, and by electron ionization. The rate coefficients shown in Fig. 11, taken from Ref. [11], are required for calculating the trapping length. Consider a deuterium atom (D^0) injected tangentially with velocity v_0 into the CBT plasma. The relative collision energy U between the D^0 and D^+ ions is very small (e.g., $U \approx 7$ keV for the $T_e = 5$ keV case of Fig. 3), while U for collisions between the D^0 and T^+ ions is very large ($U \approx 140$ keV for the same case). Approximately half the neutrals are trapped by charge-exchange with ions of the same species, and the other half by impact ionization on ions of the other species, and by electron ionization. The penetration length is given by

$$\lambda_t = v_0 \left[n_1 \langle \sigma_{cx}(u_{11}) u_{11} \rangle + n_2 \langle \sigma_{cx}(u_{12}) u_{12} \rangle + n_1 \langle \sigma_I(u_{11}) u_{11} \rangle + n_2 \langle \sigma_I(u_{12}) u_{12} \rangle + n_e \langle \sigma_e v_e \rangle \right]^{-1} \quad (21)$$

where $u_{11} = |\vec{v}_0 - \vec{v}_1|$ and $u_{12} = |\vec{v}_0 - \vec{v}_2|$ are the relative velocities between the fast neutral and ions of the two species of densities

n_1 and n_2 , σ_{cx} and σ_I are the charge-exchange and impact-ionization cross-sections, respectively, and $\langle \sigma_e v_e \rangle$ is the electron ionization rate, averaged over a Maxwellian electron velocity distribution. For each charge-exchange or impact-ionization process,

$$\langle \sigma v \rangle = \frac{\int d\vec{v}_i \sigma(u) u f_i(\vec{v}_i)}{\int d\vec{v}_i f_i(\vec{v}_i)} \quad (22)$$

where $f_i(\vec{v}_i)$ is given by Eq. (4). The three-fold integrals are reduced to a single integral by the transformation $\vec{u} = \vec{v}_0 - \vec{v}_i$.

The result is

$$\langle \sigma v \rangle = \frac{\int_0^{u_{\max}} du \sigma(u) u^2 \left(e^{C_3 u} - e^{-C_3 u} \right) e^{-M_i u^2 / 2T_b}}{\int_0^{u_{\max}} du u \left(e^{C_3 u} - e^{-C_3 u} \right) e^{-M_i u^2 / 2T_b}} \quad (23)$$

where $C_3 = M_i(v_{pi} - v_0)/T_b$, $v_{pi} = (2W_s/M_i)^{1/2}$, and $u_{\max} \gg |v_0| + |v_{pi}|$.

Figure 12(a) shows λ_t for D^0 as a function of neutral energy, with $n_e = 5 \times 10^{13} \text{ cm}^{-3}$ and $T_e = 4.0 \text{ keV}$. The dependence of λ_t on T_e enters through $\langle \sigma_e v_e \rangle$, and through the variation of $f_i(\vec{v}_i)$ with T_e . In this work, $W_{OT} = W_{OD}$, so that λ_t for T^0 is only 81% of that for D^0 . This difference is generally not significant, however, since in steady-state operation, the ions will be distributed over most of the radial profile, and in any event, one can inject D and T beams along chords that are displaced by different amounts from the magnetic axis. (Actually, the difference in λ_t -values could be compensated in a tokamak by injecting T^0 in the

direction of the current, J , and injecting D^0 in the opposite direction. Then the drift surfaces of the resulting D^+ shift outward, while those of the T^+ shift inward [12].)

Figure 12(a) also shows R_{cx}^0 , the probability of the injected neutral being trapped by charge-exchange. The average energy of the secondary neutral is in the range $0.6 - 0.9 \bar{W}$, depending on W_{OD} and T_e . This neutral will be travelling nearly parallel to the magnetic axis, and will undergo charge-exchange or ionization before escaping. Successive generations of neutrals are formed in this way, until finally ionization occurs, or the neutral escapes. The trapping process thus leads immediately to a certain loss of energy.

The charge-exchange energy loss can be estimated with the aid of Fig. 13. For illustrative purposes, we consider a plasma of aspect ratio $A = 5.0$, uniform n_e , radius $a = \lambda_t$, $T_e = 4$ keV, and $W_{OD} = 60$ keV. The injection trajectory is a chord lying slightly inside the magnetic axis, so that penetration to the axis is adequate even when $\lambda_t \lesssim a$ [12]. The total path length in the plasma is $L \approx 2(2R_0 a + a^2)^{1/2} = 6.65a$. For $W_{OD} = 60$ keV, the average energy of the charge-exchange neutral is $W_n = 34$ keV. As shown in Fig. 12(b), this neutral in turn has a mean-free-path for a charge-exchange or ionization collision of $\lambda_n = 0.75 \lambda_t$, with $R'_{cx} = 0.64$. Suppose that all succeeding neutrals have energy W_n , and that the N th generation neutral is lost from the plasma. Then the fraction of injected energy lost by charge-exchange is

$$F_{cx} = \frac{W_n}{W_0} R_{cx}^0 (R'_{cx})^{N-1} \quad (24)$$

Since $L = 6.65a \approx \lambda_t + 7\lambda_n$, N could be as large as 8. However, the injected neutrals and plasma ions have velocities that are not exactly parallel to the magnetic axis, so that $N = 4$ or 5 seems reasonable, for which F_{cx} is 0.09 and 0.06 respectively. Both Q and P_f are reduced by approximately the factor $1 - F_{cx}$.

According to our basic postulates, neutrals enter the plasma only by beam injection (recycling is not permitted), so that the above analysis accounts completely for charge-exchange loss. This process reduces Q and P_f by less than 10%. While the charge-exchange loss is smaller at larger W_0 , both Q and P_f are inherently reduced for $W_0 > 60$ keV (cf. Figs. 8 and 9).

3.2 ELIMINATION OF DECELERATED IONS

In order to maintain ion velocity distributions such as shown in Fig. 2, two conditions must be met:

(1) The rate of recycling of cold ions and neutrals must be small compared with the rate of particle input by the neutral beams. This condition can probably be met with the use of a poloidal magnetic divertor [13]. Such a divertor is required also for minimizing bombardment of the vacuum wall by outgoing particles. In any event, the rate of ion energy outflow ($\sim 3/2 n_e T_e / \tau_h$) is less than the plasma power throughput by the ratio $\tau_E / \tau_h \ll 1$. Practically all the injected energy is lost through the electron channel.

(2) Decelerated ions (ions with $W \gtrsim 2T_e$) must be removed from the plasma in a time τ that is short compared with their slowing-down time; i.e., $\tau \ll \tau_h$. The results of Ref. [4] have shown that Q is rather insensitive to the exact value of τ / τ_h , when $\tau \sim \tau_E$.

There are three processes that appear effective for the rapid removal of decelerated ions:

(i) Charge-exchange of injected neutrals. In the example of Sec. 3.1, the ratio of the particle loss rate by charge-exchange to the particle injection rate is $F_{cx} W_o/W_n \sim 1/8$. While the average energy of the escaping neutrals is relatively large, an appreciable fraction of these neutrals have low energy.

(ii) Instability-induced diffusion. The temperature of a practical CBT plasma will be sufficiently high so that electrons and ions will lie in the collisionless "banana" regime, and trapped-particle instabilities can be expected to exist [14]. While these instabilities apparently cause negligible radial diffusion of energetic ions, because of the short wave-particle correlation times [7], they should be rather effective in the elimination of thermal ions and electrons. Collective drift-type micro-instabilities cause comparable loss rates of particles and energy, so that if they are chiefly responsible for the electron energy loss rate τ_E^{-1} , one expects that $\tau^{-1} \sim \tau_E^{-1}$.

It is desirable to operate with a large "peaking factor" of the injected beams [12], so that there will be a gradual outward diffusion of the fast ions during their slowing-down period. The colder ions, which are susceptible to instability-induced diffusion, will have a relatively short diffusion length before they are captured by the poloidal divertor.

(iii) Toroidal magnetic field ripple. Ripple in the toroidal magnetic field can lead to enhanced diffusion of thermal plasma particles [15]. But ripple has the undesirable effect of enhancing

deceleration of the counterstreaming velocities in the laboratory frame [16], and is also effective in eliminating energetic ions trapped in banana orbits. Field ripple is generally small in a tokamak (e.g., $\delta B/B \ll 1\%$ near the magnetic axis), but it can be made rather large in the outer region of the plasma. Ripple-induced loss would be most beneficial in the case of large "peaking factors," so that, as explained above, only colder ions would find themselves in the region of appreciable ripple.

All three of these loss processes can be made to act in concert, so that the prompt removal of decelerated ions does seem feasible. Nevertheless, this requirement seems likely to be the most important difficulty facing successful implementation of the counterstreaming-ion plasma.

3.3 MAINTENANCE OF TOROIDAL CURRENT

The counterstreaming-ion distributions have currents associated with them, but it is easily shown that when $Z_{\text{eff}} = 1$, the total plasma current is zero unless an external electric field $\vec{E}$ is applied. We suppose that the deuterons are injected oppositely to $\vec{E}$, so that the deuteron streaming velocity is in the same direction as the electron drift velocity. The steady-state electron momentum-balance equation is

$$n_e e E + K n_D n_e (v_D - v_e) - K n_T n_e (v_e + v_T) = 0 \quad (25)$$

where $K = \text{constant} \times \ln \Lambda_e / T_e^{3/2}$ and v_e , v_D , v_T are the magnitudes of the average streaming velocities of the electrons, deuterons, and tritons, respectively. From Eq. (25), we have

$$v_e = \frac{1}{n_D + n_T} \left(\frac{eE}{K} + n_D v_D - n_T v_T \right) \quad (26)$$

The plasma current density is

$$J = -e(n_e v_e - n_D v_D + n_T v_T) \quad (27)$$

If $z_{eff} = 1$, then $n_e = n_D + n_T$, and

$$J = -\frac{e^2 E}{K} \quad (28)$$

If the toroidal ion current happens to be unbalanced, the electron velocity distribution shifts so that momentum balance of electrons is always maintained, and net current can be driven only by an external field.

(While the electric field tends to accelerate the tritons, solution of the momentum balance equations for the ions shows that this effect is overwhelmed by the electron drag on the tritons, so that continuous injection is required to maintain the triton drift velocity [17].)

If the plasma has an impurity population, or if a significant fraction of the electrons are trapped in banana orbits or in toroidal field ripple, then $z_{eff} > 1$, and a toroidal current can be maintained even when $E = 0$ [17]. (A similar situation can be established in a beam-heated thermal reactor [18].) Thus the counterstreaming-ion tokamak plasma can be operated for extremely long pulse times, with the volt-sec requirement of the

transformer core determined only by current startup. Since the refueling of the plasma is carried out solely by the injected beams, it appears that virtually steady-state operation is attainable.

The ohmic power dissipation P_{OH} for typical conditions is in the range $0.01-0.03 \text{ W/cm}^3$, so that $P_{OH}/P_b \sim 10^{-2} - 10^{-1}$. Essentially all the power input to the electrons comes from Coulomb energy transfer from the fast ions.

3.4 IMPURITY IONS

The principal effect of an impurity population in the CBT is enhanced pitch-angle scattering, which undermines the advantage of head-on collisions. A second deleterious effect is the additional Coulomb drag on the energetic ions by the impurity ions and their neutralizing electrons. For high-Z impurities, the second effect is much less damaging than the first. Preliminary results in Ref. [4] demonstrated that for $W_0 = 60 \text{ keV}$, $T_e = 5 \text{ keV}$, and iron impurity, Q is reduced by 33% and 46% for $Z_{eff} = 2$ and 3, respectively. Pitch-angle scattering can be moderated by going to higher injection energy, but again this leads to some loss in Q and P_f even when $Z_{eff} = 1$.

Radiation losses are enhanced by a high-Z impurity population, but increased radiation does not significantly affect Q or P_f in a purely beam-driven reactor [3], and can even be desirable for regulating τ_E . The penetration length of neutral beams is reduced in the presence of impurities, since the rate coefficient for impact ionization by oppositely travelling ions is proportional to Z_{eff} [19]. This effect has also the beneficial result of reducing the charge-exchange energy loss associated with trapping by the co-streaming ions.

While the impurity population can presumably be minimized by means of a poloidal divertor, which is required in any event (Sec. 3.2), it is relevant to note that $Z_{\text{eff}} \approx 1$ plasmas (after correction for electron trapping) have been attained in both the Alcator [20] and ATC [21] tokamaks, without the use of a divertor. In each case, however, the plasma energy density was much smaller than we envisage here.

4. PLASMA PARAMETERS FOR REACTOR APPLICATIONS

4.1 NEUTRON WALL LOADING

The confinement requirements of electrons and ions in CBT plasmas are sufficiently lenient, at least at low T_e , that one can conceive of $Q \geq 1$ reactors with thermal power outputs in the 5-50 MW range. But very large power output may be required for economic reasons, and a large neutron wall loading may be demanded for such applications as nuclear transmutation or materials testing. This section is particularly concerned with the problem of attaining large wall loading.

Given T_e and $W_o = W_{oD} = W_{oT}$, the fusion power density for a plasma of arbitrary cross-section is

$$P_f = 4.34 \times 10^{-4} P_{fo} \left(\frac{B_t}{qA} \right)^4 \beta_p^2 S^4 \quad (29)$$

where P_{fo} is the value given in Fig. 9, and S = plasma circumference/ $2\pi a$. In the case of a circular plasma of radius a , with uniform n_e and T_e , and a circular vacuum chamber of radius

a_w , the neutron wall loading is

$$WL = 0.40 P_f \left(\frac{a}{a_w} \right) a \quad (30)$$

If W_0 is fixed, the plasma radius cannot be arbitrarily increased in order to increase WL, because of the limited penetration of the neutral beams. Figure 14 shows WL as a function of W_0 with the plasma pressure fixed. For each W_0 , the plasma radius is $a = \lambda_t$, so that penetration is always satisfactory [12]. We take $a_w/a = 1.3$, which allows adequate space for the divertor scrape-off region. If T_e is small, WL attains its maximum value at $W_0 \approx 100$ keV, while if T_e is large, WL is maximum at $W_0 \approx 60$ keV. These plasmas are all relatively small, with $a < 100$ cm. Increasing the plasma size beyond $a \approx 100$ cm (and therefore increasing W_0) results in a larger power output, but the wall loading is reduced. It should also be kept in mind that the efficiency of neutral beam production tends to decrease at larger beam energies.

The parameters on the right side of Eq. (29) may be varied to increase P_f , but such variation generally implies an increase in n_e , so that λ_t and a must be reduced. Thus substantial changes in WL from the values shown in Fig. 14 are difficult to achieve. It appears very unlikely that a wall loading greater than 0.50 MW/m^2 can be attained.

In practice, one operates at the largest B_t that is technologically feasible, and the largest β_p and smallest q that are permitted for a satisfactory plasma equilibrium (e.g., $\beta_p < A$, $q \geq 2.5$). The plasma aspect ratio A is determined by (I) the sum

of the thicknesses of the transformer core, the toroidal field coils, and the inner blanket-shield, and (II) the minimum plasma radius required for the desired power output or wall loading. (In general, this latter requirement by itself will be such that the rather lenient conditions on τ_E and τ_h will be met.) Aspect ratios smaller than 3.5 are unlikely, because of dimension (I) and the maximum useful plasma radius (Fig.14), while aspect ratios greater than 7 are unlikely because of dimension (II).

4.2 CHOICE OF INJECTION ENERGY

This section considers the trade-offs between low and high injection energy, W_0 . First, consider reactors with large power ratings ($P_{\text{thermal}} \geq 100$ MW). In general, $\bar{n}_e a$ must increase for reactors of increasing power output, so that the largest W_0 are required in order to insure penetration of the injected neutrals. But Q decreases for $W_0 \gtrsim 60$ keV, and the neutron wall loading eventually decrease also in very large devices (cf. Fig. 14). For these reasons, there may be an economic penalty in going to a very large reactor (say, $\bar{n}_e a > 5 \times 10^{15} \text{ cm}^{-2}$), which necessarily requires a large W_0 .

In the case of reactors with small to moderate power ratings, there is always some leeway in choosing W_0 for tangential injection into a plasma of given radius, since acceptable trapping profiles are obtained for the range $0.4a \lesssim \lambda_t \lesssim 1.2a$ [12]. This margin is not as large as apparent, however, since the allowable n_e tends to decrease with increasing W_0 . The advantages of operating at the maximum end of the W_0 range are the following:

- (1) The charge-exchange loss, F_{cx} , decreases with increasing W_o .
- (2) If $Z_{eff} > 1$, angular scattering is reduced at larger W_o .
- (3) If a substantial cold-ion population cannot be avoided, at large W_o one can take better advantage of target-plasma reactions, so that Q-values will not be as severely degraded.
- (4) For the same injection power, less beam current is needed at larger W_o , which could mean an amelioration of the difficult pumping problems in the beam ducts (if neutralization efficiency can be maintained).

An important figure-of-merit is the fractional burnup of tritium per pass. A large fractional burnup implies a small tritium throughput per unit thermal power production, so that the tritium inventory can be smaller, and the demands on the tritium processing system are reduced. If f_B is the fractional burnup for tritium (once-through), then

$$\frac{n_T}{\tau_{hT}} f_B E_f = P_f = QP_b \quad (31)$$

Using Eq. (14) and $\tau_{hT} = 1.5 \tau_{hD}$, we have

$$f_B = \frac{Q W_o}{E_f} \left(1 + \frac{\tau_{hT}}{\tau_{hD}} \right) = 2.5 \frac{Q W_o}{E_f} \quad (32)$$

Figure 15 shows f_B as a function of W_O and T_e . At constant W_O , the increase in f_B with T_e is due to the increase in fast-ion slowing-down time (and concomitant increase in Q). At a given T_e , f_B increases with W_O , reflecting (1) a somewhat longer τ_h for more energetic ions, and (2) a larger average σv over the lifetime of these ions. Furthermore, with increasing W_O , a given T_e is attained at a smaller value of $n_e \tau_E$ (cf. Fig. 5).

Most of the above considerations would seem to indicate that in reactors of modest size, the energy range 50 to 80 keV is generally to be preferred to the range 30 to 50 keV. But in the end, the choice of injection energy may very well depend on the efficiency of neutral beam production, which tends to be larger at lower energies. Overall efficiencies of 60% (busbar-to-beam) have been achieved for D^0 beams at 40 keV (although present-day beams contain appreciable quantities of D_2^0 and D_3^0). While the neutralization efficiency of D^+ decreases markedly for $W_O \gtrsim 50$ keV, satisfactory overall efficiency can in principle be maintained even for $W_O \gg 50$ keV by the use of direct conversion of unneutralized ions, but with the penalty of increased unit capital cost compared with that for low-energy injection.

4.3 RADIATION TEST FACILITY

The behavior of wall materials and blankets subject to intense bombardment of 14-MeV neutrons must be investigated in the course of developing fusion reactors. As noted in Sec. 4.1, the neutron wall loading in a CBT reactor is unlikely to exceed 0.5 MW/m^2 , so that a CBT plasma cannot profitably be employed to study neutron-induced bulk-radiation effects, which require wall

loadings in the range 1 MW/m^2 and above [22]. (A target-plasma reactor has been proposed for this purpose [23].) On the other hand, the CBT does seem suitable for studies of neutron-induced radiation damage to surfaces, simultaneously with studies of the effects of energetic particle and photon bombardment.

For this purpose one requires a neutron flux of at least $3 \times 10^{12} \text{ n/cm}^2/\text{s}$ (0.07 MW/m^2), with a minimum fluence for a given test sample in the range $10^{19} - 10^{20} \text{ n/cm}^2$ [22]. The neutron source must be capable of a high duty factor, to meet this fluence requirement. It is worth emphasizing that the large test area inherent in any tokamak device insures that a large number of surface radiation experiments can be simultaneously accommodated.

The special requirements of a test reactor for surface radiation effects are (1) time-averaged neutron flux $\geq 3 \times 10^{12} \text{ n/cm}^2/\text{s}$; (2) as small a total power input as possible; (3) $Q \geq \frac{1}{2}$, in order to minimize the capital cost of the beam injectors, the number of beam lines, and the circulating power; (4) ease of access for replacement of test sections. The test reactor would be driven by an external power source, and if copper TF coils are used, these would in fact account for most of the power loss. No blanket is required, although suitable means for removing the thermal power generated (injected power and fusion power) must be provided. Tritium must be supplied from an external source.

Illustrative plasma parameters are presented in Table I. The modest elongation of the plasma cross-section is compatible with the provision of a poloidal divertor. Water-cooled-copper TF coils are specified, since perhaps only these are technologically feasible for a near-term toroidal device, and since

replacement of test specimens and diagnostics is facilitated with room-temperature coils. While an injection energy of 40 keV would be adequate, $W_0 = 60$ keV is specified for the reasons discussed in Sec. 4.2. An ideal Q larger than unity is probably required, because of the deleterious effects of charge-exchange losses, impurities, and excessive population of cold ions. The time-averaged neutron flux (50% duty factor) is 0.10 MW/M^2 , and the useful test area (on the outer part of the torus vacuum vessel) is 25 m^2 .

Long-pulse, high-duty-factor, neutral beam operation is essential for this and other CBT reactor applications. Present 40-keV injectors operate with 0.3-s pulses and duty factors of 15% [24]. It appears that steady-state operation should be feasible with large-area beams, and a 50-keV, 50-A, 2-s source is presently under development [24].

4.4 FUSION-FISSION HYBRID REACTORS

The 14-MeV neutrons from the DT reaction can be used to breed fissile material (^{233}U , ^{239}Pu) in a fertile blanket (^{232}Th , ^{238}U) surrounding the fusion plasma, and the neutron energy can be multiplied by fission reactions in this blanket [2]. In the following, we use the term hybrid reactor for any fusion-driven device having a blanket containing fertile material. The hybrid can be an economic power producer even when Q of the fusion core is smaller than unity, because of the large energy gain ($\gtrsim 10$) attainable in the blanket [2]. Alternatively, the fissile products may be removed and utilized in thermal fission reactors, a small fraction of the electrical output of which may be diverted

to make up the power deficit of the fusion reactor (a "symbiotic" system [2]). It was shown in Ref. [3] that for a beam-driven fusion reactor with negligible fission power produced in the blanket, $Q \gtrsim 0.75$ is needed for satisfactory overall power balance of the symbiotic system. (Defrayment of the capital cost of the fusion reactor requires another condition, however.)

Whether or not symbiotic operation is employed, the special requirements for the plasma of an economic hybrid reactor are (1) large power output and large duty factor, so that the annual production of fissile material and energy provide a sufficient return on the total capital investment; (2) large wall loading, to maximize the cost effectiveness of the blanket; (3) $Q \sim 1$, or smaller if the blanket has a large power multiplication. In addition to production of fissile material, the blanket must have a tritium breeding ratio of at least unity. If the DT neutron energy is not moderated, then a total thickness of 1.3 m for the blanket and shield is adequate.

The preferred plasma operating point is the large-radius end of the plateau of maximum neutron wall loading (Fig. 14), in order to meet conditions (1) and (2) above. The reduction in Q at large W_0 can be tolerated, and moderate T_e is satisfactory. Illustrative plasma parameters are given in Table I. A 30-cm scrape-off layer is provided between the plasma edge and the vacuum wall. In order to attain the highest possible toroidal field on axis, Nb_3Sn coils are specified. The injection energy of 100 keV affords adequate penetration of the plasma, whose radius is chosen to provide the largest possible wall loading for the allowable plasma pressure.

($\beta_p = 2/3$ A at $q = 3$). The transformer core is capable of 18 V-sec, which is sufficient to drive 2.6-MA pulses up to 50 s duration; much longer pulses are possible with the same core, if the plasma current can be maintained by the beams [17]. The neutron power output of 103 MW is convenient for electrical generating plants of moderate size. For fast-fission blankets such as considered in [25], and 75% duty factor, the production rates (net) of fissile material and electrical power would be in the range 2-3 kg/day and 150-900 MW, respectively, with a tritium breeding ratio of at least 1.05.

4.5 TRITIUM PRODUCTION

Tritium is presently used as a reactant in thermonuclear explosives, and for the production of ^3He by beta-decay. If the controlled fusion program continues to develop, one may expect large amounts of tritium to be required for fusion test reactors, and eventually for the start-up of a fusion power economy. In the case of a system of hybrid reactors, it may turn out to be most convenient to relegate the function of breeding tritium to separate reactors, in order to better utilize the hybrid neutrons for fissile breeding.

A DT fusion reactor can produce considerable amounts of tritium for sale, because with appropriate neutron-multiplying materials in the blanket, such as sodium or beryllium, a tritium breeding ratio of 2.0 is attainable [26], with a doubling time of the order of several months or less. If tritium production is the principal purpose of the reactor, economic viability demands large power output per unit size, large wall loading, and a large duty factor. The reactor should also have overall electrical

energy break-even — the high cost of externally-supplied electricity would probably make tritium production prohibitively expensive. With superconducting TF coils and efficiently generated neutral beams, $Q \gtrsim 2.5$ is required, considering that each neutron may generate up to an extra 10 MeV in an optimized blanket [26].

Illustrative plasma parameters for a CBT "tritium factory" are shown in Table I. Attaining $Q > 2$ demands a large T_e , and correspondingly large $n_e \tau_E$, so that a larger plasma radius is specified than in the hybrid case. The 125-MW neutron power is near the maximum one can reasonably produce in a CBT reactor with $W_0 \lesssim 100$ keV. The annual net tritium production for operation at 80% duty factor is 7.0 kg.

Table I

ILLUSTRATIVE PARAMETERS FOR COUNTERSTREAMING-ION TOKAMAK REACTORS

	Radiation Test Facility	Hybrid Reactor	Tritium Producer
R_o (m)	2.45	4.05	4.20
a (m)	0.50	0.75	0.90
A	4.9	5.4	4.7
Vertical elongation	1.25	1.25	1.25
Midplane wall radius (m)	0.75	1.05	1.20
TF coil material	Cu	Nb_3Sn	Nb_3Sn
B_t on axis (kG)	40*	55	53
B_t at coil (kG)	82*	130	130
Blanket/shield thickness (m)	0.5	1.3	1.3
I_p (MA)	1.4	2.6	3.5
q	3.0	3.0	3.0
$\bar{n}_e$ (cm^{-3})	4.3×10^{13}	4.8×10^{13}	4.3×10^{13}
$\bar{T}_e$ (keV)	4.0	5.0	8.0
$\bar{n}_e \tau_E$ ($cm^{-3}s$)	2.1×10^{12}	2.3×10^{12}	6.5×10^{12}
W_o (keV)	60	100	100
λ_t/a	0.92	0.75	0.72
Injection power (MW)	19	97	68
β_p	3.3	3.6	3.1
Q (ideal)	1.32	1.33	2.3
P_f (W/cm^3)	1.28	1.81	1.43
Total neutron flux (n/s)	8.9×10^{18}	4.6×10^{19}	5.5×10^{19}
Neutron wall loading (MW/m^2)	0.20	0.45	0.44
Total neutron power (MW)	20	103	125

* duty factor = 50%

5. CONCLUSIONS

This paper has considered those properties of a beam-driven counterstreaming-ion tokamak plasma (CBT) that are most relevant for application to DT fusion reactor design. By postulating that in steady-state there exists an equipartition of energy between the directed ion motion and the total plasma thermal motion, ion velocity distributions and plasma properties are derived that are in good agreement with known Fokker-Planck results. Fusion reactivity, power gain Q , and power density P_f have been calculated for wide ranges of injection energy and electron temperature, and the results demonstrate a trade-off between Q and P_f .

The principal practical problem in implementing a counterstreaming-ion plasma in a tokamak is the need to minimize the cold-ion population. Decelerated ions ($W \lesssim 2 T_e$) are removed chiefly by instability-induced diffusion, but also by charge-exchange on injected neutrals, and by diffusion induced by magnetic field ripple. A poloidal magnetic divertor is probably required to minimize ion recycling, to remove impurities from the torus, and to reduce bombardment of the wall by ions leaving the discharge. The trapping process of injected neutrals results in less than 10% energy loss by charge exchange. Impurities reduce Q mainly by enhancement of angular scattering, thus undermining the advantage of head-on nuclear collisions.

The counterstreaming-ion system seems advantageous for a low-power, high-duty-factor, $Q \approx 1$ neutron source. Injection energies are modest (40-60 keV); the electron energy confinement

time can be small (e.g., $n_e \tau_E \approx 2 \times 10^{12}$ at $T_e \approx 3.5$ keV); refueling is accomplished entirely by the injected beams; and the plasma current can be maintained by the beams. To insure adequate penetration by neutral beams, the maximum useful size is apparently significantly smaller than that of a TCT (target-plasma) reactor, and many times smaller than that of a thermonuclear reactor.

Plasma parameters were considered for three applications of CBT-generated neutrons: a 20 MW (neutron power) surface radiation effects reactor, with a time-averaged neutron flux of 5×10^{12} n/cm²/sec; a 100 MW fusion-fission hybrid reactor; and a 125 MW tritium producer. The lenient requirements on beam voltage and $n_e \tau_E$ would seem to make these reactor applications feasible in the near future.

ACKNOWLEDGMENTS

The author acknowledges many useful discussions with Dr. R. M. Kulsrud and Mr. Y. C. Sun concerning their Fokker-Planck results. This work was supported by the United States Energy Research and Development Administration Contract E(11-1)-3073.

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Fig. 1. Counterstreaming-ion velocity distributions parallel to the toroidal magnetic field. The two distributions have the same density, temperature, and total energy.

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Fig. 2. Contours of displaced-Maxwellian velocity distributions of counterstreaming D and T ions. The large dots indicate $f(v) = 1.0$. The two distributions have the same temperature and average energy.

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Fig. 3. (a) Mean parallel velocity, (b) temperature, and (c) average energy of the displaced-Maxwellian deuteron velocity distribution in a counterstreaming-ion plasma with injection energy of 60 keV. The large dots in (c) are the results of a Fokker-Planck calculation in Ref. [4].

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Fig. 4. Ratio of ion lifetime to electron energy confinement time in a counterstreaming-ion plasma as a function of T_e , for various injection energies W_0 .

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Fig. 5. Dependence of electron energy confinement time on W_0 and T_e . The large dots are the results of a Fokker-Planck calculation for $W_0 = 60$ keV, in Refs. [1,4]. The TCT case is for maximum fusion power density [3].

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Fig. 6. Fusion reactivity of D and T displaced-Maxwellian distributions, each of temperature T_b , with relative displacement velocity v_r . The maximum reactivity of a DT beam-target system ($W_{\text{beam}} = 125$ keV) is denoted by A. The maximum reactivity of a thermonuclear DT plasma ($T_i = 60$ keV) is denoted by B.

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Fig. 7. Fusion power gain Q for steady-state counterstreaming-ion operation in DT, taking 17.6 MeV per reaction. The large dots are the results of a Fokker-Planck calculation in Refs. [1,4].

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Fig. 8. Fusion power gain Q for steady-state counterstreaming-ion operation in DT, as a function of injection energy and T_e . 17.6 MeV per reaction.

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Fig. 9. Fusion power density P_f for counterstreaming-ion operation in DT, as a function of injection energy and T_e . Total plasma pressure = 0.19 J/cm^3 , corresponding to tokamak parameters $B_t = 50 \text{ kG}$, $A = 5.0$, $\beta_p = 3.0$, $q = 2.5$.

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Fig. 10. Q versus P_f for a counterstreaming-ion DT plasma. The solid curves are contours of constant injection energy, and the dashed curves are contours of constant $n_e \tau_E$. Total plasma pressure = 0.19 J/cm^3 .

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Fig. 11. Rate coefficients for charge-exchange, ionization by plasma ions, and electron ionization for a deuterium atom in a hydrogenic plasma.

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Fig. 12. (a) Trapping length λ_t of deuterium atoms of energy W_0 injected into a counterstreaming-ion plasma formed by injection at W_0 . R_{cx}^0 is the fraction of atoms trapped by charge exchange. (b) Mean-free-path λ_n for charge exchange or ionization of a deuterium neutral of energy W_n in a counterstreaming-ion plasma formed by injection at 60 keV. R'_{cx} is the relative probability of a charge-exchange collision.

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Fig. 13. Plan view of geometry of tangential neutral-beam injection into a torus.

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Fig. 14. Neutron wall loading in a counterstreaming-ion tokamak reactor. Plasma pressure = 0.19 J/cm^3 , and plasma radius = neutral-beam trapping length. Reactors to the left of the dotted line have plasma radius $< 100 \text{ cm}$.

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Fig. 15. Fractional burnup of tritons per pass in a counterstreaming-ion DT reactor, as a function of T_e and injection energy $W_{OD} = W_{OT}$.

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