Neutral-Beam-Injected Tokamak Fusion Reactors: A Review

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A REVIEW

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ABSTRACT

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The ion velocity distributions of tokamak plasmas sustained by injected neutral-atom beams fall into three categories, according to the partial pressure of the energetic-ion population, which in turn is determined by the plasma nT_E , the injection energy, and the fraction of plasma fueling performed by the beams. There are correspondingly three regimes of beam-driven reactor operation in D-T or D-D: thermonuclear; beam-target or TCT; ion pressure/electron pressure $\gg 1$. This review treats the physical basis and plasma engineering parameters of beam-driven tokamak reactors.

The theories of energetic-ion velocity distributions, stability, injection, and orbits are summarized. The many-faceted role of the energetic ions in plasma heating, fueling, and current maintenance, as well as in the direct enhancement of fusion power multiplication and power density, is discussed in detail for the three reactor types. The relevant implications of recent experimental results on several beam-injected tokamaks are examined. The behavior of energetic ions is found to be in accordance with classical theory,

large total ion energy densities are readily achieved, and plasma equilibrium and stability are maintained.

The status of neutral-beam injectors and of conceptual design studies of beam-driven reactors are briefly reviewed. The principal plasma-engineering problems are those associated directly with achieving quasi-stationary operation. An especially attractive feature of beam-driven tokamak plasmas is their capability for useful application in devices of relatively small size.

CONTENTS. 1. Introduction. 2. Ion velocity distributions. 3. Velocity-space stability. 4. Injection and orbits of energetic ions. 5. Heating, fueling, and current maintenance by beam-injected ions. 6. Performance of beam-injected tokamak reactors. 7. Injector systems. 8. Experimental results. 9. Conceptual designs of beam-driven reactors. 10. Conclusions and outlook.

Contents of Review Paper

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1. INTRODUCTION

The present interest in beam-driven toroidal reactors has its origins in target-plasma schemes that date back to the earliest days of the controlled fusion program. Because the D-T fusion reaction has its maximum cross section in an energy range that is easily accessible to the production of high-current ion beams, it might appear that useful fusion energy can be generated by bombarding a solid tritide target with energetic deuterons, or perhaps by directing energetic D and T beams at each other. However, as realized very early upon examination of the solid-target scheme, the slowing-down of energetic ions by ionization of target atoms and by Coulomb interaction with the cold target electrons is so rapid that less than 1% of the beam energy can be recovered in fusion energy [1-3]. (Nevertheless, solid-target bombardment remains to this date the most useful method for obtaining a steady flux of 14-MeV D-T neutrons.) The colliding-beam system suffers from the low ion density available in accelerator beams, as well as from the dominance of the Coulomb angular scattering cross section for ions in opposing beams over the fusion-reaction cross section [1-3].

Ionization losses can be eliminated by using a fully ionized plasma target, and Coulomb drag can be minimized by increasing the electron temperature. In the mid-1950's, proposals appeared for operation of magnetic-mirror machines at $T_e \geq 1$ keV, into which ion beams would be injected and trapped [4,5]. At that time a principal motivation for such schemes, which appeared to offer only limited fusion energy multiplication, was the production of large fluxes of 14-MeV neutrons, to be used for breeding ^{239}Pu and ^{233}U in fertile blankets surrounding the plasma. Interest in target-plasma schemes dwindled in the late 1950's, when it became evident that fissile material could be produced in adequate quantities by other methods. (In this

period, however, proposals first appeared for the injection of neutral beams into magnetic traps [6].) In the 1960's, it was realized that energetic ions produced in nominally thermal plasmas by instabilities, for example, could have a large integrated fusion probability while thermalizing with the bulk plasma, but this process was dismissed as impractical for the production of appreciable fusion power [7].

In the period 1968-1971, two developments led to revived interest in target-plasma schemes. The first was the general recognition that the Soviet-developed tokamak device was capable of maintaining high- T_e plasmas for appreciable energy confinement times. The second was the astounding progress being made in the development of neutral-beam injectors, to be used for injection of large fluxes of energetic particles into both mirror and toroidal containment devices. While the original intent of combining these two technologies was the heating of tokamak plasmas [8], it was soon realized [9] that if the temperature of a tritium target plasma were maintained by injected deuteron beams, then for sufficiently large T_e the beam-plasma system could have an energy multiplication factor greater than unity, with a required density-confinement-time product well below the "Lawson-breakeven" value. The required size of the tokamak, and the energy and intensity of the injected beams, appeared to be attainable with reasonable extrapolation of the state-of-the-art. More recently the old colliding-beam idea has also been revived, with the realization that in a neutralized toroidal plasma the densities of oppositely traveling deuterons and tritons can be stacked by neutral-beam injection to very high values [10]. In a toroidal device, of course, ions can be confined for many 90° scattering times.

The first injection of neutral beams into tokamak plasmas took place at the Culham, Princeton, and Oak Ridge Laboratories in 1972-3. The

behavior of the injected ions was found to be classical, and the expected plasma heating and D-D neutron production were observed. Since then, more powerful injection experiments on a number of tokamaks have resulted in substantial plasma heating, with apparently no deleterious effect on the plasma equilibrium or on the fast-ion slowing-down rate. Although experiments have been in progress for only 3 years, the interaction of injected energetic ions with the tokamak plasma ranks today among the best understood of all tokamak phenomena. This favorable performance, and the continuing progress in the development of more powerful, higher voltage, neutral-beam injectors, have led to the present confidence in the eventual success of beam-driven tokamak reactors.

This review treats the physical basis and plasma-engineering parameters of such reactors. Sections 2 to 4 summarize the theory of energetic-ion velocity distributions, stability, injection, and orbits. Section 5 discusses the heating, fueling, and current-driving functions of injected ions. Section 6 describes the detailed theoretical performance of the various types of beam-driven reactors. Neutral-beam injectors are discussed in Sect. 7. The significance of relevant experimental results obtained to date is discussed in Sect. 8. Conceptual design studies are surveyed briefly in Sect. 9.

The reader is referred to recent review articles [11,12] for summaries of the theoretical and experimental bases of tokamak operation.

1.1 SURVEY OF BEAM-INJECTED TOKAMAK REACTORS

Several methods have been proposed for heating D-T tokamak plasmas to temperatures of reactor interest [11,12]. The unique feature of plasma heating by means of injected high-energy deuterons or tritons is that these ions produce fusion energy directly, while thermalizing with the bulk plasma.

At the same time the injected ions can contribute substantially to plasma fueling. The relative importance of these three roles of neutral-beam-injected ions defines the principal regimes of reactor operation.

(1) The TCT (two-energy-component torus) consists of a near-Maxwellian "warm" target plasma whose temperature is maintained against transport and radiation losses by injected energetic ions [9,13]. Fusion energy is produced principally by reactions between the energetic ions and the target ions. When the bulk-ion lifetime is substantially greater than the plasma energy confinement time, the injected ions can provide significant fueling. The ion velocity distribution in a TCT is shown schematically in Fig. 1(a).

(2) The BDTM (beam-driven thermonuclear reactor) consists of an essentially Maxwellian bulk plasma whose temperature is maintained by injected energetic ions together with charged fusion-reaction products. Fusion energy is produced principally by thermonuclear reactions among the "warm" plasma ions. The bulk plasma is fueled almost entirely by pellet injection, or gas injection, or by plasma recycling. The ion velocity distribution is shown in Fig. 1(b). In this regime the injected energetic ions play a role similar to that of the charged fusion-reaction products.

(3) The CIT (counterstreaming-ion torus) consists of two nearly thermal distributions of ions, oppositely displaced in velocity along the magnetic axis [10,14], as shown in Fig. 1(c) for a D-T system. These distributions are maintained by beam injection parallel and antiparallel to the magnetic axis as indicated, and by minimizing the influx of cold neutrals and plasma into the hot reacting region. Decelerated ions must have a lifetime substantially smaller than that of the energetic ions. The electron temperature, T_e , is maintained by Coulomb power deposition from the energetic ions. Fusion energy is produced by nearly head-on collisions between oppositely

streaming D and T ions. All fueling is carried out by the injected ions, and maintenance of the plasma current by these ions is also possible under certain conditions. (The CIT can be regarded loosely as a type of Galilean transformation of the TCT, or else as a modified "thermonuclear" system, in which the thermal D and T components are separated in the laboratory frame of reference. The CIT is sometimes referred to as the CBT, for "colliding-beam torus.")

These three reactor regimes merge into each other, with the change of one or two parameters. One key parameter is $n_e \tau_E$, the product of electron density and bulk-plasma energy confinement time. With increasing $n_e \tau_E$, less injected beam power is required to sustain the plasma temperature, and the ratio of energetic-ion pressure to bulk-plasma pressure decreases. As shown schematically in Fig. 2, the TCT regime merges into the beam-driven thermonuclear regime, and finally into the ignited thermonuclear regime, where the plasma temperature is maintained only by charged fusion-reaction products [13].

The relation of the CIT regime to the others is more subtle, and is determined by the residence time of ions in the bulk plasma, the amount of particle recycling from the wall into the reacting plasma, and the ratio of T_e to the injection energy. At low $n_e \tau_E$ and therefore low T_e , with increasing bulk-ion population beam-target reactions become significant, and the CIT evolves into a TCT with 50:50 D-T injection. At large $n_e \tau_E$, which generally implies large T_e , ions in the counterstreaming distributions tend to thermalize with each other before slowing down. Then with increasing bulk-ion population, the CIT evolves into a beam-driven thermonuclear reactor, with ion temperature $T_i > T_e$.

The principal distinctions among the three modes of beam-driven reactor operation are summarized in Table 1. The primary attraction of the TCT and CIT regimes is their large fusion reactivities, which make possible large

fusion power densities in tokamaks of relatively small size. On the other hand, the TCT and CIT plasmas are heavily beam-driven, so that their fusion power multiplications (Q) tend to be much smaller than those attainable in the beam-driven thermonuclear regime.

It is worth noting the close relation of TCT and CIT reactors to beam-injected mirror-machine operation — and in fact neutral-beam injection was proposed first for application to magnetic mirrors [6]. As in conceptual mirror reactors, the CIT employs beams both for fueling and heating, but it is capable of much larger Q -values. The greater incidence of head-on collisions results in increased reactivity, and in a torus loss cones in velocity space can be largely eliminated. The absence of the loss-cone problem is also the principal theoretical advantage of the TCT over the mirror machine. More recently a "two-component mirror" has been proposed in which the energetic ions are confined by magnetic mirrors, but the bulk-plasma target flows unimpeded along the field lines [15]. For very large mirror ratios the calculated Q -values are similar to those of the TCT.

2. ION VELOCITY DISTRIBUTIONS

2.1 TYPES OF VELOCITY DISTRIBUTIONS

Energetic neutral-atom beams injected into a tokamak plasma supply the plasma with energetic ions that serve a number of purposes: (1) They supply energy to the bulk-plasma electrons and ions; (2) they supply particles to the plasma, thereby assisting in fueling; (3) they supply momentum, that can help maintain the plasma current; (4) they may undergo nuclear reactions with each other, or with ions of the bulk plasma. The density and velocity distribution $f_h(v)$ of the energetic ions determine the degree to which each of the four principal roles is carried out — and therefore determine the regime of reactor operation. The fast ions decelerate by Coulomb interaction with the bulk-plasma particles, and in the steady state form velocity distributions whose nature depends on the intensity of beam injection, as well as on the particle and energy confinement times of the bulk plasma. Section 4 deals with the spatial neutral-beam deposition and the orbits of injected ions. This section is concerned only with the possible forms of $f_h(v)$, and the effect of the energetic ions on the velocity distribution $f_i(v)$ of the bulk-ion population.

We assume that fast ions are retained in the plasma until they decelerate to a certain energy W_1 in one "slowing-down time" τ_s . Ions with energy $W < W_1$ are referred to as "warm" ions and have a lifetime τ_p . In the following, the subscript "h" denotes the energetic or hot ions, while the subscript "i" denotes the warm or thermal ions, which can be more or less in equilibrium with the plasma electrons. We also define two parameters that relate energy densities:

$$\Gamma = \frac{\text{hot-ion energy}}{\text{bulk-plasma energy}} = \frac{n_h \bar{W}_h}{\frac{3}{2} n_e T_e + n_i \bar{W}_i} \quad (2.1)$$

and

$$\Pi_i = \frac{\text{total ion energy}}{\text{electron energy}} = \frac{(n_h \bar{W}_h + n_i \bar{W}_i)}{\frac{3}{2} n_e T_e} \quad (2.2)$$

The quantity Γ is most useful when $n_h \ll n_i$, which generally implies that $\Gamma \gg 1$, and $\bar{W}_i \approx 3/2 T_e$. When there are relatively few ions in equilibrium with the electrons, it is more meaningful to use Π_i .

There are three general classes of velocity distributions, shown schematically in Fig. 1, which correspond to the three reactor regimes discussed in Sect. 1.2.

(1) $\Gamma \ll 1$ (BDTN). A large bulk-ion population is in equilibrium with plasma electrons. The relatively small energetic-ion population (injected ions or charged fusion products) have negligible Coulomb interaction with each other, and have negligible effect on the Maxwellian nature of f_i .

(2) $\Gamma \sim 1$ (TCT). A fairly large energetic-ion population causes significant distortion of $f_i(v)$, but the interaction among energetic ions can be treated adequately as a perturbation.

(3) $\Pi_i \gg 1$ (CIT). The energetic-ion population is comparable with or larger than the warm-ion population. Coulomb interactions among the hot ions must be treated accurately. There are essentially no ions in equilibrium with the plasma electrons. Under some conditions, it is most appropriate to describe the entire ion population of a given species by a single $f_i(v)$.

Table 2 shows the principal functions of the energetic ions in each of the three plasma regimes. In practice, the ion velocity distribution is determined partly by the fraction of the plasma ion population that is supplied by injected beams, and partly by the ratio of the fast-ion slowing-down time τ_s to the warm-ion confinement time τ_p . We can define a beam-fueling fraction $f_F = (n_h/\tau_s)/(n_i/\tau_p)$. (1) If plasma recycling is negligible, then beam injection can be the dominant fueling mechanism ($f_F \sim 1$). If τ_p/τ_s is very

small, then $\Pi_i \gg 1$. (2) If recycling is moderate, then injected beams can supply a significant fraction of the plasma ion population ($f_F \geq 0.2$). Provided that τ_s is not too small, one has $\Gamma \sim 1$. (3) If the plasma is fueled by recycling and possibly by an additional method such as intense pellet injection ($f_F \ll 1$), then the relative beam-ion population is likely to be very small, so that $\Gamma \ll 1$ even when $\bar{W}_h \gg T_e$.

2.2 VELOCITY DISTRIBUTIONS FOR $\Gamma \ll 1$

In this regime the warm ions can be treated as Maxwellian with temperature T_i , and $f_h(v)$ can be calculated from the drift-kinetic equation with a Fokker-Planck collision operator, neglecting interaction among the hot ions. (The electron distribution is always taken to be Maxwellian with temperature T_e .) Neglecting magnetic trapping effects, the Fokker-Planck equation for fast ions slowing down on Maxwellian background ions and electrons having thermal velocities v_i and v_e , with $v_i < v \ll v_e$, becomes [16-19]

$$\begin{aligned}
 \frac{\partial f_h}{\partial t}(v, \zeta, t) &= S(v, \zeta, t) && \text{(fast-ion source)} \\
 &- f_h / \tau_{cx} && \text{(charge-exchange)} \\
 &+ \frac{1}{\tau_{se} v^2} \frac{\partial}{\partial v} [v^3 + v_c^3] f_h && \text{(electron and ion friction)} \\
 &+ \frac{1}{2\tau_{se}} \left[\frac{v_e^2 m_e}{m_h} + \frac{v_c v_i^2 m_i}{v^3 m_h} \right] \frac{\partial^2 f_h}{\partial v^2} && \text{(energy diffusion)} \\
 &+ \frac{1}{2\tau_{se}} \frac{m_i}{m_h} \frac{Z_{eff}}{[Z]} \frac{v_c^3}{v^3} \frac{\partial}{\partial \zeta} \left[(1-\zeta^2) \frac{\partial f_h}{\partial \zeta} \right] && \text{(pitch-angle scattering)} \\
 &- \frac{eE^*}{m_h} \left[\zeta \frac{\partial f_h}{\partial v} + \frac{(1-\zeta^2)}{v} \frac{\partial f_h}{\partial \zeta} \right] && \text{(E-field acceleration and distortion of } f_h)
 \end{aligned}$$

(Eq. 2.3)

$$Z_{\text{eff}} = \sum_j \frac{n_j Z_j^2}{n_e} \quad \text{and} \quad [Z] = \frac{1}{n_e} \sum_j n_j Z_j^2 \frac{m_i}{m_j} \quad (2.4)$$

where the summation is over all background-ion species, with i denoting the dominant species. In Eq. (2.3), τ_{se}^{-1} is the rate of slowing-down on electrons (τ_{se} is called the "Spitzer slowing-down time" [20]); $\zeta = v_{\parallel}/v$, where $v_{\parallel}$ is the velocity parallel to the magnetic field; $E^* = E_{\phi}(1 - Z_h/Z_{\text{eff}})$, where E_{ϕ} is the tokamak electric field; $\tau_{\text{cx}}(v)$ is the charge-exchange lifetime; S is the injection rate of fast ions; v_c is the "transition" (or "critical") velocity of fast ions at which the electron Coulomb friction equals the bulk-ion Coulomb friction [17-19,21]. The "transition energy" $W_c = \frac{1}{2} m_h v_c^2$ is

$$W_c = 14.8 T_e A_h \left[\frac{1}{n_e} \frac{1}{\ln \Lambda_e} \sum_j \frac{n_j Z_j^2 \ln \Lambda_j}{A_j} \right]^{2/3} \quad (2.5)$$

Here A_h and A_j denote atomic weights, and the Coulomb logarithms are calculated as in Ref. [22]. The transition energy is typically 15-20 T_c .

From Eq. (2.3) one sees that energy drag is dominated by electrons when $v > v_c$ and by bulk ions when $v < v_c$. The angular scattering term is due to collisions with bulk ions and is small when $v > v_c$ (unless $Z_{\text{eff}}/[Z] \gg 1$) but is comparable with the bulk-ion friction term when $v < v_c$. For injection velocity $v_h > v_c$ with $Z_{\text{eff}} \sim 1$, beam ions degrade in energy with little angular scattering until $v \approx v_c$. The ions then isotropize and degrade further in energy at comparable rates.

In the following, we are concerned only with the steady-state $f_h(v)$, which prevails after one slowing-down time. Exact solutions to Eq. (2.3) for a monoenergetic source, but neglecting energy diffusion terms, have been obtained

by a number of authors by expansion of f_h in Legendre polynomials [16-19,23]. Figure 3 shows $f_h(v)$ for energetic deuterons injected at 100 keV ($v_b = 3.1 \times 10^{18}$ cm/s) tangentially to the magnetic axis of a deuterium plasma with $T_e = T_i = 3.0$ keV, $E^* = 0$, $\tau_{cx} = \infty$, and (a) $Z_{eff} = 1.0$, (b) $Z_{eff} = 4.0$. The increased pitch-angle scattering at the larger Z_{eff} is apparent. In this solution, energy diffusion during slowing down is accounted for only in the formation of the tail above the injection velocity. An analysis for isotropic energetic-ion distributions that retains energy diffusion [24] shows that diffusion has a rather small effect on $f_h(v)$ for $v < v_b$, unless T_i is relatively large. Actually, this effect can significantly influence the power gains of driven D-T reactors, when v_b/v_i is not too large (see Sect. 6.3). Omission of energy diffusion leads to errors in Q of order T_i/W_b . The exact solution to $f_h(v)$ for tangential injection, when all the Fokker-Planck diffusion coefficients are retained, has been calculated numerically [25]. The velocity distributions parallel and perpendicular to the magnetic field for $W_b = 200$ keV and $T_e = 10$ keV are shown in Fig. 4.

The effect of the toroidal electric field when the fast ions are injected parallel to E (co-injection) is to enhance $f_h(v)$ at the larger velocities, and to increase the width of the tail above v_b [16]. For oppositely directed ions (counter-injection), $f_h(v)$ is decreased at larger velocities. Charge-exchange collisions increase the probability of ion loss while slowing down, and therefore reduce the population of lower-velocity ions [16,23].

If energy diffusion below the source is ignored, one can derive an expression for $F_h(v)$, the angle-averaged $f_h(v)$, from a simple test-particle picture. This calculation applies regardless of the extent of pitch-angle scattering during the slowing-down process. With a source at velocity v_b and a sink at velocity v_m , the steady-state flow of ions through velocity space $v_m < v < v_b$ is described by

$$F_h(v) v^2 \frac{dv}{dt} = \text{constant} \quad (2.6)$$

The Coulomb deceleration dv/dt is

$$\frac{dv}{dt} = c_1 v \frac{n_e \ln \Lambda_e}{T_e^{3/2}} + \frac{1}{v^2} \sum_i c_{2i} n_i \ln \Lambda_i \quad (2.7)$$

where the summation is over the warm-ion species, and c_1 and c_{2i} are constants. Inserting Eq. (2.7) into (2.6) gives $F_h(v) \propto (v^3 + v_c^3)^{-1}$.

A zero-order approximation to the energy diffusion is obtained by noting that faster members of the bulk ions and electrons accelerate injected ions above v_b , so that when $T_e \sim T_i$ an energetic tail is formed with slope approximately equal to that of the bulk-plasma distribution. The angle-averaged distribution becomes

$$F_h(v) = c_3 (v^3 + v_c^3)^{-1}, \quad v_m \leq v \leq v_b \quad (2.8a)$$

$$F_h(v) = c_3 (v_b^3 + v_c^3)^{-1} \exp[-m_h (v^2 - v_b^2)/2T_e], \quad v > v_b \quad (2.8b)$$

where c_3 is the normalization constant. Neglect of energy diffusion at lower velocities is not serious when $v \geq 2v_i$ [24].

By integrating Eq. (2.7), one obtains a simple expression for the fast-ion slowing-down time [20,21], which is the time to completely thermalize with the bulk ions:

$$\tau_s = \frac{\tau_{se}}{3} \ln \left[1 + \frac{v^3}{v_c^3} \right] \quad \text{where} \quad \tau_{se} = \frac{6.3 \times 10^8 A_h T_e (\text{eV})^{3/2}}{Z_h^2 n_e \ln \Lambda_e} \quad (2.9)$$

The fast-ion density is then $n_h = S \tau_s$, where S is the injection rate per unit volume.

Energetic ions injected at a large number of angles with respect to the magnetic axis tend to form a velocity distribution of the type given by Eq. (2.8). Evidently, charged fusion-reaction products also fall into this category. In almost all practical situations $\Gamma \ll 1$ and these very energetic ions are far too few to interact with each other. If these ions can be confined, their velocity distribution should be described reasonably well by Eq. (2.8). For these ions $v_b \gg v_c$ when $T_e \approx 30$ keV, so that most of their energy goes directly to plasma electrons [26].

2.2.1 Energy Clamping. It is possible in principle to "clamp" the fast-ion energy for times of order of the energy diffusion time [13], and to clamp the fast-ion velocity (i.e., directional motion) for times of order of the angular scattering time. Energy clamping requires a continuous power input to the fast ions to compensate for the bulk-plasma Coulomb drag [Eq. (2.7)]. The two most readily demonstrated means of energy clamping in a tokamak use the toroidal electric field (for co-injected ions) [27] or adiabatic major-radius compression [13,28]. For ions injected parallel to the magnetic field with $v > v_c$, Eq. (2.3) shows that the Coulomb drag is counterbalanced when $eE^* \sim m_h v / \tau_{se}$. If compression clamping of parallel-injected ions is used, then conservation of angular momentum prescribes a compression rate $C^{-1} dC/dt = v^{-1} dv/dt$, where dv/dt is given by Eq. (2.7). Other proposed clamping methods include a rigid displacement of the plasma column at the transit frequency of hot ions around the torus ("wobble clamping") [29], and the application of resonant electromagnetic waves [30].

The energy clamping technique can potentially increase Q in pulsed operation (Sect. 6.3), and also augment the effectiveness of neutral-beam heating (Sect. 5.1). Reference [25] shows how to modify the Fokker-Planck equation in order to investigate the evolution of $f_h(v)$ during compression clamping.

A typical compression period is of order $\tau_s/2$ [30], during which time appreciable angular scattering and energy diffusion of the fast ions are found to occur.

Any technique that can be used for energy clamping can probably be adopted for the production of energetic ions from the less collisional members of the thermal-ion population. Application of radiofrequency waves at twice the ion cyclotron frequency of a minority ion species seems well suited for this purpose [31].

2.3 VELOCITY DISTRIBUTIONS FOR $\Gamma \sim 1$

When $n_h/n_e \geq 0.1$, both the high-energy tail of $f_i(v)$ and the tail of $f_h(v)$ above the injection velocity v_b can be significantly distorted. Calculation of these effects has so far been rather limited. Perturbation analyses have been used to estimate the distortion of $f_h(v)$ due to collisions among the hot ions [16,32]. Distortion of $f_i(v)$ is due to momentum transfer from the fast ions at a rate larger than the collisional relaxation rate of accelerated ions with the warm bulk ions. Figure 5 taken from Ref. [33] shows that this distortion is significant for $(n_h/n_i)(v_i/v_c) \geq 0.02$, or $n_h/n_i \geq 0.1$ if $T_i \sim T_e$. The energetic alpha-particle population of an ignited D-T reactor turns out to be too small to produce significant distortion when $T_i \gtrsim 30$ keV [33]. The calculation of Ref. [33] neglects interaction among the accelerated warm ions themselves, an effect that would lead to a more Maxwellian-like bulk-ion population.

It appears that a numerical solution of the full Fokker-Planck equation will be necessary to evaluate accurately these distortion effects — which could result in increased fusion reactivities, especially when W_b is relatively low. The distortion effect caused by an appreciable energetic-ion population

often makes it difficult in experimental work using charge-exchange diagnostics to distinguish the decelerated hot-ion population from accelerated warm ions, if both species are hydrogenic. Consequently, it may be difficult to determine the plasma bulk-ion "temperature." With increasingly intense injection, however, it becomes no longer meaningful to refer to ion temperature.

2.4 VELOCITY DISTRIBUTIONS FOR $n_i \gg 1$

The continuous evolution of $f_i(v)$ and $f_h(v)$ with increasing n_h/n_e has yet to be worked out for the entire range $0 \leq n_h/n_e \leq 1.0$. However, numerical Fokker-Planck solutions have been obtained for $n_h/n_e \sim 1.0$ [10,14,34]. The most interesting case, at least for near-term application, appears to be CIT operation [Fig. 1(c)] which utilizes both co- and counter injection and requires fairly rapid removal from the hot reacting region of ions decelerated below W_1 . (See Sect. 6.2 for a discussion of the setting-up and maintenance of CIT plasmas.) The numerical results indicate that when the warm-ion lifetime is less than τ_{Ee} , the hot ions form displaced Maxwellian-like distributions, as shown in Fig. 6. These distributions can be understood qualitatively from the following model. The strongest interaction between oppositely injected ions is pitch-angle scattering. This long-range interaction produces differences in the velocities among ions injected together, so that they can now undergo self-diffusion. Because of their small relative velocities, these ions form a "thermal" distribution among themselves, while slowing down on the plasma electrons. (The perpendicular "temperature" of the ions tends to be somewhat larger than the parallel "temperature," because pitch-angle scattering is the principal means of interaction between the two oppositely traveling distributions.) When $\tau_p < \tau_E$ and $T_e \ll W_b$, the CIT distribution functions can be calculated with useful accuracy by means of a very simple model, in which it is assumed that the average ion drift energy equals the total plasma thermal energy [35].

The relative time scales for "equilibration" among the ions of each distribution (τ_{hh}) and the slowing down on electrons (τ_{se}) is [20]

$$\frac{\tau_{hh}}{\tau_{se}} \approx 2 \left(\frac{T_h}{T_e} \right)^{3/2} \left(\frac{m_e}{m_h} \right)^{1/2} \frac{1}{Z_h} \frac{\ln \Lambda_e}{\ln \Lambda_h} \quad (2.10)$$

Provided that the beam "temperature" $T_h \lesssim 4T_e$, for deuterons we have $\tau_{hh} \ll \tau_{se}$. The Fokker-Planck solutions show that this condition on T_h/T_e is generally satisfied [14,34]. However, if T_e is very large, then the counter-streaming distributions tend also to thermalize with each other in a time $< \tau_{se}$, thereby forming a nearly isotropic distribution at very high temperature, as shown in Fig. 6(C). The condition for ion thermalization before deceleration is, approximately, that the relative injection velocity ($2v_b$) be less than the "critical velocity." For oppositely injected D and T beams, this condition is $W_b \gtrsim 4T_e$, where W_b is the deuteron injection energy.

If the tokamak machine parameters such as magnetic field permit very dense plasmas, moderate-energy neutral beams injected tangentially may have inadequate penetration (see Sect. 4.1). In this case it may be most profitable to set up a nearly isotropic energetic-ion plasma by direct beam injection at a number of angles. If $W_b \gtrsim 4T_e$, a Maxwellian-like ion population at very high temperature and fueled entirely by injected ions might be formed and sustained in this way.

3. VELOCITY-SPACE STABILITY

When $\Gamma \approx 1$ and $f_h(v)$ is isotropic, in steady state the total ion velocity distribution is monotonically decreasing with velocity [cf. Eq. (2.8) for $\Gamma \ll 1$], and therefore absolutely stable to velocity-space modes. More practical distributions, such as those shown in Fig. 3 and 4, are anisotropic for $v \geq v_c$ with regions where $\partial f_h / \partial v_{\parallel}^2 > 0$, and possibly $\partial f_h / \partial v_{\perp}^2 > 0$, so that one might expect that a variety of unstable modes could arise. It is a remarkable feature of beam-driven tokamak plasmas that for steady-state conditions of practical interest, the ion velocity distributions are stable. This favorable result is due in part to the breadth of the velocity distributions, and in part because the injection velocities are usually smaller than the Alfvén velocity. It is only in the case of a transient peaked distribution such as occurs just after the start of injection, or when the fast-ion energy is "clamped" (Sect. 2.2.1), that the possibility of velocity-space instability is serious.

The lifetime and slowing-down history of energetic ions may be affected by instabilities caused by the energetic-ion population itself, or by instabilities naturally present in the tokamak plasma (such as drift waves), which in turn may be influenced by the presence of fast ions. The sources of free energy introduced by the fast-ion population are found in the nonMaxwellian ion distribution ($\partial f_h / \partial v_{\parallel}^2 > 0$, or $\partial f_h / \partial v_{\perp}^2 > 0$), an unfavorable density gradient of the fast ions ($\partial f_h / \partial r > 0$), or possibly the magnetic field of the axial currents induced by streaming ions. In this section, we summarize briefly the results of investigations that have dealt primarily with instabilities in velocity space that arise from the nonMaxwellian ion distributions (i.e., infinite-medium modes), and that could lead to enhanced slowing-down or pitch-angle scattering.

3.1 PARALLEL INJECTION, $\Gamma \approx 1$

The infinite-medium instabilities of TCT-type plasmas have been treated in detail by the authors of Refs. [36,37], using hot-ion distribution functions that represent approximate steady-state solutions to Eq. (2.3) in the region $v > v_c$, where $f_h(v)$ is expected to be anisotropic. Their approach is to isolate the source of the instability and look for a resonance of the fast ions with possible bulk-plasma modes. They find that the longitudinal ion sound mode ($k_{\parallel} \gg k_{\perp}$) is always stable, even if the bulk ions are assumed to have zero temperature. The slow Alfvén wave is stable, provided that the injection velocity v_b is less than the Alfvén velocity v_A . The fast Alfvén (whistler) mode is stable if $v_b < 2.6v_A$. Since $v_b/v_A \approx (\beta_w W_b/2T)^{1/2}$, where β_w is the beta of the bulk plasma, the first condition is violated only if $\beta_w \geq 0.1$ — an unlikely prospect in tokamaks.

If $Z_{\text{eff}} \gg 1$, so that pitch-angle scattering is enhanced and $\partial f_h / \partial v_{\perp}^2 > 0$, it appears that modes with $k_{\perp} \gg k_{\parallel}$ resonant at harmonics of the beam-ion cyclotron frequency ω_{ch} could be excited transiently [36]. The chief effect of such an instability would be to increase velocity diffusion near the injection energy. However, it appears that the steady-state $f_h(v)$ is stable. Reference [38] considers isotropic distributions of the form $W e^{-\beta W}$ which might correspond to the transient situation near the start of multi-angle injection, or to clamped distributions before energy diffusion has become significant. For very peaked distributions (small s and β), they find that lower-hybrid modes can be destabilized.

It appears that the ordinary electromagnetic ion cyclotron wave could possibly be driven unstable during intense beam injection [39], but its appearance seems unlikely, because of the variation in cyclotron frequency along the guiding-center trajectories of the fast ions.

The interaction of the anisotropic velocity distribution with the density-gradient energy reservoir has also been investigated [36,37]. Drift waves and trapped-electron modes cannot be induced by the fast ions, provided that $(\nabla n_h / \nabla n_e) > 0$. Drift modes, however, may be destabilized if the neutral-beam penetration is sufficiently poor so that $dn_h/dr > 0$. No new trapped-ion modes are induced by the beam [40]. It appears that the only finite-medium mode that could be excited by the energetic ions is the ion branch of the drift-Alfven mode, which can be made unstable if $n_h(r)$ is too peaked near the magnetic axis [41]. This instability is due to a resonance $k_{\parallel} v_A = k_{\perp} v_d$, where v_d is the VB-drift velocity of the fast ions. The instability is localized to radii smaller than the resonance position, and its effect is to broaden the fast-ion profile. (This instability can also be excited by charged fusion-reaction products for which $v_b > v_A$.)

While perfectly clamped energy distributions of the form $f_h(v_{\parallel}) \sim \delta(v_{\parallel} - v_b)$ can be expected to be unstable, the advantages of energy clamping can be largely retained for the relatively broad $f_h(v_{\parallel})$ that is likely to prevail in practice [30].

3.2 PERPENDICULAR INJECTION, $\Gamma \approx 1$

Modes with $k_{\perp} \gg k_{\parallel}$ and with frequencies near the ion cyclotron harmonics could be driven unstable by perpendicular injection, particularly in the transient period within one slowing-down time of the start of injection [36, 42]. The drift-cone mode is likely to be stabilized by the large bulk-plasma population, although the convective loss-cone mode could possibly be made unstable. But in steady-state operation, it is difficult to find conditions where instability persists, unless the lower-energy portion of the distribution is depleted by charge exchange or large orbit losses of decelerating

ions. The energy reservoir $\partial f_h / \partial v_\perp^2 > 0$ could perhaps be enhanced by energy clamping when employed to delay the slowing down of injected ions.

3.3 INTERACTION WITH NONRESONANT BACKGROUND MODES

Estimates show that trapped-particle modes and drift waves naturally present in the background plasma are barely influenced by the fast ions [36]. Furthermore, if the wave frequency $\omega_* \ll k_\parallel v_h$, where v_h is a typical fast-ion velocity, the radial diffusion rate of the fast ions induced by the low-frequency background modes is smaller than that of the bulk ions by a factor $\sim T_i/W_b$ [36]. This result is due to the short correlation time between the fast ions and the background waves.

The interaction of stable kink modes with the fast ions has not been investigated theoretically, but such modes may be expected to influence the orbits of ions with small $v_\parallel/v$. When $f_h(\vec{v})$ is strongly peaked, Alfvén waves can be excited near rational magnetic surfaces, even when $v_b < v_A$ [43]. Parallel injection leading to $p_\parallel > p_\perp$ has a stabilizing influence on the localized interchange modes, provided that $dp_\parallel/dr < 0$ [44]. But perpendicular injection might reduce the stability of the resistive interchange modes [44].

3.4 STABILITY OF CIT/CBT PLASMAS

The CIT system appears to have considerable free energy, since ideally the ion velocity distribution is highly anisotropic ($p_\parallel \gg p_\perp$), and $\partial f_h / \partial v_\parallel^2 > 0$ for $|v_\parallel| < u_b$, where u_b is the mean streaming velocity (cf. Fig. 6). Nevertheless, for all practical conditions the steady-state CIT velocity distributions are stable to electrostatic and electromagnetic infinite-medium modes [45].

This stability is due primarily to the large thermal spread of the ion distributions.

Fokker-Planck calculations have shown that the ion velocity distributions $f_h(v)$ can be approximated by shifted Maxwellians (cf. Fig. 6):

$$f_h(v) = \frac{n_h^{3/2}}{(2\pi)^{3/2} T_{h\perp} T_{h\parallel}} \exp[-m_h (v_{\parallel} + u_b)^2 / 2T_{h\parallel} - m_h v_{\perp}^2 / 2T_{h\perp}] \quad (3.1)$$

We define $\Delta W = 2u_b v_{ti} m_h$, where $v_{ti} = (2T_{h\parallel}/m_h)^{1/2}$. Then $\Delta W/W_b = 4(T_{h\parallel}/W_b)^{1/2}$, where $W_b = 1/2 m_h u_b^2$. The Fokker-Planck calculations for steady-state operation reveal the following relations: $T_{h\parallel} \geq 1.5 T_e$, $T_{h\perp} > T_{h\parallel}$, $v_{ti}/u_b > 1/3$, $\Delta W/W_b > 4/3$.

The stability of counterstreaming-ion distributions to electrostatic and electromagnetic velocity-space modes has been investigated in conjunction with both magnetospheric phenomena and laboratory experiments [39,46-48]. The stability criteria are summarized in Table 3. Since the streaming energy density can be treated formally as if it were a thermal pressure [49], the usual constraint on β_p for steady-state tokamak equilibrium (i.e., $\beta_p \geq R_0/a$) allows only a small value of u_b/v_A . This situation contrasts markedly with TCT operation, where most of the plasma mass resides in the background ions, so that a fraction of these ions can stream at $v > v_A$ without violating the tokamak equilibrium limit. Table 3 also indicates the range of parameters of "ideal" CBT/CIT distributions, which are bound to be the least stable. That is, a substantial population of colder ions enhances the slowing-down rate of energetic ions, resulting in a smaller u_b , and also increases ΔW and $T_{h\parallel}$. (At the start of injection, the ion distributions are more peaked, with a smaller ΔW and a larger u_b , so that instabilities can

be excited. This transient phase, however, lasts only one or two slowing-down periods, typically less than 0.5 s.)

Turning briefly to MHD modes that could result from the anisotropic pressure, it is clear that the "firehose" instability cannot arise in tokamak operation, where $p_{\parallel}/B^2 \ll 1$. As mentioned previously, the anisotropy $p_{\parallel} > p_{\perp}$ has a stabilizing influence on localized interchange modes [44].

The results surveyed in Sections 2 and 3 show that by control of the intensity of beam injection and the confinement time of plasma particles and energy, it is possible to vary the steady-state ion distribution of a tokamak plasma from near-Maxwellian with a small energetic component, to counter-streaming displaced Maxwellians. Although velocity-space instabilities may arise during perpendicular injection, or during parallel injection with v_b much greater than the Alfvén speed, there appear to be no unstable modes that would significantly affect the prospects of beam-driven tokamak reactors.

4. INJECTION AND ORBITS OF ENERGETIC IONS

4.1 NEUTRAL-BEAM TRAPPING

Although the direct introduction of energetic ions from an external source into a tokamak plasma is not inconceivable, penetration across the magnetic field to the edge of the plasma is accomplished most straightforwardly by means of a beam of neutral atoms. Significant penetration into the plasma is also accomplished most readily by use of neutral beams. Liouville's Theorem [50] when applied to magnetic confinement of injected energetic particles places no restriction on the containment of neutral-beam-injected ions, because of the change in charge-to-mass ratio that occurs upon ionization. Liouville's Theorem, however, apparently forbids the containment of particles injected directly as monatomic ions, although it turns out that this restraint can be circumvented somewhat (see Sect. 4.2).

All experimental tokamak devices and conceptual reactor designs that specify energetic ions for heating depend on neutral beams to deliver these ions to the interior of the plasma. Given acceptable injector efficiency, the most important restrictions on the choice of neutral-beam energy W_b are that the neutrals must penetrate adequately before ionization, and that essentially all the energetic particles must be trapped in the plasma. Energetic neutrals are trapped by charge-exchange and impact ionization on plasma ions, and by electron ionization. The attenuation of the neutral beam flux I_b is given by

$$I_b(x) = I_{b0} \exp\left[-\frac{1}{v_b} \int K(x) n_e(x) dx\right] \quad (4.1)$$

where x is the distance along the beam path, and

$$K = \sum_i \frac{n_i}{n_e} \langle \sigma_{cx} v_{bi} \rangle + \sum_i \frac{n_i}{n_e} \langle \sigma_I v_{bi} \rangle + \langle \sigma_e v_e \rangle \quad (4.2)$$

Here, $\sigma_{cx}(v_{bi})$ and $\sigma_I(v_{bi})$ are the charge-exchange and impact-ionization cross sections, respectively, v_{bi} is the relative velocity between the fast neutrals and a bulk-plasma ion, and $\langle \sigma_e v_e \rangle$ is the electron ionization rate, averaged over a Maxwellian electron velocity distribution. The reactivities $\langle \sigma_{cx} v_{bi} \rangle$ and $\langle \sigma_I v_{bi} \rangle$ are averaged over the plasma ion distribution. For thermonuclear and TCT-type plasmas, $v_{bi} \approx v_b$. Figure 7(a) shows the energy dependence of these reactivities [51]. When $n_e(x) \approx \text{constant}$, we can define a trapping length λ_t for exponential attenuation, namely, $\lambda_t = v_b / (Kn_e)$.

The ratio λ_t/a_p determines the beam trapping profile in a manner worked out in detail by the authors of Ref. [52]. Figure 8 shows schematically the process of tangential beam injection. While the electrons remain on the flux surface on which they are born, the fast ions follow guiding-center orbits whose poloidal motion is approximately circular, but centered about a "stagnation axis" which is displaced from the magnetic axis because of the curvature drift in the toroidal field. (See Ref. [53] for an excellent review of particle orbits in a tokamak.) The stagnation axis is located typically several cm outside (inside) the magnetic axis for ions injected parallel (antiparallel) to the plasma current I_p . As shown in Fig. 8, for injection just inside the magnetic axis, the orbits of co-injected ions (injected parallel to I_p) are shifted outward in major radius, while the orbits of counter-injected ions (injected antiparallel) are shifted inward. Figure 9 shows illustrative "spatial shape factors" $H(r)$ for co- and counter-injection. Here $H(r)$ is the density averaged over the magnetic surface of newly ionized fast ions. Even when λ_t/a_p is significantly less than unity, $H(r)$ can be peaked on the magnetic axis, because the volume of the magnetic shells decreases

with decreasing r . Thus for a parabolic $n_e(r)$ and co-injection, $H(r)$ is peaked at $r = 0$ for $\lambda_t/a_p \geq 1/4$, where λ_t is calculated at $n_e(0) = n_{ec}$. For counter-injection, axial peaking is not as marked and some ions are found on unconfined orbits, because of the unfavorable shift in major radius of the ion drift surfaces [21,52]. The path length for tangential injection is approximately $2a_p (2R_o/a_p)^{1/2} \sim 5a_p$, so that for $\lambda_t/a_p \geq 1$, at least 98% of the co-injected neutrals and 85% of the counter-injected neutrals are ionized and trapped in the plasma.

Axial peaking of the beam trapping profile is desirable for a number of reasons. The efficiency of plasma heating is increased when the beam deposition is peaked on axis, since the diffusion length from the source of heat to the cold plasma boundary is then maximized. In beam-driven reactors, peaking in a hot central region will maximize the fusion gain, while the remainder of the plasma is heated principally by conduction from the hot core (as well as by alpha deposition), and serves as an insulator for the primary fusion-production region. Fast ions deposited in the outer region of the plasma are more susceptible to loss by charge exchange and by neoclassical diffusion; the banana excursions $\Delta r/a_p$ are often not negligible. Finally, alpha particles produced in the central region are less likely to be lost on the limiter. While excessive peaking of the beam deposition could have an adverse effect on MHD equilibrium, and possibly give rise to a drift-Alfven instability, tokamak plasmas with sufficient size and density for $n_e \tau_E \geq 10^{13} \text{ cm}^{-3} \text{ s}$ are likely to be uncomfortably opaque to neutral-beam penetration, a subject to which we now turn.

4.1.1 Effect of Impurities on Beam Penetration. The trapping processes [Eq. (4.2)] depend only on the fast-neutral velocity, so that D^0 with energy W_b is equivalent to H^0 with energy $1/2 W_b$, or to T^0 with energy $3/2 W_b$. If the

energy per nucleon $W_{bl} \geq 50$ keV, the impact-ionization cross section on impurity ions of charge Z_I is expected to be larger than that of Fig. 7(a) by the factor Z_I^2 [54]. If $v_{bi} \approx v_b$, the second term on the right-hand side of Eq. (4.2) can be written $Z_{eff} \sigma_I v_b / n_e$, where σ_I is the value for protons. Thus at the higher W_b , where impact ionization is dominant, one has $\lambda_t \propto (n_e Z_{eff})^{-1}$. (The charge-exchange cross section of impurity ions is also expected to increase somewhat [55], but this effect has not yet been resolved). Figure 7(b) shows the product $n_e a_p$ as a function of injection energy for $\lambda_t = v_b / K n_e = a_p / 4$. As noted previously, for parabolic and similar profiles with peak density n_{ec} an acceptable beam deposition profile is obtained with $\lambda_t \geq a_p / 4$, where λ_t is calculated at n_{ec} . Figure 10 shows that for $W_{bl} \geq 75$ keV, one has [56]

$$W_{bl} \text{ (keV/nucleon)} \approx 4.5 \times 10^{-15} n_{ec} a_p Z_{eff} \quad (4.3)$$

With increasing plasma opacity $\langle n_e Z_{eff} \rangle a_p$, the required W_b for adequate penetration becomes uncomfortably large for neutral beams formed from positive ions (see Sect. 7). If the fast neutrals are trapped in the outer region of the plasma and their energy deposited there, sputtering of the first wall by hot-particle bombardment is aggravated. The resulting enhancement of impurity influx further reduces beam penetration, so that the wall sputtering and plasma Z_{eff} continue to increase. This phenomenon has been called the "impurity-trapping instability," and may prevent heating of the main body of the plasma. Several computer simulations of this effect have been described recently [57].

The problem of penetration into plasmas of high opacity can be alleviated by using increasingly oblique injection; near-perpendicular injection would allow nearly a factor of 2 reduction in W_b for a given $H(r)$ [58],

and offers the additional advantage that access to the torus is more convenient. Potential difficulties with perpendicular injection are that (1) micro-instabilities are excited more easily, particularly at the start of injection; (2) ions with large transverse energies born near the plasma edge often have banana orbits that intercept the limiter [17]; (3) ions with large $W_{\perp}$ may be trapped in "superbanana" orbits that escape the plasma (see Sect. 4.3.2).

The injection of helium atoms has been proposed as a method of achieving more favorable penetration at $W_b \sim 100\text{--}200$ keV [59]. For $W_b \geq 100$ keV, the efficiency of production of He^0 beams is much larger than that for D^0 beams formed from positive ions, while the penetration distances λ_t are comparable. Two disadvantages of He^0 injection are reduced effectiveness of cryopumping in the beam ducts, and dilution of the reacting-ion population of the plasma ($\text{D-}^3\text{He}$ reactors excluded!).

4.2 RIPPLE-TRAPPING INJECTION

A new injection technique that takes advantage of the vertical VB-drift of ripple-trapped ions was proposed recently to increase the effective penetration length of energetic neutral beams [60]. As shown schematically in Fig. 10(a), the local magnetic wells of a nonaxisymmetric toroidal field will persist in the presence of a plasma current when the peak-to-peak vacuum ripple δ exceeds $(2r/R)(\sin\theta/Nq)$, where N is the number of toroidal-field coils [61]. If δ_* is the actual well depth, energetic ions with $v_{\parallel}/v < \delta_*^{1/2}$ are trapped in these wells and drift vertically with the $\vec{B} \times \nabla B$ velocity (superbanana orbits). Consider the configuration of Fig. 10(b), where energetic ions formed from neutral beams injected vertically from the bottom are mirror-trapped in the high-ripple region at $z < 0$ and drift upward. The ripple is made to decrease rapidly at $z > 0$, in order to form a "symmetric trapping"

region [see Fig. 10(a)], where ions with small $v_{\parallel}/v$ must follow banana orbits, and have no time-averaged vertical drifts. In order to reach the midplane without first scattering out of its super-banana orbit, an ion must have energy (cgs units)

$$W_{\min} \text{ (keV)} = 2.8 \times 10^{-9} \left(\frac{\langle n Z_{\text{eff}} \rangle a R B_t}{\delta_0} \right)^{2/5} \quad (4.4)$$

where δ_0 is the ripple at $z = 0$. Figure 11(a) shows $W_{\min}$ as a function of δ_0 . For plasmas of large opacity, the required $W_{\min}$ with moderate ripple is an order of magnitude smaller than that required by the conventional injection method. Because of the diffusive nature of the pitch-angle scattering process, the symmetric trapping region at $z > 0$ is essential to trap all injected ions with $W \geq W_{\min}$, and to prevent further upward drift of these ions until they thermalize.

The feasibility of the ripple-assisted penetration technique is supported by orbit calculations with a Monte Carlo guiding-center code that shows that ions with energy in the range given by Eq. (4.4) are captured in the central plasma region, and are confined until they thermalize [60]. An example is shown in Fig. 11(b). The vertically asymmetric ripple also gives rise to nonsymmetric banana orbits that gradually drift outward [62]. In the absence of Coulomb drag, all ions injected into the central plasma region by ∇B drift would escape fairly rapidly to the outer region by this banana-diffusion process, as expected by Liouville's Theorem [50]. Coulomb drag does, however, allow injected ions in a limited energy range near $W_{\min}$ to penetrate to the central region, and also to thermalize there. The upper end of this energy range is increased greatly by installing toroidal-field ripple in only those locations that are needed for injection.

The nearly monokinetic hot-ion distribution function in the ∇B -drifting region introduces a strong source of free energy that could result in loss-cone-type instabilities [60]. Drift-cone modes are probably stabilized by the dense warm-ion background. Convective loss-cone modes can be stabilized by radial energy convection and magnetic shear, provided that $\Delta W_b/W_b$ can be made $\geq 10\%$ in this region [63].

The ∇B -drift injection technique offers in principle a means of injecting ions directly into a tokamak [60,64] — that is, without relying on neutralization at all for penetration. However, no convincing way is known for transporting 50 to 500-keV ions from a source at low magnetic field, into the high-field tokamak region. Furthermore, the stability of a dense ion beam drifting perpendicular to the magnetic field in a region outside the plasma, where the damping mechanisms mentioned above are absent, is dubious [63]. Finally, the charge-exchange loss of ions drifting slowly through a significant distance in a region near the plasma edge can be considerable.

4.3 FAST-ION LOSS PROCESSES

There are several mechanisms by which energetic ions can be lost from the torus before they thermalize. These loss processes generally reduce the performance of beam-driven reactor plasmas.

4.3.1 Charge-Exchange Loss. A tokamak discharge may contain a significant population of neutral atoms produced by the inward diffusion of neutrals from the wall region, or by the charge-exchange trapping of injected energetic neutrals. The first mechanism gives rise to a significant neutral density ($n_0 \sim 10^8 - 10^9 \text{ cm}^{-3}$) in most present-day tokamak plasmas, but tokamak reactor plasmas will be sufficiently opaque so that in the hot central region, one expects $n_0 \gtrsim 10^7 \text{ cm}^{-3}$ [57]. The Q -values of TCT-type reactors are reduced by less than 10% for $n_0/n_c \sim 10^{-6}$ [65], so that for typical plasmas densities, $n_0 \gtrsim 3 \times 10^7 \text{ cm}^{-3}$ is acceptable. Charge-exchange loss is then serious only if

poor neutral beam penetration results in large fast-ion densities in the outer plasma region, where n_0 is large. Charge-exchange loss on impurity ions may also be significant [66], but this problem remains relatively uninvestigated.

For most types of beam-heated reactors with $W_b \geq 100$ keV (D^0), relatively few charge-exchange ions are introduced during the process of neutral-beam trapping (cf. Fig. 7). For CIT-type reactors, however, the relative velocity of injected neutrals with the ion population streaming in the same direction is very small, so that approximately 50% of the injected neutrals are trapped by charge exchange. Taking into account that neutrals can be lost from the plasma only by several generations of charge-exchange, it is found that up to 10% of the injected energy may be lost as a direct result of neutrals introduced by the trapping process [35].

4.3.2 Velocity-Space Loss Regions. Particles following banana orbits that extend outside the discharge region will be scraped off by a limiter or by divergent field lines. Fast ions acquire such orbits either by perpendicular neutral-beam injection, or by pitch-angle scattering from passing orbits. The authors of Ref. [17] have calculated the drift orbits of particles with all possible pitch angles passing through a given radius on the outside of the torus and in the equatorial plane. Typical loss regions for an axisymmetric torus are shown in Fig. 12. Evidently, counter-injected particles ($v_{||} < 0$) need be scattered through a smaller angle than co-injected particles, in order to find themselves in the loss cone. While this loss process can be important in present-day tokamaks ($I_p \sim 100$ -200 kA), it is unlikely to be serious for reactor plasmas ($W_b \geq 200$ keV, $I_p \geq 1$ MA), since the poloidal gyroradius $\propto W_b^{1/2}/I_p$. However, a substantial fraction of charged fusion-reaction products can be lost in this way unless $I_p \geq 3$ MA [67].

The loss of energetic ions is also enhanced when $Z_{\text{eff}} \gg 1$, since fast ions then undergo several cycles of trapping and detrapping in banana orbits, and diffuse across the plasma to the edge. For this neoclassical diffusion process to proceed faster than thermalization, at $I_p \geq 1$ MA one requires typically $Z_{\text{eff}} > 5$.

If the toroidal field is sufficiently nonaxisymmetric to produce local magnetic wells of depth δ_* , then energetic ions with $v_{\perp}/v < \delta_*^{1/2}$ are trapped in these wells and drift vertically. The step size before detrapping is $\Delta_z = \delta_* \tau_{90} v_{\text{drift}} \propto W^{5/2}/Z_{\text{eff}}$. If $\Delta_z \geq a_p$, ripple-trapped ions drift out of the plasma. Loss due to ripple trapping increases at first with Z_{eff} , because more ions are scattered into the ripple-trapping band, but the loss decreases when Z_{eff} becomes sufficiently large so that $\Delta_z < a_p$ [68]. This loss process, which is the inverse of the ripple-trapping injection scheme, is most serious for energetic ions in the outer region of the plasma, where δ_* tends to be higher [61].

(As noted in Sect. 4.1.1, the neutral-beam penetration length depends on Z_{eff} . It has been proposed to determine the local value of this quantity by introducing toroidal-field ripple, and measuring the fluxes of energetic ions escaping transverse to the path of fast neutrals injected perpendicularly in the horizontal midplane [69].)

4.3.3 Fusion Burn-Up. Typically, only a tiny fraction of energetic ions undergo fusion reactions while slowing down. For example, in a $Q = 1.0$ TCT plasma driven by 200-keV deuteron beams, only about 1% of the deuterons are burned up during deceleration. However, the integrated fusion probability during slowing-down increases with both W_b and T_e . For example, about 5% of the 1.0-MeV tritons produced from the D-D reaction react while thermalizing, if $T_e \sim 10$ keV.

Summarizing the importance of these various loss processes in reactor-size plasmas ($I_p \geq 1$ MA), we find that charge-exchange loss is unimportant except in CIT-type reactors. Banana orbits are likely to intercept loss cones only in the case of the very energetic charged fusion-reaction products, or in the case of perpendicular injection with poor neutral-beam penetration. Loss of energetic ions because of toroidal-field nonaxisymmetry (superbanana orbits) can be important, but in principle can be eliminated by suitable design of the TF coils to give small ripple at the plasma.

4.4 EFFECTS ON PLASMA EQUILIBRIUM

4.4.1 Beam-Induced Currents. Although the fast-ion current I_h is larger than the total injected current I_b by the stacking factor $\langle v_{\parallel} \rangle_s / 2\pi R_0 \sim 10^3 - 10^4$, I_h is generally only a fraction of the ohmic-heating current. On injection, the electric current of the energetic ions parallel to the magnetic field is cancelled identically by induced electron currents. These return currents relax with a very long time constant, equal to the bulk-plasma skin time [70], and even in steady state the ion current tends to be partly shielded by electron counter-currents. To investigate the possible effect of unshielded injected-ion currents, the authors of Refs. [37,71] have carried out a spatially two-dimensional computer simulation, where a monoenergetic group of 2500 "superparticles" is injected instantaneously, and the electron counter-current is removed in some arbitrary time. The (collisionless) guiding-center motions of the super-particles are followed in time and used to compute contributions to the toroidal current. For the case of co-injected beams, the plasma surface expands somewhat, but the flux surfaces are undistorted, even when the beam-ion current is comparable with the ohmic-heating current. Computations with counter-injected beams alone indicate

serious deformation of the magnetic surfaces, with appreciable outward drift of the injected particles. This behavior may be thought of as due to repulsion between the beam-induced current and the ohmic plasma current. When both co- and counter-injected beams are introduced in equal numbers, the undesirable effects of the counter-injected beams alone are reduced markedly.

The above considerations account only for electromagnetic effects, and in particular they neglect the momentum exchange between the fast ions and the background plasma. In the steady state the beam current is at least partially cancelled by plasma electrons which are accelerated by momentum transfer from the hot ions. But when $Z_{\text{eff}} > 1$, a net current can be produced by the injected ions together with plasma electrons driven by these ions, even when $E_{\phi} = 0$ [72,73], a phenomenon that may have considerable practical utility (see Sect. 5.3). The induced toroidal current is lowered significantly by collisional isotropization and banana-orbit trapping of the fast ions [74], or if oblique injection must be employed to ease neutral-beam penetration or accessibility.

Momentum transfer from the fast ions to bulk ions could also give rise to a large toroidal plasma rotation and a corresponding radial electric field [73,75]. For unidirectional injection, the plasma toroidal flow velocity can be comparable with the sound speed, if the only damping mechanism is perpendicular viscosity. This potential problem can be minimized by utilizing oppositely directed beams — or very oblique injection. Toroidal flow can also be damped effectively by a toroidal-field ripple $\sim 1\%$ [73,76], because of the friction between ripple-trapped particles and passing particles being accelerated by the fast ions.

4.4.2 Pressure. To a good approximation, the injected energetic-ion component in a large tokamak produces the same effects as an equal increment in the pressure of the bulk plasma. If a system of parallel and antiparallel

injection is used, there results a pressure anisotropy, but the modification of the MHD equilibrium solution for the tokamak is very small, even for CIT-type plasmas [77,78].

The effect on gross plasma stability of the anisotropic pressure caused either by parallel or perpendicular beam injection is considered in Ref. [44]. As mentioned in Sect. 3.3, parallel injection giving rise to $p_{\parallel} > p_{\perp}$ (as in CIT plasmas) has a stabilizing influence on the localized interchange mode, so that q -values as low as 0.8 are possible. For perpendicular injection giving rise to $p_{\perp} > p_{\parallel}$ (as in the VB drift-injection scheme), there is a reduction in stability of the resistive interchange mode.

4.4.3 Induced Electric Fields. Since the fast ions are displaced from the magnetic surfaces where they are born, a radial electric field is established when $n_h(r) = 0$, or when a large number of fast-ion orbits intersect the limiter region [42,75]. While the resultant poloidal rotation is strongly damped by ion viscosity in the poloidally inhomogeneous magnetic field [79], the toroidal rotation is damped only by toroidal-field ripple — or by some nonclassical process. With perpendicular injection the radial field can be enhanced, because of greater deviation of the ion orbits from the magnetic surfaces of birth, and because an even larger fraction of the fast ions tends to be lost on the limiter. Furthermore the charge distribution on a flux surface is nonuniform, because the trapped-ion orbits are found mainly on the large major-radius side of the torus; this effect gives rise to a poloidal electric field that may result in radial plasma convection [75].

In the VB drift-injection technique, the inward drift of the energetic ions (with the electrons remaining behind) may set up a radial electric field opposite in direction to the usual ambipolar field. Calculations show that for injection currents sufficient for heating, the reversed radial field is

unlikely to be dangerously large [62], but the use of large injection currents for fueling could lead to serious changes in plasma equilibrium.

5. HEATING, FUELING, AND CURRENT MAINTENANCE BY BEAM-INJECTED IONS

Table 2 lists the several functions served by the injected energetic ions in various operating regimes. This section elaborates on the heating, fueling, and momentum-related aspects of the injected ions, while Sect. 6 describes their role in enhancing the fusion reactivity of the plasma.

5.1 PLASMA HEATING

5.1.1 Power Balance. When $v_b \ll v_e$, the injected ions decelerate by Coulomb drag on plasma electrons and bulk ions according to Eq. (2.7). The electron drag exceeds the bulk-ion drag when $v > v_c$, where v_c is the "critical" velocity defined in Sect. 2.2. Figure 13 shows the fractions of injected power and momentum transferred to bulk-plasma ions, as a function of W_b/W_c [19,21]. If the equilibration time τ_{ei} of electrons and bulk ions is comparable with or less than the smaller of the energy confinement times τ_{Ee} , τ_{Ei} , then $\bar{W}_i \approx 3/2 T_e$, regardless of W_b/W_c .

In most of this discussion we employ a global (i.e., zero dimensional) model for the plasma, so that the values of density, temperature, and fast-ion energies represent various volume-averaged quantities. Plasma electrons, and ions with energies $< W_1$ lose energy to the outside world by particle diffusion, heat transport, charge-exchange (ions) and radiation (electrons). These loss processes are often represented most conveniently by the single quantity, τ_{Ee}^{-1} or τ_{Ei}^{-1} . The power balance per unit volume is then

$$\frac{3}{2} \frac{n_e T_e}{\tau_{Ee}} + \frac{\sum n_j \bar{W}_j}{\tau_{Ei}} = P_b + P_\alpha + P_{OH} \quad (5.1)$$

where $P_{OH} \propto T_e^{-3/2}$ is the ohmic power dissipation, P_α is the power deposition of charged fusion-reaction products, and P_b is the injected-beam power density. For a D-T plasma, $P_\alpha = 0.2 F_\alpha Q P_b$, where $Q \equiv P_f/P_b$ is the fusion power multiplication factor, and F_α is the fraction of fusion-alpha energy deposited in the plasma. As before, we assume that injected ions are perfectly confined if $W > W_1$.

For $T_e \geq 3$ keV, plasma current $J_{OH} \approx 300$ A/cm², and $Z_{eff} \approx 3$, $P_{OH} \approx 0.1$ W/cm³. This power density is generally much smaller than $P_b + P_\alpha$, so that P_{OH} will be neglected in the following.

When $n_h/n_e \ll 1$, we can use the quantity Γ [Eq. (2.1)]. Since $P_b = S W_b = n_h W_b / \tau_s$, we have

$$\Gamma = \frac{\left(P_b + P_\alpha \frac{\tau_{E\alpha}}{\tau_s} \frac{\bar{W}_\alpha}{W_\alpha} \frac{W_b}{W_h} \right) \frac{\bar{W}_h}{W_b} \frac{\tau_s}{\tau_E}}{P_b + P_\alpha} \quad (5.2)$$

where $W_\alpha = 3.5$ MeV, and we have taken $\tau_E = \tau_{Ee} = \tau_{Ei}$. Generally, the factor multiplying P_α in the numerator of Eq. (5.2) is close to 1.0. Then

$\Gamma \approx (\bar{W}_h/W_b)(n_e \tau_s/n_e \tau_E)$. Since $n_e \tau_s$ depends only on T_e and W_b , the value of Γ — which determines the type of reactor operation (cf. Table 2) — is determined by the $n_e \tau_E$ of the plasma and the desired operating temperature.

Since $\tau_{ei} \sim \tau_s$, the condition that $\bar{W}_i \approx 3/2 T_e$ is that $\Gamma \approx \bar{W}_h/W_b \sim 1/2$.

It is evident from Eq. (5.2) that when $P_\alpha \gg P_b$ (i.e., $Q \gg 5$), $\Gamma \approx \Gamma_\alpha \approx (\bar{W}_\alpha/W_\alpha)(\tau_{s\alpha}/\tau_E)$. Figure 14 shows Γ_α for D-T plasmas at ignition (i.e., $P_b = 0$) [81]. Evidently, Γ_α is comparable with values for beam-heated plasmas, and in particular for plasmas beam-heated to ignition. This

fundamental equivalence of plasma heating by fusion-reaction-injected ions and by neutral-beam-injected ions is one of the major attractions of neutral-beam heating. The equilibrium condition at ignition can be written explicitly as

$$\frac{3n_e^2 T_e}{n_e \tau_E} = n_D n_T \langle \sigma v \rangle E_\alpha F_\alpha \quad (5.3)$$

where $\langle \sigma v \rangle$ is the thermonuclear reactivity at $T_i = T_e$ and $E_\alpha = 3.5$ MeV.

If the plasma contains a significant high-Z impurity content, the loss rate τ_{Ee}^{-1} may be determined principally by line radiation and recombination radiation. Nevertheless, a given T_e can always be maintained if P_b is sufficiently large and the stability of the high- Γ plasma is favorable. Radiation loss (i.e., reduced $n_e \tau_E$) precludes ignition ($P_b = 0$, $\Gamma = \Gamma_\alpha$) when $n_e \tau_E$ drops below the value given by Eq. (5.3).

When the energetic-ion population becomes relatively large, Eq. (5.1) is still valid, but in Eq. (5.2) Γ should be replaced by Π_i . If $n_i \ll n_e$, as in the ideal CIT (CBT) regime, then with both D and T injection and neglecting P_α , the electron temperature is given by

$$\frac{3}{2} \frac{T_e}{\tau_{Ee}} = \frac{C_1 \ln \Lambda_e}{T_e^{3/2}} \left[\frac{n_D}{2} (\bar{W}_{hD}^{-3/2} T_e) + \frac{n_T}{3} (\bar{W}_{hT}^{-3/2} T_e) \right] \quad (5.4)$$

where $C_1 = 1.06 \times 10^{-13}$ (T_e, W in keV). In this type of plasma operation, where fueling is carried out entirely by the injected beams, both n_e and T_e are determined by $n_e \tau_{Ee}$ and the ion source strength.

5.1.2 Applications with $v_b \geq v_e$. In low- T_e plasmas, the injected beam velocity v_b may be comparable with v_e , so that Eq. (2.7) is invalid. The exact

calculation of dv/dt for these conditions is given in Ref. [22]. Beam injection into low- T_e tokamak plasmas, where ohmic heating is still strong, has two important purposes: to aid in current start-up, and to enhance the plasma density.

In large tokamaks ($I_p \geq 1$ MA), the flux swing in the OH (ohmic heating) transformer required for current start-up is determined mainly by the plasma inductance. However, if a significant impurity content is present at start-up, the resistive volt-sec caused by enhanced radiation and ionization losses can be substantial [82]; in extreme cases all the available OH flux-swing can be consumed during start-up, so that the discharge is extinguished. It appears that this problem can be alleviated greatly by initiating neutral-beam injection as soon as I_p is sufficiently large to confine fast ions and to produce a T_e so that $P_b/P_{OH} \geq 1$. For example, with $W_b = 100$ keV and $P_{beam} \geq 10$ MW, injection could well begin at $I_p = 300$ kA, loop voltage = 30 V, corresponding to $T_e \approx 100$ eV for devices with $R_o \approx 2.5$ m.

The maximum density obtainable in an ohmic-heated plasma is apparently limited by the current density J_{OH} [83], but it is found experimentally that substantially higher densities are attainable if gas puffing is employed together with intense neutral-beam injection [84,85]. It is thought that this density enhancement is made possible by Coulomb power transfer from energetic ions to the cold electrons formed from newly ionized gas in the outer region of the plasma. The additional heat source available to cold electrons reduces the tendency of the discharge region to shrink during gas puffing.

5.1.3 Energy Clamping. The previous section indicated how the neutral-beam injectors can be used to alleviate the requirements on the OH power supply for plasma heating at low temperature. Conversely, external circuitry can be used to aid the injected neutral beams in plasma heating, if energy-clamping techniques are employed (cf. Sect. 2.2.1)). For example, a major-radius

compression on the time-scale τ_s not only heats the bulk plasma, but can nearly sustain the hot-ion energy, which eventually must be deposited into the bulk plasma. The potential advantage of this technique is that if the fast-ion velocity is directed preferably along B_t , these ions are energized as C^2 , rather than as $C^{4/3}$ for thermal (isotropic) particles. The additional energy comes from the equilibrium-field circuit.

The proposed "wobble" technique, in which energy is introduced into the hot ions by oscillation of the vertical field at the ion transit frequency around the torus [29], seems more suitable for quasi-stationary heating. Again, the equilibrium-field circuit provides the additional energy.

At low T_e , where the toroidal electric field E_ϕ is significant, the OH circuit can give significant energy to co-streaming ions, and this effect has been observed experimentally (see Sect. 8). For temperatures of reactor interest, however, E_ϕ should be insignificant. Finally, a very recent tokamak experiment has indicated that high-powered radiofrequency waves near the lower hybrid frequency couple well to the perpendicular velocity components of injected energetic ions [85a].

5.2 FUELING

Present tokamak plasmas are fueled by plasma recycling from the limiter and wall, and by gas puffing, but it is not clear that these processes will be effective in devices of practical reactor size. Suggested fueling methods for plasmas of high opacity include high-velocity solid pellets [86], liquid jets [87], and atomic clustres [88]. At the present time, there is no assurance that the injection velocity required for penetration can be attained by any of these methods.

Injected neutral-atom beams can at least partly fuel the tokamak plasma, the chief drawback being the substantial energy investment in each particle. This difficulty does not exist in CIT-type operation, where most of the plasma ions are required to be energetic. Here we are concerned with TCT-type

plasmas, as well as with thermonuclear plasmas that are operating near ignition conditions.

Consider a beam-driven plasma with $\Gamma \approx 1$ ($n_i \approx n_e$) in which a fraction γ_b of the plasma energy losses is made up by the injected energetic ions. Then the beam-fueling fraction (Sect. 2.1) is

$$f_F = \frac{n_h/\tau_s}{n_i/\tau_p} = \frac{3T_e}{W_b} \frac{\tau_p}{\tau_E} \gamma_b \quad (5.5)$$

For $\tau_p \sim 2$ to $3 \tau_E$ [12] and $\gamma_b \sim 1$, complete beam fueling is possible when $W_b \gtrsim 9 T_e$. Tokamaks of 1-MA size presently coming into operation (Sect. 9) can apparently be beam-fueled entirely by their prescribed 40-keV heating beams, when $T_e \approx T_i \gtrsim 5$ keV. But in order to fuel beam-driven reactor plasmas ($T_e = 5$ to 10 keV) of opacity $\gtrsim 2 \times 10^{16} \text{ cm}^{-2}$, it seems necessary to use VB-drift injection so that relatively low W_b in the range of 100 keV can be employed (see Sect. 4.2).

An application of this latter technique may be necessary also to obtain the high plasma densities now being proposed as most favorable for ignition devices [89,90]. As mentioned earlier the plasma density can greatly exceed the limiting value with ohmic heating alone, when gas puffing is employed during intense beam injection [84,85]. But even higher densities seem likely, if significant density build-up comes also from decelerated energetic ions formed from low-energy neutral beams of high intensity.

For reacting tokamak plasmas operating near ignition conditions, $\gamma_b \ll 1$, so that one requires $\tau_p/\tau_E \gtrsim 10$ for $f_F \sim 1$. This requirement is unlikely to be met, and in any event the energetics seems unfavorable for high-Q operation: For 100-keV injection of D and T, a 2% burn-up is required just to recover the energy per atom expended in the injector, assuming 70% beam-production efficiency and 40% blanket conversion efficiency.

5.3 CURRENT MAINTENANCE

The momentum transferred to the electrons by injected energetic ions results in a current which together with the current associated with the streaming ions can comprise a significant fraction of the total plasma current [72]. After the return currents have been dissipated (see Sect. 4.4.1), the total beam-induced current in the steady state is [73,77,91]

$$J_b = \frac{E_\phi}{\eta_{ohm}} + e \sum_h Z_h n_h \langle v_{h\parallel} \rangle \left[1 - \frac{Z_h}{Z_{eff}} + G \left(\frac{r}{R} \right)^{1/2} \frac{Z_h}{Z_{eff}} \right]. \quad (5.6)$$

The second term on the right-hand side is the current associated with the energetic ions, while the third term is due to the equilibrium drift velocity of the electrons being accelerated by the energetic ions and decelerated by Coulomb drag on warm bulk ions. The fourth term, where the constant $G \sim 2$ [73], takes into account the electron population trapped in banana orbits. When Z_h and Z_{eff} are unity and the trapped electron population is negligible in the region of large fast-ion density, the total plasma current is zero unless an electric field is applied by the OH transformer. Otherwise, the current of the streaming energetic ions is essentially cancelled by the electron "back current" that is driven by momentum transfer from the fast ions. (Actually, an extremely small current density is still possible when $Z_{eff} = 1$ [92]). It seems likely that Z_{eff} will always be somewhat larger than unity, but if necessary a small impurity-ion population could be introduced by injection of helium beams, for example [93]. The population of trapped electrons could also be increased by introducing a significant toroidal-field ripple. Either of these procedures is likely to be palatable only for low- $n_e \tau_E$ reactors.

As indicated in Fig. 13, the injection energy W_b must be substantially greater than W_c in order that a large fraction of the beam momentum be

transferred to plasma electrons. The analysis that gives Eq. (5.6) neglects the drift velocity of the warm background ions, but in fact these ions are accelerated by Coulomb friction with the drifting electrons and can seriously reduce the net current. Consequently, a lower-energy beam ($W_b \lesssim W_c$) that transfers momentum mainly to bulk ions must be injected in the opposite direction, in order to oppose the electron frictional acceleration of the bulk ions [72]. However, if the plasma opacity is such that the high-energy beam has a favorable deposition profile, it seems unlikely that the low-energy beam can penetrate. This dilemma may perhaps be avoided by spontaneous damping of the bulk-ion momentum in a rippled field [73,76], or by the anomalous damping process that has been observed in some experiments (see Sect. 8).

The problem of bulk-ion momentum is essentially eliminated in CIT operation, where all plasma ions are injected by the beams, whose relative momenta can be adjusted to produce a large net current. For oppositely directed D and T beams with $W_b \sim 60$ keV, $W_T \sim 90$ keV, $n_D \sim n_T \sim 3 \times 10^{13} \text{ cm}^{-3}$, and $Z_{\text{eff}} = 1.5$, J_b can be 50 A/cm^2 [34]. If the required plasma current $\gtrsim 80 \text{ A/cm}^2$, it seems possible to arrange the relative neutral-beam intensities so that use of the OH transformer is unnecessary after the current start-up stage. Because plasma fueling is performed entirely by the beams, steady-state operation is possible in principle.

For TCT-type plasmas in particular, the beam-induced current is reduced by trapping of fast ions in banana orbits. This effect has been included by solving the bounce-averaged Fokker-Planck equation [74] which includes the gradient in the toroidal magnetic field. It is found that on a given magnetic surface, J_b is reduced by a factor $\sim [1 - (r/R)^{1/2}]$. It remains to carry out an accurate evaluation of the beam-induced current for a practical device, including spatial effects such as the radial variation of electron trapping and the beam deposition profiles. A related engineering analysis must include

the interaction of this current with the ohmic-heating regulation system [94]. Even in cases where the OH power supply is responsible for essentially all the plasma current, the fast ions can be employed to control the radial profile of current density by tailoring the beam deposition profile to produce the desired $T_e(r)$.

Another possible application of the momentum transferred from fast ions is to induce a net outward diffusion of impurity ions. Collisions of banana-trapped impurity ions with energetic ions injected antiparallel to the plasma current result in an outward drift of the impurity banana orbits [95]. It appears that extremely large beam currents are required for this purpose [96], so that the scheme would be acceptable only for low-Q reactor operation.

6. PERFORMANCE OF BEAM-INJECTED TOKAMAK REACTORS

6.1 FUSION REACTIVITY AND $n_e \tau_E$

The $n_e \tau_E$ requirement for a "conventional" thermonuclear plasma is prescribed by the Q -value desired, ranging from a "Lawson-type" $n_e \tau_E$ -criterion for $Q = 1$ [97], to the larger requirement for ignition ($Q \rightarrow \infty$). (Here Q is defined as fusion power/injected heating power.) For plasmas heated by reacting ion beams, on the other hand, there can be a large range of $n_e \tau_E$ -values to obtain a given Q at a given bulk-plasma temperature, because of the new flexibility in determining the total reactivity $\langle \sigma v \rangle$ of the plasma. Defining $U_t = (3/2)n_e T_e + n_i \bar{W}_i$ as the bulk-plasma energy density, the injected power density required to sustain U_t against losses is $P_b = U_t / \tau_E$. The fusion power density is $P_f = P_b Q = n_D n_T \langle \sigma v \rangle E_f$ where $E_f = 17.6$ MeV, so that

$$n_e \tau_E = \frac{Q U_t n_e}{n_D n_T \langle \sigma v \rangle E_f} \quad (6.1)$$

In a majority of operating modes, $n_D n_T \approx 1/4 n_c^2$. It follows from Eq. (6.1) that the larger $\langle \sigma v \rangle$, the smaller $n_e \tau_E$ can be for a given Q -value. For $Q \geq 10$, P_b is necessarily small, so that thermonuclear reactions in the bulk plasma make up the dominant contribution to $\langle \sigma v \rangle$. But for Q -values in the range 1 to 5, where $n_e \tau_E$ is small and P_b is large, one has considerable choice in the type of reactor-plasma operation. (We note that $Q \geq 5$ is probably required for a successful pure-fusion reactor, while $Q \sim 1$ is adequate for applications such as fissile breeding or materials testing; see Sect. 9.) Low- $n_e \tau_E$ highly driven reactors are often referred to as "power amplifiers."

Figure 15 illustrates the fusion reactivities of ideal CIT [35], TCT (D-T) [98] and thermonuclear plasmas. The CIT and TCT systems have reactivities at least several times the thermonuclear value for T_i up to about 25 keV.

At moderate temperature the beam-driven systems give $Q = 1$, for example, with a reduction in the $n_e \tau_E$ requirement of at least one order of magnitude from that given by the "Lawson criterion" [97]. In the purely beam-driven modes larger Q is obtained only by raising T_e , which increases the fast-ion slowing down time, and therefore by raising $n_e \tau_E$. But if T_e must be kept constant, Q can also be increased in the TCT case by reducing the tritium enrichment of the bulk plasma and increasing $n_e \tau_E$ (for example, by going to a larger plasma), so that bulk-plasma reactions contribute importantly to $\langle \sigma v \rangle$ (see Sect. 6.4).

In the following we are concerned principally with D-T plasmas because of their large reactivity and associated large power density, although the possibility of employing other fuels is discussed briefly in Sect. 6.6. We begin with low- $n_e \tau_E$ reacting plasmas, since these are of the size closest to present tokamak experiments, and indicate how performance evolves as $n_e \tau_E$ increases.

6.2 COUNTERSTREAMING-ION REACTORS (CIT or CBT)

In CIT (CBT) operation, injected neutral beams are used to stack large densities of oppositely traveling D and T ions. The steady-state counterstreaming velocity distributions shown in Figs. 6(a) and 16 can be set up by introducing all ions by means of neutral beams injected parallel and antiparallel to the toroidal magnetic field. At the beginning of a cycle the energetic atoms are trapped on an Ohmic-heated low density ($n_e \sim 10^{13} \text{ cm}^{-3}$), low temperature ($T_e \sim 1 \text{ keV}$) tokamak plasma. During injection for several slowing-down times (a period typically less than 0.5 s), n_e is built up to the range 3 to $6 \times 10^{13} \text{ cm}^{-3}$, T_e reaches a steady-state value of at least several keV, and the ion population becomes dominated by energetic ions —

provided that the inward flow rate of cold particles from the wall region is kept small compared with the rate of particle injection by the beams. In the steady state, injected energetic neutrals are trapped by charge exchange and impact ionization with the ions in the energetic distributions. In calculating λ_t for energetic-ion plasmas, it is important to average $\langle \sigma_{cx} v_{bi} \rangle$ and $\langle \sigma_I v_{bi} \rangle$ over the plasma ion velocity distribution [35]; see Eq. (4.2).

The CIT density build-up is to be distinguished from proposed thermal-plasma schemes in which n_e is increased by continuous ionization of background gas by beam-heated electrons [99]. The CIT scheme depends on a very low value of cold particle recycling into the reacting plasma. Very low recycling has been obtained experimentally in the Alcator device [100], for example, where the density drops continuously in time unless puffs of gas are injected into the plasma. In the ATC device the same result is obtained by titanium gettering of the wall [101]. A magnetic "unload" divertor [102] is expected to serve a similar role in inhibiting recycling.

6.2.1 Ideal CIT Performance. By suitable choice of the injection energy, the relative displacement velocity of the counterstreaming distributions can be made near the optimal value for generating fusion power, which occurs principally by nearly head-on collisions between counterstreaming D and T ions. If completely decelerated ions are removed from the plasma in a time τ_p that is short compared with the fast-ion slowing down time, and if recycling is prevented, virtually all plasma ions will be found in the toroidally directed beam-like distributions. (When $T_e \geq W_b/4$, however, all the energetic ions tend to thermalize with each other before completely decelerating, so that most of the directed motion is lost [34]; see Sect. 2.4.) The velocity distributions $f_h(v)$ of Figs. 6(a) and 16 were calculated numerically from the Fokker-Planck equation, assuming that energetic ions are perfectly confined

until they decelerate to an energy $W_1 = 2T_e$, and that ions with energy $W < 2T_e$ are lost at a rate $\tau_p^{-1} > \tau_{Ee}^{-1}$ [14]. T_e is determined by $n_e \tau_{Ee}$ according to Eq. (5.4), and is shown in Fig. 17(a).

The fast-ion slowing-down time satisfies

$$\frac{\tau_s}{\tau_{Ee}} \approx \frac{W_b - W_1}{3/2 T_e} \gg 1 \quad (6.2)$$

Evidently, the total ion lifetime τ_h must be several times τ_{Ee} . In considering diffusion mechanisms for ion loss, we note that the plasma conditions of interest fall into the so-called trapped-electron regime [12]. Theory predicts that the particle diffusion rate caused by trapped-electron instabilities is several times less than the energy diffusion rate [103], so that if τ_{Ee} is determined in large part by this instability, it is reasonable to expect that $\tau_h \gg \tau_{Ee}$ as required for CIT operation.

The fusion reactivity of the ideal CIT system approaches that of a beam-target (TCT) system requiring 3 times the injection voltage (see Fig. 15). For a given T_e the reduced bulk-ion friction on the energetic ions in the CIT results in Q -values somewhat larger than that of a TCT [14,35,104]. (Here we calculate Q by integrating fusion reactivity over the D and T velocity distributions, and dividing by the injected power.) Figure 17(b) shows that for injection energy of 60 keV, $Q = 1$ ("break-even") is attained, ideally, at $n_e \tau_E \approx 2 \times 10^{12} \text{ cm}^{-3} \text{ s}$, corresponding to $T_e \approx 3.3 \text{ keV}$, with $n_e \tau_h \approx 1.5 \times 10^{13} \text{ cm}^{-3} \text{ s}$. If fusion alpha particles are retained in the plasma, the ideal Q reaches 5 when $n_e \tau_E \sim 3 \times 10^{13} \text{ cm}^{-3} \text{ s}$ [34].

Figure 18 shows the 14-MeV neutron production rate from ideal CIT plasmas of circular cross section [105] as a function of plasma radius and W_b , when $n_e \tau_E$ is such that $T_e = 4.0 \text{ keV}$. For a given a_p the neutron output is proportional to $n_e^2 \langle \sigma v \rangle$. At small W_b , n_e is determined by the neutral-beam

penetration requirement, which is taken here as trapping-length $\lambda_t = a_p$. At large W_b , n_e is limited by the allowed plasma pressure, which is chosen so that the poloidal beta of the plasma is $\beta_p = (2/3)R_o/a_p$. Therefore, for each radius there is an optimal W_b for maximum neutron output. Evidently for tokamak plasmas of modest size ($a_p = 40\text{-}50$ cm), and using beam voltages at which high-power neutral beams are now available ($W_b \sim 40\text{-}60$ keV), neutron intensities in the range $10^{18}\text{-}10^{19}$ s⁻¹ are attainable.

6.2.2 Practical Limitations to Performance. The degree to which ideal CIT performance can be attained depends most crucially on minimization of the warm-ion population. The proportion of warm ions depends first on the residence time τ_p of decelerated ions in the reacting region, and second on the influx of neutrals and cold plasma to this region. For small machines, τ_p should be acceptably small, and the most severe limitation on CIT performance is expected to be set by cold-particle influx. When $Z_{\text{eff}} = 1$ about half the injected energetic atoms are trapped by charge exchange with ions streaming in the direction of beam injection, and this process is responsible for loss of up to 10% of the injected power [35]. A potentially much more serious problem is the influx of recycled neutrals from the scrape-off region surrounding the discharge, which results in the maintenance of a cold plasma surrounding the hot reacting region. In principle recycling can be practically eliminated by means of an "unload" magnetic divertor [102] that exhausts rapidly from the torus those ions diffusing out of the reacting plasma. But provision for a magnetic divertor is a complex and costly affair, and divertor effectiveness remains to be demonstrated. (Nevertheless, recent results with a "bundle" divertor in reducing recycling are encouraging [106].) An alternative approach is to install baffle-protected gettering surfaces above and below the plasma in an elongated vacuum chamber. The density of the

neutral gas in the scrape-off region ($T_{\text{gas}} \sim 10 \text{ eV}$) can be kept sufficiently small ($\approx 10^{11} \text{ cm}^{-3}$), provided that a sizeable fraction of the scrape-off layer has direct access to the gettering region [34].

In order to determine the proximity of a practical operating point to the ideal CIT regime, an analysis with a radial transport model is evidently needed. Initial calculations of reactor performance have been made with a coupled Fokker-Planck radial transport code [34,107]. Figure 19 shows how Q decreases with relative cold particle influx into the reacting region. With increasing cold-particle influx, beam-target reactions eventually become the dominant source of neutrons. In the limiting case of interest, the plasma is operated in the TCT regime, driven and fueled by injected D and T beams [108]. For $W_b \sim 60 \text{ keV}$, the Q -value (i.e., neutron production for a given beam power) is reduced by a factor of 5 to 7 from their values in ideal CIT operation. Furthermore, $n_e \tau_E$ must be increased to maintain the same T_e .

While a small impurity content in a CIT plasma facilitates maintenance of the plasma current by the energetic ions (see Sect. 5.3), an accompanying deleterious effect is the enhancement of pitch-angle scattering, which undermines the advantage of head-on collisions. Evidently this effect is less important when thermal ions make up a substantial fraction of the total ion population. A second deleterious effect is the additional Coulomb drag on energetic ions by the impurity ions and their neutralizing electrons.

The angular scattering rate is reduced at increased injection energy. At larger W_b one can also take better advantage of target-plasma reactions, so that Q -values are not as severely degraded in the event that the warm-ion population is substantial. Finally, the charge-exchange loss that accompanies beam trapping decreases with increasing W_b . For these practical reasons the optimal energy in CIT operation is probably 80 to 100 keV for D⁺ (120 to 150 keV for T⁺), where the efficiency of neutral-beam injectors is still high

(see Sect. 7.1). At still larger W_b the hot-ion reactivity is seriously reduced, while for a given P_b the injection rates of fast particles and momentum also decrease.

6.2.3 Limitation to Neutron Flux. Attaining Q -values larger than 3 to 4 demands operation at large T_e , so that $\tau_{hh}/\tau_{se} < 1$ [cf. Eq. (2.10)], and the ion distributions approach isotropic ones, as in Fig. 6(c). Reactor operation then begins to resemble the beam-heated thermonuclear mode (Sect. 6.4). Operation in the range $Q \sim 1$ to 3 with $n_e \tau_E < 10^{13} \text{ cm}^{-3}$ seems realistic for small machines, with W_h in the range 40 to 100 keV (D°). However, the small plasma size permitted for adequate neutral-beam penetration with tangential injection ($\langle n_e \rangle_p \approx 5 \times 10^{15} \text{ cm}^{-3}$) severely restricts the attainable 14-MeV neutron wall loading, ϕ_w . Here $\phi_w = 0.8 \times P_f \times (\text{plasma volume/wall area})$, where

$$P_f = n_e^2 \frac{n_D}{n_e} \frac{n_T}{n_e} \langle \sigma v \rangle E_f \quad (6.3)$$

and n_e is limited by

$$n_e (T_e + \frac{2}{3} \frac{n_i}{n_e} \bar{W}_i + \frac{2}{3} \frac{n_h}{n_e} \bar{W}_h) \approx 2.5 \times 10^{13} \left(\frac{B_t S}{qA} \right)^2 \beta_{pmax} \quad (6.4)$$

Here B_t is in kG, T_e and W are in keV, $\beta_{pmax} \sim R_o/a_p$ is determined by MHD equilibrium or stability, and S is the plasma circumference divided by $2\pi a_p$. ($S = 1$ for a circular plasma.) For practical CIT parameters, $P_f = 0.5-1.5 \text{ W/cm}^3$ and $\phi_w = 0.1-0.5 \text{ MW/m}^2$ [35,105].

An alternative approach to an energetic-ion reacting plasma is to establish a less anisotropic ion distribution by means of very oblique injection. For $W_b \sim 100 \text{ keV}$ (D°), ϕ_w would be degraded by relatively little from its value in counterstreaming operation, and the factor of 1.5 or more

increase in n_e now permitted for acceptable penetration would result in a substantial increase in ϕ_w .

6.3 TARGET-PLASMA REACTORS (TCT)

6.3.1 Ideal TCT Performance. The "ideal" TCT consists of a pure tritium target plasma whose temperature is maintained by injected energetic deuterons, which undergo fusion reactions with the bulk tritons while thermalizing. The energy multiplication factor of the injected beam ions of velocity $v(t)$ is

$$Q_b = \frac{\int_0^{\tau_s} n_T \sigma(v) v E_f dt}{W_b} = \frac{\int_0^v n_T \sigma(v) v E_f \frac{dv}{dt} dv}{W_b} \quad (6.5)$$

In early work on TCT performance [9,13,109,110], T_i was taken to be zero and dv/dt was calculated from Eq. (2.7). The resulting energy gains are shown in Fig. 20(a), assuming that $n_T = n_e$ (i.e., $\Gamma \ll 1$). The increase in Q_b with T_e is due to the continuous reduction in electron drag and therefore a continuous increase in τ_s . Note that the optimal injection energy is considerably larger than the energy W_M for maximum reactivity ($W_M = 125$ keV for D-T). (The larger Q -values in Ref. [110] are apparently due to erroneous values of the D-T cross section.)

For finite T_i the Q -value is significantly modified for several reasons:

(1) The beam-target reactivity is increased at $W < W_M$ and decreased at $W > W_M$ [111]; (2) Enhanced energy diffusion results in a modification of $f_h(v)$ [24,65] (see Sect. 2.2); (3) Coulomb drag by the bulk ions is reduced. The first two effects are extremely significant at low values of W/T_i , and lead to large increases in Q_b for $W_b < 100$ keV [65,112]. Figure 20(b) taken from Ref. [112] shows Q_b when $T_i = T_e$, the most practical situation when $\Gamma \approx 1/2$.

(Ref. [65] indicates even higher Q_b at low values of W_b/T_i ; the discrepancy may be due in part to the use of different values for the fast-ion lower cut-off energy.) As in all calculations of Q for beam-injected plasmas, the plasma density enters only through $\ln \Lambda$ factors. Even at the highest T_e it is possible to reach Q -values no greater than 3 in ordinary TCT operation.

(The injection of triton beams into a deuterium target plasma ($T \rightarrow D$) gives lower values of Q_b than for $D \rightarrow T$ [108]. As expected, the $T \rightarrow D$ Q -values peak at $W_b \approx 50\%$ larger than the optimal W_b for $D \rightarrow T$. At very low T_e where electron drag is the dominant slowing-down mechanism, the Q -values are about the same. But the friction on background ions is 50% larger for $T \rightarrow D$ than for $D \rightarrow T$, and in the limit of very large T_e , $Q_{T \rightarrow D} \approx 2/3 Q_{D \rightarrow T}$.)

In order to achieve maximum Q for the entire beam-plasma system, we require that $\tau_h > \tau_s$ and that all external power to the plasma be provided by the injected beams (i.e., Eq. 5.1 with $P_{OH} \approx 0$). Let us now consider the maximization of fusion power density. The injected beam power can be increased with no change in T_e if τ_E is continuously reduced. Since the plasma pressure must be limited, however, the resultant increase in Γ necessarily leads to a decrease in n_e and therefore in target density n_T [cf. Eq. (6.4)]. Furthermore, charge neutrality demands that the large fast-ion density itself must reduce n_T/n_e . Simultaneous solutions of the equations of power balance, pressure limitation, charge neutrality, and Q_b demonstrate that for a given T_e there is an optimal value of n_h/n_e to maximize P_f , accompanied by optimal values of Γ and $n_e \tau_E$ [113]. These results for $W_b = 200$ keV are shown in Fig. 21. Evidently the attainable P_f is nearly inversely proportional to T_e . Under the restriction of constant pressure, lower temperatures allow one to increase the plasma and beam densities, and therefore the fusion rate. Furthermore, the beam-plasma reactivity is relatively insensitive to plasma temperature (in marked contrast with the thermonuclear case; cf. Fig. 15).

Another important feature is that for maximum P_f , τ_E must be a factor 1.3-2 times smaller than $\tau_s \leq \tau_h$, which is equivalent to the condition that $\Gamma = 0.5$ to 1. This range of Γ has in fact been achieved in recent beam-injection experiments, with apparently classical slowing-down of the energetic ions (see Sect. 8.3).

Figure 21 shows that $Q_b = 1$ ("breakeven") is attained at $n_e \tau_E \approx 10^{13} \text{ cm}^{-3} \text{ s}$ when $T_e = 5 \text{ keV}$. This $n_e \tau_E$ -value is an order of magnitude smaller than that required by the Lawson criterion for a thermonuclear plasma at similar temperature, a result that follows directly from the much larger fusion reactivity of the TCT system.

For many applications, economic considerations tend to favor the maximization of P_f and the related quantity ϕ_w , since the cost of 14-MeV neutrons from the reactor is roughly inversely proportional to P_f . Nevertheless, it is also important that Q be larger than about 0.5, since the required investment in neutral-beam injectors, and the quantities of electrical power and tritium that must be circulated, are nearly inversely proportional to Q . While the largest P_f is attainable at small T_e , there are several practical reasons for preferring the range $T_e = 5\text{-}10 \text{ keV}$: (1) Q_b increases with T_e ; (2) $P_b = P_f/Q_b$ increases extremely rapidly at small T_e , so that a substantial fraction of the vacuum wall area might have to be taken up by neutral-beam apertures. (3) The large values of n_e required to maximize P_f at low T_e make beam penetration difficult unless the plasma radius is very small, or $W_b \gg 200 \text{ keV}$, or VB-drift injection is employed. Consequently, conceptual design studies to date (see Sect. 9) have employed the 5 to 10-keV temperature range.

6.3.2 Effect of Impurities on TCT Performance. Impurity ions have four principal effects on TCT performance: (1) Reduction of neutral-beam

penetration length [54,55]. (2) For a given plasma pressure, the triton-target population must be depleted. (3) Radiation losses are increased, but these losses are easily compensated by the large injection power density characteristic of high- Γ TCT operation. (4) Neoclassical diffusion of the fast ions is enhanced by a factor Z_{eff} . The last two effects tend to be inconsequential for $Z_{\text{eff}} \geq 5$, provided that adequate beam penetration is maintained. Methods of alleviating the penetration problem are discussed in Sections 4.1 and 4.2.

Let us consider briefly the triton depletion problem. When $Z_I/A_I = 1/3$, then $[Z] = 1$, v_c is unchanged, and the Coulomb drag on the energetic ions is identical with that of a pure triton-target plasma. Then $\tau_s = \text{constant}$ and Q_b is reduced from its ideal value by simply the factor n_T/n_e . If $[Z] > 1$, the Coulomb drag is increased, and vice versa. For a given impurity concentration Z_{eff} , the equations of power balance, pressure, and Q_b can again be optimized for maximum P_f [113]. Figure 22 shows the variation of $P_{f \text{ max}}$ and Q_b for $T_e = T_i = 6.0$ keV and a range of impurity concentration. Evidently P_{max} is reduced by a smaller factor than the reduction in Q_b , for a given impurity content. Provided that effective beam penetration techniques are available, stringent impurity control would seem to be unnecessary in a TCT whose main purpose is neutron production at low Q [110a]. Even for $Z_{\text{eff}} = 10$ caused by ions sputtered from the vacuum wall (e.g., iron), the reduction in P_f is at most 30%.

6.3.3 Energy Clamping. Section 2.2.1 considered $f_h(v)$ when external power is applied to the fast ions to balance the frictional drag of the bulk-plasma, and thereby "clamp" the fast-ion energy. The effectiveness of clamping actually decreases with time because of energy diffusion, but it is interesting

to consider the ideal maximum Q that can be attained by this technique. For a single ion, one has

$$Q_c = \frac{n_T \sigma v E_f \tau_c}{W_b + |dW_b/dt| \tau_c} \quad (6.6)$$

where τ_c is the clamping period. For $\tau_c \gg \tau_s$, $Q_c \rightarrow n_T \sigma v E_f / |dW_b/dt|$. Evidently the ideal fast-ion energy W_{cl} is the value that gives maximal ratio of fusion reaction rate to plasma drag. Figure 23 shows the variation of W_{cl} and $(Q_c)_{\max}$ with T_e , when $T_i = 0$ [13,37]. The optimal injection energy is only 53 to 65% of the optimal value in the conventional TCT scheme (Fig. 20), while the fusion gain in clamped operation can be nearly twice as large. When T_i is finite, σv is reduced at $W = W_{cl}$ so that $(Q_c)_{\max}$ is reduced, but finite T_i results in an increase in Q_c at values of W somewhat lower than W_{cl} [114].

Reference [30] describes a graphical means for estimating Q_c/Q . Defining $S(W) = n_T \sigma v E_f / |dW/dt|$, it is easily shown that

$$\frac{(Q_c)_{\max}}{Q_{\max}} = \frac{W_b S(W_{cl})}{\int_0^{W_b} S(W) dW} \quad (6.7)$$

where W_{cl} and W_b are the optimal injection energies for clamped and unclamped operation, respectively. Evidently, energy clamping affords the largest improvement for fusion reactions with highly peaked $\sigma(v)$.

In practice, $f_h(v)$ must have finite width because of energy diffusion and pitch-angle scattering (and perhaps also microinstability), but even for a very broad $f_h(v)$, a significant increase in Q is possible [25,30,115]. Several techniques have been proposed for energy clamping. Slow major-radius compression [28,30,116] is especially powerful with ions having large $v_{||}/v$,

and the effectiveness of this technique has been demonstrated on the ATC device (see Sect. 8). Compression-expansion cycles utilizing fast-ion clamping and direct energy recovery comprise a possible mode of reactor operation [116]. Clamping by the toroidal electric field [27] appears also to be effective, but substantial enhancement of Q seems unattainable by this method because of the accompanying large ohmic dissipation [117]. An attractive technique is a rigid displacement of the plasma column at the transit frequency around the torus of phase-bunched fast ions [29], but this method has been scarcely investigated. Another proposed method is the application of harmonic ion-cyclotron resonant heating; the production of fast-ion tails by this means has been demonstrated on several devices [118], but the method seems likely to be useful only for systems of very low Γ .

As discussed in Sect. 5.1.3, the principal use of energy clamping is most likely to be the enhancement of the heating effectiveness of injected fast ions -- rather than the attainment of enhanced Q -values. Clamping techniques that rely on the ions having large $v_{\parallel}/v$ can possibly be applied also to CIT plasmas, while application to "exotic-fuel" plasmas employing ions of $Z \geq 2$ will have to make use of techniques that favor more isotropic distributions.

6.3.4 50:50 D-T Operation. Maintaining a highly enriched tritium plasma in the face of continuous deuteron injection is a formidable problem, and probably requires the continuous injection of tritium pellets or liquid jets. In addition, all ions escaping the plasma must be removed from the torus by means of an unload divertor with zero recycling coefficient, in order to avoid a gradual build-up of deuterons in the bulk plasma [113]. Therefore approximating ideal TCT operation ($D \rightarrow T$) requires perfection of both pellet and magnetic divertor technologies. Other proposed fueling techniques include gas

injection, low voltage beam injection [119], or simply restricting operation to short pulses with fueling carried out between pulses. Analysis of these methods [108] finds them to be unattractive for practical reactor operation.

In the absence of pellet injection technology, the optimal tritium strategy appears to consist of treating the deuteron and triton plasma components in a perfectly symmetric manner. That is, D° and T° beams are injected with the same velocity, and with sufficient intensity to maintain a 50:50 D-T bulk plasma — which is fueled by the injected beams together with plasma recycling. The values of Q and P_f are necessarily smaller than those attainable in the ideal D-T situation, although significant fusion power is also produced from thermonuclear reactions, and especially from reactions among the energetic D and T ions themselves [104,108,113]. In the limiting case where the superthermal-ion density becomes a large fraction of the electron density, one approaches the "energetic-ion" class of reactors (such as the CIT). The advantages of this technique are that proper fueling is provided automatically, the power output is well-defined and constant during the pulse (in the absence of impurity buildup), and the fractional burn-up of tritium is relatively large. The length of the power pulse is limited fundamentally by the flux swing in the transformer core.

The 50:50 D-T mode of operation of heavily beam-driven plasmas is analyzed in detail in Ref. [108], which employs fusion reactivities calculated separately for energetic-ion reactions, target-plasma reactions, and thermonuclear reactions. Multi-angle injection is specified, so that the hot ions are assumed to have isotropic velocity distributions. For conditions where $Q \sim 1$, beam-target reactions are dominant, although energetic-ion reactions make a substantial contribution to P_{out} when $n_h/n_e \gtrsim 0.2$. With increasing n_h/n_e , a given Q -value $\lesssim 1$ can be maintained with $n\tau_E$ -values substantially below $10^{13} \text{ cm}^{-3} \text{ s}$, as illustrated in Fig. 24.

When tritium is injected via energetic neutral beams (rather than by pellets or gas), the fractional burn-up of tritium per pass is maximized. For $T_e = 6$ to 10 keV and $W_b = 200$ keV, the burn-up rate once-through the reactor is 2 to 3% [108].

6.3.5 Maximum Neutron Wall Loading. Certain potential applications of tokamak fusion reactors require the largest possible 14-MeV neutron flux, despite accelerated damage rate to the reactor first wall (see Sect. 9). Given T_e and W_b (which may be determined in practice by efficiency of beam production), two critical reactor parameters that determine the neutron wall loading ϕ_w are the magnetic field B_M at the TF coils, and the plasma height-to-width ratio κ . (An early analysis [120] showed that for a circular plasma at constant B_M the maximum injected power must decrease with Z_{eff} , which lowers the n_e for acceptable penetration.) In the analysis of Ref. [121], the distance R_1 between the major axis and the inner edge of the plasma is fixed, a procedure that is justified a posteriori. Given B_M , κ , and a definite mode of plasma operation — that is, fixed T_e , W_b , and $n_h/n_e = 1$ — one seeks the value of a_p that gives maximum ϕ_w . There are two constraints on the plasma density, n_e : (i) $n_e a_p \leq C_1$, where C_1 is a constant determined by adequate neutral-beam penetration (see Sect. 4.1). (ii) The plasma pressure must not exceed the maximum allowed by MHD equilibrium [see Eq. (6.4)]. At large a_p constraint (i) determines n_e . By decreasing a_p we can increase n_e (and therefore ϕ_w) until constraint (ii) is reached. Maximum ϕ_w is attained when [121]

$$n_{eM} = \frac{C_1}{a_p} = C_2 \frac{a_p (1 + \kappa)^2 B_M^2}{(R_1 + a_p)^3} \quad (6.8)$$

where C_1 and C_2 are constants. When $a_p \ll R_1$, several approximate results are readily derived: (a) $\phi_w \propto B_M$. (b) $\phi_w \propto \kappa$. (c) Plasma current $I_p \propto 1/B_M$, and independent of κ . (d) Injected power $P_{\text{beam}} \propto \kappa$, and independent of B_M . The nearly linear dependence of ϕ_w on B_M or κ differs strikingly from that of an ignited thermonuclear reactor, since in the latter case the plasma radius need not be reduced with increasing n_e to accommodate neutral beams.

With $b/a = 2$, $B_M = 14$ T, $W_b = 200$ keV, $Q \approx 1.0$, one has maximum $\phi_w \approx 3.5$ MW/m². But the large injected power and relatively small plasma volume that are required to achieve such large neutron currents aggravate the problems of beam access and particle throughput. In practice the maximum neutron intensity from a TCT plasma is likely to be determined by the torus pumping capability that is required to control n_e in the face of an injected particle input of some 10^{22} ions/s [108,122]. A viable solution to the pumping problem may possibly be found with a poloidal or bundle magnetic divertor that can exhaust ions at high velocity from the torus [122]. However, divertor installation poses the still largely unresolved practical problems of shielding, power requirements, and forces associated with the divertor field coils.

6.3.6 Spatial Profile Effects on Performance. Most analytic studies of beam-driven tokamak plasmas have employed spatially homogeneous plasma models (except for compression). If $n_h(r)$ is sufficiently peaked so that a majority of hot ions slow down in a region where T_e is larger than the spatially averaged value $\langle T_e \rangle$, then Q is expected to be substantially larger than the value computed for a temperature $\langle T_e \rangle$. Reference [123] investigates the performance of a one-dimensional TCT plasma with radial profiles $T_e(r) = T_{ec} (1 - r^2/a_p^2)$ and $n_h(r) \propto \exp(-r^2/r_s^2)$, where r_s is a constant. The hot-ion density is one-half its peak value at $r_H = r_s \sqrt{\ln 2}$. With $n_e \tau_E = 1.5 \times 10^{13} \text{ cm}^{-3} \text{ s}$, the required $\langle T_e \rangle$ for $Q = 1$ decreases from 6 keV to 3 keV when r_H/a_p decreases from 0.7 to

0.3. This advantage is not easily realized in a reactor, where one wants generally the largest possible plasma density in order to maximize power output, so that with the conventional injection method at least, the neutral-beam deposition profile is unlikely to be peaked significantly on axis.

The one-dimensional transport codes that have been used to investigate the performance of beam-driven tokamak plasmas [57,124] predict Q -values that are in essential agreement with the zero-dimensional analytic or numerical models. These codes have proved most useful for determining the temporal evolution of plasma conditions, accounting in particular for charge exchange and the effects of impurity build-up that results from wall sputtering. The particle and thermal diffusion coefficients in these codes are somewhat arbitrary, however, often being derived from ad hoc "adjustments" to untested theoretical transport models in order to give energy confinement times in essential agreement with "reasonable" extrapolation from presently observed tokamak performance. To date these codes have not included the transfer of decelerated ions from the hot to warm populations; the hot ions simply evaporate at a prescribed $W \sim 2T_i$. Coupled Fokker-Planck radial transport codes are needed to treat accurately the fate of injected energetic ions [107]. When experimentally verified transport coefficients are finally obtained, such "hybrid" codes will be invaluable for the optimization of reactor design.

6.4 BEAM-DRIVEN THERMONUCLEAR REACTORS (BDTN)

6.4.1 Power Density versus Q . Figure 21(a) compares P_f for a 50:50 D-T thermonuclear plasma with the maximum P_f in a TCT plasma of the same total plasma pressure. If (nonreacting) heating beams must be used to maintain the temperature of the thermonuclear plasma, it is assumed in that figure that τ_E is sufficiently large so that the beam pressure is negligible. For TCT operation

at $Q \sim 1$ (say $T_e = 6$ keV), P_f is a factor of 4 larger than the maximum P_f attainable with the thermonuclear plasma in the optimal temperature range for the latter (12-15 keV). Nevertheless, thermonuclear operation does permit the attainment of the largest Q -values. When reacting ion beams are used to maintain the plasma temperature, one can demonstrate a continuous evolution from "pure TCT" operation to large- Q thermonuclear operation.

The total fusion gain Q of a TCT can be raised beyond the values shown in Fig. 20 by increasing $n_e \tau_E$ as well as the proportion of D in the bulk plasma, so that a larger fraction of the fusion energy comes from thermonuclear reactions [13,126]. The overall power gain is

$$Q = \frac{n_{iD} n_{iT} \langle \sigma v \rangle E_f + Q_b P_b}{P_b} \quad (6.9)$$

where P_b is given by Eq. (5.1) and Q_b by Eq. (6.5). An increase in $n_e \tau_E$ allows a reduction in P_b , so that the relatively small thermal fusion power can eventually surpass $Q_b P_b$. With increasing $n_e \tau_E$ the optimal value of n_{iD}/n_{iT} for maximizing Q increases toward 50%. Figure 25(a) shows the evolution of Q and P_f for $T_e = T_i = 8$ keV [113]. Evidently maximum P_f is attained in "pure TCT" operation, with an $n_e \tau_E$ that allows only small Q . As $n_e \tau_E$ increases and n_{iD}/n_{iT} approaches unity, P_f goes over to the purely thermonuclear value. The decrease in Γ with increasing confinement time is illustrated schematically in Fig. 2.

Figure 25(b) demonstrates the Q vs P_f trade-off for a number of plasma temperatures. Each point on a constant- T_e curve corresponds to a particular value of $n_e \tau_E$, with n_{iD}/n_{iT} optimized for maximum Q . The important feature is that the plasma conditions for maximum P_f differ markedly from those for maximum Q . Reactor applications such as materials testing and fissile-fuel breeding typically demand large neutron flux, but can tolerate relatively low

Q ; for such applications the purely beam-driven regimes on the right-hand side of Fig. 25(b) are most attractive. On the other hand, "stand-alone" fusion power plants generally require $Q \geq 5$, so that reactor operation must be in the nearly pure thermonuclear regimes on the left-hand side. At very large $n_e \tau_E$, $P_b \rightarrow 0$ and the plasma temperature is sustained entirely by alpha-particle heating.

(This analysis is not peculiar to tokamak operation or even to toroidal reactors, but in fact is applicable to all fusion plasmas driven by reacting ion beams, when the injected beams provide the dominant external power input, and when loss cones in velocity space are relatively unimportant. See Ref. [125], for example.)

6.4.2 Driven Thermonuclear Reactors with Impurities. The energy confinement time τ_E in Figs. 24 and 25 includes implicitly all sources of power loss from the plasma. For $n_e \tau_E \geq 5 \times 10^{13} \text{ cm}^{-3} \text{ s}$, line radiation from impurities could be the dominant power loss, and the impurity ions also reduce the plasma reactivity. These effects can easily prevent ignition [127], but useful large- Q operation is still possible when reacting ion beams are injected to supplement the alpha particle heating. Reference [126] shows, for example, how Q decreases with increasing iron content of a D-T plasma operating near ignition. For typical machine parameters ignition is unattainable with an iron content of 0.15%, but even with 0.5% impurity ($Z_{\text{eff}} = 4.2$), $Q = 15$ is obtained when the plasma is driven by 200-keV deuteron beams. If the beams can be generated with 70% efficiency, then the net plant efficiency [128] is an acceptable 30%.

Even in the absence of an impurity-ion population, the BDTN mode offers the advantage over an ignited plasma of thermal stability control. Effective control is probably achieved with $P_b/P_\alpha \approx 0.2$, or $Q \approx 25$, a regime sufficiently close to ignition so that the net plant efficiency is essentially unaltered [128]. Whatever the motivation for a high- Q quasi-stationary power amplifier,

the principal concern must be to ensure adequate neutral-beam penetration into the rather opaque high- $n_e \tau_E$ plasma.

6.5 HEATING TO IGNITION

Neutral-beam injection into tokamak plasmas was first proposed as a means of heating to ignition temperature [8]. The use of reacting beams for heating is even more profitable, because the alpha particles produced by beam-target and beam-beam reactions reduce the external start-up energy requirement [13, 81]. For short-pulse ignition devices, the neutron production from beam-induced reactions can contribute significantly to the overall energy balance for the cycle. Figure 26 shows the fraction of the plasma energy content at ignition supplied by the injected beams [81]. Evidently, this fraction can be as large as 20%. Taking into account the beam-generated neutrons, the net electrical energy required for heating is reduced typically by 20 to 40%. Still further advantage over other proposed heating methods could be realized by applying an energy clamping technique to replace some injector capacity with energy stored in the ohmic or equilibrium-field circuit (see Sect. 5.1.3).

6.5.1 Minimum Beam Energy for Ignition. The energy confinement time of present Ohmic-heated plasmas with $q \sim 4$ appears to follow the scaling $\tau_E \sim K \langle n_e \rangle a_p^2$, where the constant $K \approx 8 \times 10^{-19}$ cm-s [89,100]. The indicated dependence on density and radius is not yet explained by any theoretical model. One can postulate that energy confinement in the higher temperature regimes of reactor interest can be represented as $\langle n_e \rangle \tau_E \approx K n_e^2 a_p^2 G(T_e, T_i, B_t)$, where the function G remains to be determined by experiment. If G does not decrease significantly with T_e or T_i then the plasma opacity required for $\langle n_e \rangle \tau_E = 3 \times 10^{14} \text{ cm}^{-3} \text{ s}$ is $n_e a_p \approx 2.8 \times 10^{16} \text{ cm}^{-2}$. Figure 7 indicates that the required

D^o beam energy is 300 keV when $Z_{\text{eff}} = 1$. Efficient injection at this energy, or at the higher values required when $Z_{\text{eff}} > 1$ or if the temperature dependence of G is unfavorable, can be obtained only with negative ion beams, an undeveloped technology (see Sect. 7). Alternatively, W_b can be reduced substantially if the ripple-trapping injection technique proves successful. The latter method also provides a means of heating to ignition the very dense reactor plasmas proposed in Ref. [89], where $\langle n_e \rangle_a = 5 \text{ to } 10 \times 10^{16} \text{ cm}^{-2}$ in the examples given therein. For this application asymmetric ripple produced by auxiliary coils would be most desirable, since the coils could be turned off during the thermonuclear burn, thereby eliminating the ripple enhancement of the diffusion of alpha particles born on banana orbits (see Sect. 4.2).

The total beam power required for ignition of circular cross-section plasmas is $P_{\text{ign}} \approx 6\pi^2 T_{i\alpha} (n_e a_p)^2 R_o / n_e \tau_E$, where $T_{i\alpha}$ is the temperature at which $P_\alpha \sim P_b$. Calculated values of P_{ign} are in the range 60 to 300 MW.

6.5.2 Dynamic Start-Up. A dynamic start-up technique has been proposed in which 100-keV beams are used to heat the cold outer layers of an expanding-limiter plasma [130]. Provided that radiation losses are relatively small, the temperature of the inner layers remains high during the heating stage, the inner plasma being insulated by the beam-heated outer layer. A related technique proposes to use even lower energy beams ($W_b \sim 10\text{-}20 \text{ keV}$) to build up the plasma density layer by layer by deceleration of the injected energetic ions [131]. A variant of this technique proposes somewhat higher energies ($W_b \sim 60 \text{ keV}$) with oppositely directed D and T beams in order to produce alphas and neutrons in highly reactive CIT-like operation, while the plasma layers are being formed [132]. These proposed dynamic start-up methods offer a potential means of reaching ignition with state-of-the-art neutral beams, but they may also introduce two serious problems: Excessive charge-exchange

loss of energetic ions, and aggravated sputtering of the first wall by hot neutrals escaping from the beam-maintained edge region. A Monte Carlo treatment of the charge-exchange neutral flux induced by 80-keV beam injection has found that the escaping neutrals impinge on relatively small regions of the first wall [133].

Reactor plasmas with size much larger than the minimum required for ignition can best be heated by first igniting a fractional-size plasma core by means of injected beams and adiabatic compression, and then allowing the bulk of the plasma to be heated by a propagating thermonuclear burn driven by alpha energy deposition in the ignited core [81,134]. By this technique the largest plasmas can be ignited with minimum neutral-beam voltage and minimum total energy investment. For expansion ratios ~ 3 , the injected beam need supply only a few per cent of the thermal energy of the final ignited plasma. Large energy gains (~ 100) are attainable even in a pulsed-ignition device, with the discharge terminated soon after the burn front reaches the wall [52]. If it is desired to trade off fusion power output for reduced tritium breeding requirements, as well as reduced 14-MeV neutron radiation damage, then the deuterium enrichment of the fuel mixture is made an increasing function of plasma radius.

While pulsed operation has the disadvantages of introducing enhanced thermal fatigue of the wall and blanket and possibly the need to install additional storage capability, the feasibility of pulsed high-Q operation as suggested here does not depend on the successful development of techniques for plasma refueling, ash removal, impurity control, or preservation of thermal stability — techniques that are essential for quasi-stationary ignited operation. Finally, we recall that the attainment of ignition at all depends on sufficiently benign electron loss mechanisms so that $T_e > T_i$ can be maintained. As shown in Fig. 13(b), for temperatures of practical tokamak

interest the bulk of the alpha energy is deposited in plasma electrons, which must heat the ions by Coulomb relaxation. If either radiation loss or electron thermal conductivity tends to "clamp" T_e at a relatively low value, then only driven systems with the injected energy going mainly to plasma ions are likely to be interesting. Energetic-ion plasmas such as the CIT comprise a particular class of such hot-ion, warm-electron tokamak reactors [135].

6.6 ADVANCED-FUEL BEAM-DRIVEN REACTORS

Employment of the TCT operating mode for "advanced fuel" cycles such as D-D, D- ^3He , D- ^6Li , and p- ^{11}B offers the same advantages in reducing the $n_e T_e$ requirement for achieving small Q-values as it does for D-T. Nevertheless, extremely large values of W_b and T_e are required to obtain even $Q_b = 0.5$ with non-DT fuels, as shown in Fig. 27 [136]. In principle substantial net energy production is possible even with $Q_h \ll 1$ for the neutron-free reactions (D- ^3He or p- ^{11}B) when coupled with high-efficiency direct energy conversion techniques, which recover all the energy of the unburned injected ions as well as of the charged reaction products [137,138]. However, the large circulating power and inevitably imperfect injector efficiency are likely to make such operation economically impractical unless $Q_b > 1$.

(Note that a D- ^3He plasma must have a substantial D-D reaction rate, and therefore cannot be truly neutron-free, although the relative D-D reaction rate can be made small by injection of ^3He beams into a D plasma [139]. The most practical future of the D- ^3He reaction, however, probably lies in its role as a catalyst in a D-D reactor, where the ^3He is one of the reaction products. The reaction may also prove useful for low-activation tokamak experiments in which the p and ^4He reaction products are to be used to simulate the plasma effects of D-T fusion alphas [140].)

Energy clamping provides substantially enhanced Q-values for those reactions such as $p\text{-}^{11}\text{B}$ with notably peaked fusion cross sections [cf. Eq. (6.7)]. Compression clamping is especially suited for the neutron-free reactions, since decompression at the end of the clamping period results in highly efficient direct energy recovery via the equilibrium-field circuit [30,137]. A reactor operating on compression-expansion cycles may also be useful with catalyzed D-D [137]; when the T and ^3He reaction products are burned, 62% of the total fusion energy produced is carried by charged particles.

The reactivity advantage of the CIT mode cannot be used profitably with reactions involving ions with $Z > 1$, because of the rapid pitch-angle scattering (proportional to $Z_i^2 Z_j^2$) that destroys the frequency of head-on collisions. But $Q = 1$ can be obtained in catalyzed D-D with $T_e = 70$ keV and $W_D = 300$ keV, an energy several times smaller than useful for TCT operation in D-D.

In a pressure-limited device such as a tokamak, the large particle energy required for advanced-fuel operation restricts n_e to low values. Although a beam-driven advanced-fuel power system with acceptable net plant efficiency could conceivably be demonstrated, the low values of P_f resulting from small n_e and small $\langle\sigma v\rangle$ make the advanced-fuel cycles unlikely contenders with the D-T cycle.

7. INJECTOR SYSTEMS

There are three critical areas in the high-power injection of energetic neutrals into tokamak plasmas: (1) The production of energetic neutrals with satisfactory efficiency; (2) the efficient transport of energetic neutrals into the torus; (3) adequate penetration of the energetic neutrals into the central plasma region. The third topic has been discussed in Sect. 4, and the first two are the subject of this section.

7.1 INJECTOR PERFORMANCE REQUIRED FOR REACTOR APPLICATIONS

Table 4 shows the desired performance of reactor-quality neutral-beam injectors. It is assumed here that only the conventional injection method is used, so that very large beam energies are required for the thermonuclear plasmas — even larger than indicated if $Z_{\text{eff}} \geq 3$. The four cases have similar total beam intensities (I_{beam}). Since fluxes of $0.02 \text{ A-equiv./cm}^2$ seem feasible for beam ducts with characteristic dimension $\sim 0.5 \text{ m}$ [141], the total beam-duct area ($1\text{--}5 \text{ m}^2$) is typically only about 1% of the torus wall area.

The overall injector efficiency η_I is the ratio of the beam power injected into the torus, to the net electric power consumed by the injector. Regardless of the plasma Q -value attainable, it is desirable to have the largest possible η_I , in order to minimize the investment in beam injectors, as well as the quantity of electric power that must be circulated. The values of η_I shown in Table 4 are similar to those stipulated in conceptual design studies (Sect. 9). While CIT and TCT neutron sources might be used for similar purposes, the much higher injection power typical of TCT systems demands a larger η_I for economic viability. The large value of η_I in the beam-driven thermonuclear case reflects the need to maintain large overall power multiplication; a smaller η_I can be tolerated in the ignited-plasma case,

where the injectors are used during only a small fraction of the total burn cycle.

The injectors should be capable of continuous operation. The neutral-beam composition should be at least 90% monatomic, since fractional-energy ions generally contribute little to fusion production or heating of the central plasma. These ions also tend to enhance the problem of wall sputtering, since their energy is deposited in the outer plasma region [57].

Figure 28(a) shows schematically the components of an ideal neutral-beam injector system, which has $\eta_I \geq 0.9$:

- (1) A stable ion beam of current density at least 0.2 A/cm^2 is extracted from an ion source with excellent "gas efficiency" $\eta_g = \text{ion current/atom equiv.-current}$; e.g., $\eta_g > 0.50$.
- (2) Only monatomic ions are extracted.
- (3) The power dissipation in the ion source is negligible compared with the acceleration power at $W \geq 20 \text{ keV}$ (e.g., $< 2 \text{ kW/ampere-extracted}$).
- (4) The ion source, extraction electrodes, and acceleration electrodes can be operated steady state.
- (5) The neutralization efficiency η_n is nearly 100%.
- (6) No extracted ions are neutralized or otherwise lost before reaching the neutralizer, and no energetic neutrals are ionized or otherwise lost during transit in the beam duct.

Positive-ion sources have been developed that meet requirements (1) and (3) [142-145]. The steady-state requirement (4) seems readily attainable with some modification of these same sources, while requirement (6) can be approached with careful design. Sources that satisfy requirement (2), while also satisfying (1), (3), and (4), are not yet known, so that molecular ions must be separated out in practical injectors [Fig. 28(b)]. The most serious drawback of positive-ion injectors, however, is the rapid decrease in η_n for

$W_b \geq 60$ keV (D^0), as shown in Fig. 29. At large W_b , direct recovery of the energy of unneutralized ions is required to achieve acceptable η_I [146,147]. While negative ions of any energy can be neutralized with very high η_n , no negative-ion source is known that satisfies requirement (1) or (3). Nevertheless, for energies above about 200 keV, negative-ion beams would seem to be essential for achieving $\eta_I \geq 0.5$.

Two additional methods that have been suggested for improving the efficiency of injectors utilizing positive-ion beams are recirculation of unneutralized ions through the neutralizer [148], and deflection of the unneutralized ions into a second neutralizer (followed by a second beam duct) [149]. However, direct energy recovery techniques would seem to be the most promising near-term method for improving injector efficiency in the energy range 60 to 200 keV. In particular, the in-line direct conversion system with the ion source grounded, shown schematically in Fig. 30(a), has performed well in initial tests [150]. This technique is relatively inexpensive, since no provision need be made for bulky bending magnets. Figure 30(b) shows η_I -values obtainable with direct recovery methods [151], when the ion source composition is 75% D^+ and 25% D_2^+ (as in the present Berkeley source [142]). If the ion source composition could be increased to 90% D^+ , then the overall injector efficiency would be about 65% for 150-keV beams, and 50% for 200-keV beams.

7.2 PERFORMANCE OF PRESENT-DAY BEAM INJECTORS

Over the past several years, neutral-beam injectors of improved performance have been pressed into service almost as fast as developed. Indeed, the major injector development programs are being pursued with application to specific confinement devices in mind. Smaller research programs in areas

such as direct energy recovery and the production of negative ion beams have more long-range objectives.

The characteristics of some of the most advanced beam injectors are shown in Table 5. The high-current positive-ion beams are extracted from the plasma source by means of a grid with multiple circular or linear apertures. On leaving the source the dense accelerated ion population is space-charged neutralized by cold electrons. In present injectors the accelerated ions are charge-exchanged by the gas that escapes from the source, thereby producing the energetic neutral-atom beam. The achieved gas efficiency and source power dissipation are acceptable. For applications in the 40 to 80-keV range, the most urgent requirement is to enhance the proportion of monatomic ions in the extracted beam, a task that may require the development of a new ion source. Another approach is to operate the walls of present sources at a temperature of order 1500°C. For applications in the 80 to 200-keV range, high efficiency demands also the perfection of direct electrical energy recovery techniques [146,147]. An increase in neutral-beam current above that shown would be desirable for reactor applications (e.g., 100 A-equiv./source), but in any event a given beam line can be fed from several sources. Multi-source operation has in fact been successfully implemented in both the TFR [152] and 2XIIB [153] devices (using perpendicular injection in both cases). In the latter device, a mirror machine, up to 260 A from a total of 12 separate sources at 15 to 19 keV (4.4 MW) was injected into a target plasma of only 14-cm diameter and 100-cm length; these requirements are probably more rigorous than those demanded by any beam-injected tokamak.

Negative-ion sources are presently in an early stage of development. The Brookhaven negative-ion source [154] summarized in Table 5 is similar to the Novosibirsk source [155]. It is not clear whether direct negative-ion extraction from magnetrons or similar tubes will ever be practical for the

generation of high-current beams. The most promising technique may turn out to be the conversion of low-energy (~ 2 keV) positive-ion beams, by double charge-exchange in a sodium or magnesium gas cell [156]. This conversion process is only 10-20% efficient, so that about 20 kW/A must be invested to produce the negative-ion beams (but direct energy recovery could also be used at this stage). For $W_b \geq 200$ keV, overall efficiency would be greater than 50%, especially if a plasma cell is used for beam neutralization [146]. A serious unresolved problem in the transport of high-current D^- beams is the need to supply a space-charge neutralizing population of positive ions while minimizing the population of electrons in the beam region.

7.3 BEAM LINES

Present-day injectors are operated as if they were ideal [Fig. 28(a)]; they employ no species separation, no direct recovery, and only limited pumping (if any) along the beam ducts. The tightness of the present injector layout is advantageous in that the neutralizer can utilize gas flowing from the ion source, thereby alleviating the overall pumping problem. In most experiments, the gas loads have been manageable during the relatively short injection period of 10 to 150 ms. If the molecular components (D_2^+ and D_3^+) were to be separated out of the beam, however, elaborate pumping would be required in front of the neutralizer, and a separate gas source would probably be required to serve the neutralizer.

Figure 31 shows an elevation view of a proposed beam line for the TFTR device [157]. Even this rather simplified design — there is no separation of fractional-energy components and no electrical energy recovery — has a total length of 8 m, and dwarfs the tokamak vacuum vessel. The principal distinctions between the system illustrated and present-day practice are the

provision for direct cryopumping of the beam line, and deflection of the un-neutralized ion beam.

Achieving adequate pumping of the separator regions and the beam duct leading to the torus may, in the end, be the most serious obstacle facing the practical implementation of efficient quasi-steady-state neutral-beam injection. There are two potential problems associated with inadequate pumping in the beam ducts. (1) Ionization of energetic neutrals by impact with background gas; the consequent wall bombardment leads to gas release and further ionization [141]. (2) Flow of cold gas into the torus. The latter effect may be useful for raising the plasma density in some types of operation, but is undesirable for applications such as CIT reactors. The first experimental tokamak where the beam ducts have sufficient diameter (25 cm) to maximize beam transmission ($\eta_{t2} \geq 0.8$) is the DITE device at Culham [158]. The DITE beam lines are also equipped with titanium getter pumping. The first beam lines to be equipped with a cryopumping capability prototypical of reactor systems are those being fabricated at Oak Ridge for the PLT device at Princeton [159].

In a tokamak reactor, the cross sectional area of the beam ducts will be sufficiently large ($\sim 0.3\text{-}1\text{ m}^2$) so that they can also serve as the main torus pumping system. A serious potential problem is the streaming of neutrons up the beam lines, which results in a gradual degradation of insulator performance in the ion source and direct-recovery electrodes. However, the periodic replacement of electrode structures seems quite feasible, and in any event the electrodes could be protected by the use of bending magnets. A more serious concern is neutron penetration of the walls of the beam lines in the region of the superconducting toroidal-field magnets. Additional shielding is required around the beam lines in order to protect these coils.

The total cost of neutral-beam injectors designed for the energy range $w_b = 40$ to 150 keV is variously estimated at \$500 to \$1500 per kilowatt of neutral beam delivered into the plasma. At least three-quarters of this cost is due to the power supplies. The total cost of a powerful injector system for a large beam-driven tokamak can be comparable with the cost of the tokamak itself (although it would account for only a fraction of the total project cost including buildings, tritium handling, heat conversion, etc.). That roughly equal outlays for injectors and tokamak are required is not particularly dismaying, since one system is of little practical use without the other. The dominant power supply cost can perhaps be minimized by using direct rectification of the electric grid. However, it seems that significant reduction in injector costs can be obtained only by maximizing the beam production efficiency with the use of high-quality monatomic-ion sources and efficient direct energy recovery systems.

8. EXPERIMENTAL RESULTS

8.1 FEASIBILITY DEMONSTRATIONS REQUIRED

Table 6 summarizes the plasma performance that would be regarded as a favorable feasibility demonstration for the plasma component of each type of beam-driven tokamak reactor. These conditions must be met for quasi-stationary operation, that is, for an experimental time $\Delta t \gg \tau_s$. It is assumed that injector performance and neutral-beam penetration are satisfactory. (Except for the plasma current entry, the conditions listed in Table 6 are applicable to all toroidal configurations.)

The demands on energetic-ion performance increase with decreasing demands on bulk-plasma performance. The requirements on the CIT ion population are the most severe — although the required T_e and $n_e \tau_E$ are the most modest. The most critical feature of CIT operation is that plasma recycling must be sufficiently limited so that the influx of cold plasma and neutrals into the hot reacting region is small compared with the rate of particle influx from the neutral beams. Optimal performance of both TCT and CIT reactors demands that the fast-ion slowing-down rate must be near-classical, charge-exchange losses must be small, and plasma equilibrium must be maintained when the energetic-ion pressure is very large. The fueling and plasma-current control provided by the beams can be regarded perhaps as bonuses, although they would be necessary for steady-state operation. For the thermonuclear regime, the fast ions need only supply heat to the bulk plasma, and their slowing-down rate need not be classical. Ignition is feasible only for very large $n_e \tau_E$ and very small Z_{eff} .

8.2 PLASMA HEATING

The first neutral-beam injection experiment was performed on the CLEO tokamak (Culham) in late 1972, but the beam power was too small to produce

measurable heating [160]. In 1973 ion heating was observed in the ATC device (Princeton) [161], the ORMAK device (Oak Ridge) [162], as well as in CLEO [163] with injected powers of 40-80 kW at $W_b = 14-30$ keV. By early 1976 T_i had been doubled in experiments on the ATC [84], ORMAK [164], and TFR (Fontenay) [165] devices utilizing 200-400 kW injected power.

Table 7 shows the most impressive experimental parameters that have been reported as of May 1976. The energetic-ion density has been increased to produce $\langle \Gamma \rangle$ up to 0.6 (spatially averaged value) with no adverse effect on the plasma equilibrium. (Early ORMAK experiments indicated some equilibrium problems during counter-injection, when the plasma current was not sufficiently large to retain most of the orbits of counter-injected ions [162].) While injection tangential to the magnetic axis is generally used, in TFR near-perpendicular injection has been employed, also with apparently no adverse effects, although oscillations at harmonics of the ion cyclotron frequency have been observed in the outer region of the plasma [165]. In these experiments the bulk-ion temperature T_i is measured principally by the standard charge-exchange neutral method. In the case of a deuterium bulk plasma with H^0 or He^0 injection, T_i has also been determined from the neutron production rate. In ORMAK [164] and TFR [165], T_i has been raised from about 800 eV to record values of 1.5 keV or more. In the ATC experiment shown in Fig. 32(a), T_i was increased from 200 eV to 400 eV in the precompression plasma, and then to 1200 eV by a major-radius compression $C = 2.2$ [84]. The 70% increase in n_e during the beam heating resulted finally in a factor of 50 increase in the bulk-ion energy density. In these three devices beam injection has produced $\langle T_i \rangle \geq \langle T_e \rangle$, so that the large bulk-ion temperature is sustained wholly by the injected beams. The total energy of the warm ions and energetic ions (taking into account impurity depletion of the hydrogenic species) exceeds the electron energy by as much as 50%; i.e., $\Pi_i = 1.5$.

The observed ion temperature increases are consistent with calculations that include energetic-ion losses by charge-exchange on background neutrals, and by orbits that intercept limiters. The latter effect has been especially pronounced for counter-injected ions in ORMAK [162], while in ATC up to 60% of the injected power is lost by charge exchange. As shown in Fig. 32(b), the bulk-ion temperature rise tends to saturate in these experiments after 10 to 15 ms (depending on the device), which provides a measure of τ_{Ei} . In all experiments reported thus far, the increase in T_i has been proportional to the injected power, indicating no degradation in τ_{Ei} . The observed τ_{Ei} is often consistent with neoclassical ion heat transport coupled with charge-exchange losses [164,166]. (In addition to the four experiments discussed herein, a recent injection experiment on the T-11 device at Kurchatov has resulted in $\Delta T_i \sim 150$ eV for injected power ~ 150 kW [167].)

Measurements of T_i from the bulk-plasma neutron yield ($H^0 \rightarrow D^+$) are generally in agreement with the charge-exchange determination [161,162,168]. When the neutron-deduced temperature was too small, the discrepancy has been attributed to displacement of the deuteron population by hydrogen gas flowing from the injector beam lines [165], although this explanation is still only a qualitative one. In those cases where the neutron-deduced T_i was too high, the discrepancy has been attributed to the absence of a large neutral population in the hot central region [164]: That is, the neutron flux gives T_i near the plasma center for a spatially peaked temperature profile, but in a dense plasma the observed charge-exchange spectrum is representative of a region somewhat displaced from the center. This latter phenomenon will become increasingly troublesome in larger, more opaque plasma experiments, and will necessitate greater reliance on neutron diagnostics and Thomson scattering to determine T_i .

In all experiments to date the injected power has been comparable with or less than the Ohmic heating power. Because $W_b \sim W_c$ most of the injected energy has been transferred to the bulk ions. Only in TFR has there been definite evidence for a significant increase in central T_e [165] ($\Delta T_e \sim 200-300$ eV), but this result may reflect mainly a reduction in power drain from electrons to ions, because of the increase in T_i during injection. While radial profiles of T_e have been broadened in ATC [84] and ORMAK [169], the inability to noticeably raise the central T_e has been a somewhat disconcerting feature of most of the higher-powered experiments. Relevant hypotheses are that T_e is "clamped" by excitation radiation from a stubborn high-Z impurity population, or that the electron thermal conductivity has the steeply falling temperature dependence characteristic of diffusion caused by the trapped-electron instability [103] ($\tau_{Ee} \propto T_e^{-7/2}$). A definite increase in the level of high-Z impurities during intense beam injection has been observed in some experiments [165,170].

On the other hand, in ORMAK and ATC with simultaneous beam injection and gas influx it has been found possible to raise n_e to a value nearly twice that attainable with ohmic heating and gas injection alone; thus the electron energy density has been nearly doubled. Apparently cold electrons in the outer region of the plasma are being heated directly by the energetic ions, or the additional energy deposited in the central region flows quickly to the cold-electron outer region. In either event the customary shrinkage of the current channel during gas influx [12,83] can be warded off. It is relevant to note that neutral-beam injection experiments on the FM-1 toroidal multipole device at Princeton have succeeded in raising T_e from 10 eV to 40 eV, with a concomitant factor of 5 increase in n_e by ionization of background gas [171]. These results provide support for schemes to build up large plasma densities in toroidal containment devices by means of beam injection alone [99].

The "first generation" heating experiments on several different tokamaks have produced generally consistent results. A rule of thumb is that the central ion temperature increases as 1-2 eV per injected kilowatt. The plasma heating rates (20-50 eV/ms) themselves are actually quite suitable for reactor purposes; the principal ingredient needed for attainment of reactor-like temperatures is longer energy confinement time. The beta-values achieved with beam heating have already become significant. In ORMAK, $\langle \beta_p \rangle \approx 1.5 \approx 0.4 R/a_p$, while $\langle \beta_t \rangle \sim 0.5\%$ and $\beta_t(0) \sim 1.5\%$ [164]. Similar values are obtained in the injection-plus-compression ATC plasma [84]. These results indicate that achieving the nominal MHD-limiting value of $\langle \beta_p \rangle \sim R/a_p$ is feasible with increased beam power, or larger τ_E .

8.3 SLOWING-DOWN RATE AND NEUTRON PRODUCTION

The slowing-down process of the energetic ions is found to be essentially classical at the hot-ion fractional pressures achieved thus far ($\langle \Gamma \rangle \approx 0.5$). The only evidence for velocity-space instability has been the ion cyclotron modes observed during perpendicular injection on TFR [165], but there is no indication of anomalous slowing-down. The processes of slowing down and angular scattering have been observed by the charge-exchange neutral spectrum and by the D-D neutron production rate. Figure 33(a) gives an example of a charge-exchange neutral spectrum measured parallel and perpendicular to the magnetic axis in ORMAK [162]. Figure 33(b) shows a comparison of the measured and calculated neutral spectrum in ATC [23], where the neutral density is rather high ($\sim 5 \times 10^8 \text{ cm}^{-3}$ in the central region). The substantial charge-exchange loss in ATC is evident by the rather steep fall-off of the distribution at energies below the injection value (14.2 keV). Classical slowing-down spectra have also been obtained on the CLEO [163] and TFR [165] devices. The

slowing-down process has also been found to be classical in low-temperature toroidal multipole experiments [172,173], where the dominant frictional force is exerted by plasma electrons.

Recent detailed measurements of $Z_{\text{eff}}(r)$ from the rate of change of the charge-exchange neutral spectrum at various angles to the direction of an incident "pencil beam" are in good agreement with gross values of Z_{eff} deduced from Thomson scattering measurements [174]. The pencil-beam experiment gives for the first time rather detailed information on $Z_{\text{eff}}(r)$ and $q(r)$, and demonstrates that the angular scattering rate is classical.

For pulsed D^0 beams injected into deuterium plasmas, the time-resolved neutron yield provides an additional measurement of the slowing-down process. The measured neutron yields agree with the calculated values to within a factor of 2 [161,165,175]. An uncertainty of this magnitude is expected because the fast-ion orbits cover significant variations in plasma temperature and density, and because of uncertainties in the neutron measurement itself. In calculating the neutron yield, it is important to take into account depletion of the target-ion population by impurity ions, and also to include the fusion reaction rate among the fast ions themselves [176].

Figure 34 shows neutron-decay measurements from early ATC experiments with 40-kW injection [175]. While the neutron decay rate for parallel injection is 3-4 ms, for antiparallel injection the decay time is of order 1 ms [168]. These decay times are accounted for by the fast-ion slowing-down rates observed by charge-exchange measurements, which in turn are consistent with the classical ion slowing-down rate taking into account the predicted accelerating (or decelerating) effect of the toroidal electric field [161]. (The electric-field effect results in significantly less plasma heating by counter-injected ions, even if well confined [169].) When the plasma is compressed by $C = 2.2$ immediately following injection, the neutron production

rate increases by 2 orders of magnitude — because of the increase in $n_e^2 a_p^2 R$ $\propto C^2$, the increase in T_i (which increases the beam-target reactivity), and especially because of the compressional energization of the fast ions. For tangential beam injection, the expected $W_b \propto C^2$ has been measured [161,174]. The observed fast-ion energization by adiabatic compression and by the toroidal electric field (for co-injected ions) support the feasibility of energy clamping in two-component reactors (see Sect. 6.3.3).

In the range of injection energies currently used ($W_b = 10\text{--}30$ keV), the D-D fusion reactivity rises steeply with relative collision velocity v_{12} . For example, $\sigma(v_{12})v_{12}$ increases by a factor of 700 when v_{12} increases from 1×10^8 cm/s (10 keV) to 2×10^8 cm/s, corresponding to a head-on collision between two 10-keV ions. With multi-angle D^0 injection or with plasma $Z_{\text{eff}} \gg 1$, large-relative-velocity collisions between energetic D^+ can be the dominant source of neutrons when $n_h/n_e \geq 1/10$ [176], corresponding to injection powers above about 200 kW in small devices. This effect provides a partial explanation of the intense neutron production often observed when D^0 is injected into a target plasma that is nominally hydrogen [177]. The observations have little relevance to the feasibility of CIT operation, however, because the dominant Coulomb interaction in present experiments is still between beam and background ions.

8.4 EQUILIBRIUM EFFECTS

Perhaps the most surprising result to date has been the very small observed toroidal plasma rotation, even in the presence of unbalanced injection from a high-power single injector. This lack of substantial rotation was reported both in CLEO [163] and in ATC [161] experiments. Neither impurity-line Doppler shifts nor the rotation frequencies of MHD perturbations have indicated a toroidal rotation velocity greater than about 5×10^5 cm/s,

in cases where at least several times larger velocity would have been expected from the momentum input [73,75]. Some evidence of a substantial toroidal flow velocity for co-injection has been observed in ORMAK [162], but it is still not clear what part of this measurement may be due to effects such as "runaway" ions or large banana orbits [178]. Since neither charge-exchange loss nor the actual toroidal-field ripple can account for the observed lack of rotation in CLEO and ATC, there appears to be some anomalous damping mechanism at work. Presumably this mechanism would also mitigate the undesirable effects of the radial electric field that could be set up by perpendicular injection, particularly with use of the ripple-trapping injection technique.

In early ORMAK experiments, counter-injection tended to degrade the plasma equilibrium: MHD oscillations increased, the plasma position shifted considerably, the central T_e decreased, more impurities entered the plasma [162,169]. These undesirable effects were thought to be due to the large fraction of counter-injected ions that entered the velocity-space "loss region" (cf. Fig. 12), because of enhanced angular scattering by impurities ($Z_{eff} \geq 5$), and because of the relatively small plasma current ($I_p \sim 100$ kA). On the other hand, co-injection had the opposite effect on the plasma equilibrium, and was generally benign. The adverse effects of counter-injection largely disappeared with an increase in I_p to about 170 kA. Experiments with near-perpendicular counter-injection on ATC showed inefficient heating, also resulting from the expected rapid transport of the fast ions into the velocity-space loss region (as well as direct injection into loss orbits), but there were no adverse effects on the plasma equilibrium [168]. These experiments with counter-injected ions have done much to verify the theory of fast-ion orbits in tokamaks.

In tokamak experiments thus far there has been no direct evidence of a current perturbation by the beams themselves and none has been expected, because of the relatively small beam currents, and because of the pitch-angle scattering that occurs during "stacking" (see Sect. 4.4.1). But in a recent injection experiment in a toroidal multiple device, a plasma current apparently driven by injected ions has been observed [173]. In tokamak experiments the current density profile has no doubt been affected by the beam-induced changes in the T_e profile. Direct evidence of a beam-induced shift in the location of the $q = 1$ surface has been obtained in ATC by observation of changes in the relaxation oscillations set up by the internal kink mode [179].

8.5 FUTURE EXPERIMENTS

Experiments to date with $\langle T \rangle$ up to 0.6 have demonstrated the expected heating of bulk ions, the fast-ion slowing-down rate is found to be classical, and the expected beam-target neutron production is observed. Fast-ion orbits can be well-confined with reasonable plasma currents, and in some devices charge-exchange losses have been made negligible. Two experimental results unexpected from previous theoretical considerations are pleasant surprises — the small observed toroidal plasma momentum, and the large ramp-up in plasma density that is made possible with intense beam injection. For tokamaks of low current density, neutral-beam injection is as necessary for raising density as it is for raising temperature. It appears that neutral-beam heating will become as much a staple feature of tokamak operation as Ohmic heating itself.

Nevertheless, upon examining the required feasibility demonstrations listed in Table 6, it is evident that many of the crucial tests lie ahead. In particular, injection at sufficiently large W_b and with sufficient power

to raise T_e (i.e., $P_b \gg P_{OH}$) must yet be demonstrated. The injected power must be raised by a large factor to drive an appreciable fraction of the plasma current, or to provide a significant fraction of plasma fueling. Finally, the plasma impurity level in most present injection experiments is intolerably high.

Large increases in injection power, and hopefully larger $n_e \tau_E$, will be possible in a new generation of beam-injection experiments, making use of tokamaks recently commissioned or presently under construction. Table 8 lists important parameters of these devices, all of which will be equipped with high-power beam injectors. In most cases the injection pulses will be sufficiently long (~ 300 ms) so that plasma conditions will reach a steady state. The feasibility of TCT operation should be largely settled with experiments in PLT and ORMAK Upgrade. If successful, these experiments would imply also the feasibility of beam-driven thermonuclear operation, which requires chiefly an increase in $n_e \tau_E$. A ripple-trapping injection experiment is planned for PLT and possibly for an upgraded version of TFR.

The total ion energy (including fast ions) in present injection experiments is already comparable with or even exceeds the electron energy (i.e., $\Pi_i \geq 1$). In the new devices W_b will be sufficiently low so that at least half the beam energy will be transferred to plasma ions. Therefore a continuation of the present trend toward higher Π_i seems likely, and with it enhanced interest in hot-ion, warm-electron reactor plasmas such as the CIT. The two divertor experiments (DITE, PDX) are especially relevant for testing the heavily driven CIT and TCT operating modes, since an "unload" divertor should be effective in the rapid pumping of ions that diffuse out of the discharge, thereby reducing the rate of recycling. PDX will also be equipped with elaborate getter-pumping surfaces housed in large compartments above and below the discharge region [180], so that even without divertor operation, the influx of cold particles into the discharge should be reduced greatly.

9. CONCEPTUAL DESIGNS OF BEAM-DRIVEN REACTORS

The enormous cost of a thermonuclear tokamak reactor operating at ignition might preclude construction of such devices for commercial application in the near future — even following a successful laboratory demonstration of ignition conditions. The large fusion reactivity obtainable in relatively small beam-driven tokamak devices, on the other hand, does provide an opportunity for implementing tokamak reactors of manageable size and investment, for applications where very large Q -values are not essential. Such reactors would have fusion power outputs in the range 10 to 300 MW. Recently a number of conceptual designs of heavily beam-driven reactors have appeared. These devices are intended for specialized purposes, such as "break-even" demonstrations, the testing of fusion-engineering components and special materials, or fissile fuel breeding.

9.1 PROTOTYPICAL REACTORS

A number of large tokamaks proposed for completion in the period 1981-1984 [181] are intended to bridge the gap between present plasma physics experiments (including those of Table 8) and full-fledged reactors. The most advanced designs are those of the TFTR (Tokamak Fusion Test Reactor) in the U.S.A. [124,182] and the JET (Joint European Torus) [183]. The TFTR has been approved (and initial funding provided) for completion by 1981, but the fate of JET is uncertain as of this writing. The TFTR D-T plasma is expected to attain "break-even" ($Q = 1$) conditions and reactor-like fusion power densities by means of the TCT approach, with the injection of 20 MW of 120-keV D^0 beams. Important objectives of both devices are to investigate the equilibrium and transport properties of high-temperature (5-10 keV), high- β_p tokamak plasmas, and to provide experience in the use of prototypical fusion reactor systems

such as tritium handling and remote handling of tokamak components. Both machines are to be equipped with a modest major-radius compression capability ($C \approx 1.5$) in order to augment the beam heating, or to enhance Q by energy clamping (see Sect. 6.3.3). Figure 35 shows a drawing of the TFTR device and Table 9 gives its principal parameters.

Two other proposed very large tokamaks are the JT-60 in Japan, and the T-20 in the U.S.S.R. It appears that construction of the JT-60 device will proceed with completion in 1981. This machine has parameters similar to those of the TFTR, but will not operate with tritium. The proposed T-20 device is intended to serve as a prototypical fusion-fission (hybrid) reactor, and accordingly will be equipped with blanket modules for studies of neutron radiation damage, neutronics, and coolant effectiveness. It is expected that this device will be completed shortly after 1982. The design parameters of all these machines as of July 1975 are summarized in Ref. [181]. In each case the use of copper toroidal-field coils dictates modest duty factors ($\sim 1\%$).

9.2 MATERIALS AND ENGINEERING TEST REACTORS

Interesting practical applications that could follow successful TCT or CIT operation in the TFTR and other devices include an "engineering test reactor" to serve as a test bed for fusion reactor technologies, and a "materials test reactor" intended to produce intense 14-MeV neutron radiation on sample wall materials and blankets. The former device need have neither very high neutron flux nor a duty factor approaching unity. On the other hand the materials test reactor must deliver a high neutron fluence, which in turn implies large neutron flux and large duty factor. For this application the most important plasma consideration tends to be maximization of P_f , in order to minimize the device size (i.e., cost) for a given neutron wall loading ϕ_w .

At the same time economic considerations demand that $Q \geq 1/2$, since the investment in neutral-beam injectors, and the amount of electrical power and tritium that must be circulated, are inversely proportional to Q . Finally, the essential large duty factor dictates a minimum burn time per cycle of several tens of seconds.

Given these general constraints, practical TCT reactors tend to have similar plasma parameters, such as those listed in Table 1. Table 9 compares the TFTR parameters with those specified in an illustrative conceptual design of a TCT engineering/materials test reactor [184]. This device would take advantage of the 1 MW/m^2 14-MeV neutron current that can be obtained over a useful test area of many tens of m^2 , in order to investigate bulk-radiation effects, as well as the heat-transfer performance of a large variety of wall materials and blankets. The device would have modular construction, and could also serve as a test-bed for fusion reactor systems such as tritium breeding and processing. No electrical energy conversion would be undertaken; the 300 MWe required by the beam injectors would be taken directly from the power grid. This power drain could be reduced by operation at a larger plasma temperature to give larger Q , provided that n_e can be maintained and the increased $n_e \tau_E$ requirement can be achieved. This approach is currently being followed in a more detailed design of a TCT materials test reactor [185].

Small engineering test facilities have been proposed using a CIT plasma as the fusion driver [186,187]. The CIT operating mode offers the most powerful means of generating a large fusion-neutron flux in the smallest-sized tokamaks with modest confinement parameters. Table 9 gives key parameters for the conceptual design of Ref. [187], which is a more detailed version of that described in Ref. [186]. While the neutron current of $6 \times 10^{12} \text{ n/cm}^2/\text{s}$ is only a fraction of that attainable in TCT devices, it is quite adequate for a fusion-engineering test facility. The neutron flux is also suitable for large-area

surface radiation testing and for elaborate neutronics experiments. The specified 60-keV beam injectors are essentially state-of-the-art, the TCT problem of target-ion fueling is bypassed, and the power drain is small. It should be emphasized, however, that the crucial problems of high-throughput pumping and the minimization of plasma recycling have been addressed only casually in this design.

9.3 FUSION-FISSION HYBRID REACTORS

Another potential area of application of beam-driven reactors comprises devices that would contribute to various facets of the nuclear fuel cycle in a fission power economy [188-190]: the breeding of fissile material, power production from a fissionable blanket, and the burning of long-lived actinide (transuranic) wastes. A fissionable blanket of natural or depleted uranium would allow a high- P_f , low- Q fusion driver to serve as a power generator, while producing plutonium for use in thermal fission reactors such as light-water-moderated reactors. In general, the most attractive features of such a hybrid reactor relative to the LMFBR (liquid-metal fast breeder reactor) are that the hybrid requires no start-up fissile inventory, and is capable of producing for sale at least 5 times as much plutonium per unit of thermal energy [190]. For satisfactory overall plant efficiency in stand-alone operation ($\eta_p \geq 0.30$), one requires $Q \geq 1$ for the fusion driver and a beam injection efficiency of 70%. Adequate production rates of plutonium and electrical energy demand a plant factor of at least 75%, which in turn requires a burn time per pulse of at least 100 s. The condition on large plant efficiency can be relaxed, however, if the hybrid is to serve primarily as a producer of make-up fissile fuel for a system of thermal fission reactors [188-191].

If tritium fueling can be accomplished by injection of pellets or liquid jets, then only D^0 beams are required, and $Q \sim 1$ is attained at $\langle T_e \rangle \approx 5$ keV and $n_e \tau_E \sim 10^{13} \text{ cm}^{-3} \text{ s}$ (Sect. 6.3.1). Table 10 gives the parameters of a TCT hybrid intended originally for burning actinide wastes [192], but which would also be suitable for use with a uranium blanket. An elevation view of the reactor cross section is shown in Fig. 36. With $\eta_I = 0.70$ and plant factor of 75%, this design utilizing a depleted-uranium, fast-fission blanket would have an average net electric power production of 950 MWe, and a net Pu-239 production of 9 kg/day [193] (the tritium breeding ratio is 1.05). The overall plant efficiency would be 30%. If the presently observed energy confinement scaling [89] $\tau_E \approx (3-4) \times 10^{-19} n_e q^{1/2} a_p^2$ continues to hold in non-Ohmic-heated plasmas, or decreases only moderately with T_e , then $n_e \tau_E \approx 10^{13} \text{ cm}^{-3} \text{ s}$ can be achieved readily in a device of this size.

If pellet injection is not available, one must inject both D and T high-energy neutral beams for fueling as well as for neutron production (see Sect. 6.3.4). In this case attaining $Q \geq 1$ requires operation at larger $n_e \tau_E$, with a higher T_e and a somewhat smaller neutron production rate (cf. Fig. 24). A beam-driven hybrid reactor operating in this mode is proposed in Ref. [194]. This design specifies $n_e \tau_E \approx 3 \times 10^{13} \text{ cm}^{-3} \text{ s}$ and $T_e \approx T_i \approx 12$ keV, so that approximately one-half of the fusion power comes from thermonuclear reactions. The burn time is expected to be 300 s, with about 15 s required in each cycle for exhaust, refill, start-up. When appropriate maintenance and blanket-fuel recycling are taken into account, the plant factor is approximately 80%.

As evident from Tables 9 and 10, most of the expected plasma parameters of the TFTR are already comparable with those of illustrative TCT reactor applications. The chief distinction between TFTR operation, and operation of an economically practical TCT neutron source, are pulse length and duty factor. The plasma engineering problems associated with obtaining pulse times

≥ 50 s include steady-state beam injection, particle exhaust, tritium fueling, and long-time impurity control. Most of the TCT conceptual designs specify the use of a poloidal magnetic divertor to provide the required high-throughput pumping capability (see Sect. 6.3.5).

The problems of steady-state operation can be alleviated somewhat in the CIT regime, as explained in Sect. 5.3. A CIT hybrid with a blanket designed to maximize the production of plutonium is reported in Ref. [195], and the device parameters are given in Table 10. In both CIT examples the Q -value reflects the nonideal effects discussed in Sect. 6.2, and the fusion-neutron power of 175 MW in the hybrid example is probably the largest one can reasonably design for. The blanket configuration in Ref. [195] was optimized to maximize the production of fissile fuel to be used in a system of LWR's (following the method of Ref. [191]). The net electrical output of the entire CIT/LWR system is 4300 MWe at an estimated delivery cost of 27 mills/kWh.

Whether any fusion-fission device can be economically competitive with other potential producers of fissile fuel, such as the LMFBR or spallation-neutron sources, remains to be seen. More detailed tokamak designs are required to determine accurately both capital cost and operating cost, as well as environmental acceptability. The optimum choices of blanket configuration and coolant are not yet clear. For another example, breeding of U-233 in a thorium blanket is a much less productive enterprise than Pu breeding [189-191], but may possibly be attractive if the U-233 is sold for use in fission reactors of very high conversion ratios, such as the HTGR and CANDU reactors. In this application, however, the fusion breeder is likely to be a small net consumer of electrical power.

In a variant of the hybrid reactor, it has been proposed to transform long-lived actinide waste products from fission reactors into shorter-lived radioactive waste products [188,189]. For this purpose, high-energy neutrons are most effective. Since the actinides are transformed primarily by fission

reactions, many tens of MeV are generated per fusion neutron. In the TCT actinide waste burner with parameters shown in Table 10 and Fig. 36 [192], the total thermal power production is about 15,000 MW. This study also finds [192] that to achieve desirable rates of actinide depletion, the 14-MeV neutron current should be at least 5 MW/m^2 ; this flux is realizable in principle in TCT operation (see Sect. 6.3.5), but only after the resolution of severe problems of neutral-beam access and particle exhaust.

The broadening of the fusion-neutron spectrum because of the large relative energy of the reactants in beam-driven reactors may have noticeable influence on the probability of neutron-multiplying reactions in the blanket [190]. Angle-averaged neutron energy spectra from TCT plasmas are reported in Ref. [196]. It is found that the width of the spectrum (FWHM) increases rapidly with W_b . At $W_b = 150 \text{ keV}$, the FWHM is 1.3 MeV for $D \rightarrow T$ and 0.65 MeV for $D \rightarrow D$.

9.4 LOW-Q "PURE FUSION" REACTORS

There have as yet appeared no conceptual designs of low-Q "stand-alone" tokamak power reactors, but it is interesting to define the minimum characteristics of such devices. The gross electrical power output is

$$P_{\text{out}} = \eta_c P_{\text{beam}} (1 + 0.8 M_Q + 0.2 Q) \quad (9.1)$$

where η_c is the thermal-to-electric conversion efficiency of the blanket and heat exchanger, and M is the energy multiplication factor of the fusion neutrons in the blanket. The electrical power required to drive the reactor is $P_{\text{in}} = P_{\text{beam}} (\eta_I^{-1} + f)$, where $f P_{\text{beam}}$ is the power consumed by non-injector systems such as blanket cooling or the charging of energy storage systems to be

used during current start-up. Typically, $f \approx 0.2$. The plant efficiency η_p is defined

$$\eta_p = \frac{(P_{out} - P_{in})/P_{beam}}{(1 + 0.8 MQ + 0.2 Q)} = \eta_c \frac{(\eta_I^{-1} + f)}{1 + 0.8 MQ + 0.2 Q} \quad (9.2)$$

Here it is assumed that the thermal rejected energy from the injectors is totally lost. If there are no fissionable or exotic (e.g., beryllium) neutron-multiplying materials in the blanket, then M is limited to about 1.4. If the minimum acceptable plant efficiency for any purpose whatever is 20%, then small Q 's are permitted only for very large values of η_o and η_I . Taking $\eta_I = 0.70$ (which is feasible with 100-keV D^o beams, for example [197]) and $\eta_c = 0.40$, we find that $\eta_p = 0.20$ is possible only for $Q \geq 5.3$. Such Q -values are not attainable in TCT operation, except for pulsed energy-clamping under extremely idealized conditions. However, $Q \sim 5$ is possible in either CIT or beam-driven thermonuclear operation, but both modes require (ideally) $n_e \tau_E \approx 7 \times 10^{13} \text{ cm}^{-3} \text{ s}$ with $T_e \approx 15 \text{ keV}$. Such plasmas are likely to be very large.

The required Q -value for significant net electrical production decreases rapidly with increasing η_c . An attractive means of enhancing η_c is by means of a potassium-steam binary cycle [198], which should result in $\eta_c = 0.55$. Then $\eta_p = 0.20$ could be obtained with Q as small as 2.8. Recently it has been proposed to combine driven tokamak reactors with a hydrogen turbine, metal-hydride-compressor energy conversion system [199]. The thermal output of the reactor is used as the high-grade heat source. A low-grade thermal input, which could be the reject heat from an existing power plant, or a geothermal heat source, is required to operate the compressors. With this conversion system it seems possible to reach $\eta_c \approx 0.80$ (the low-grade heat input is not included), so that $\eta_p = 0.30$ could be obtained with Q as small as 1.7.

Low-Q "pure fusion" tokamak reactors must use superconducting toroidal-field coils, and superconducting OH and EF coils as well, in order to have acceptable η_p . But devices intended for the applications discussed in Sects. 9.2 and 9.3 can certainly employ normal-conductor pulsed-field coils (copper or aluminum). Some of the EF coils, including the divertor-field coils, can then be placed inside the TF coils, in order to afford better control over plasma shaping, and to reduce the required ampere-turns and associated magnetic forces. The power dissipation during stationary discharge conditions would be at most several tens of megawatts.

The reactor examples of Tables 9 and 10 have plasma betas of order 5%. If substantially larger values of beta eventually become feasible, or if q can be significantly lowered, then B_t can be reduced greatly. In that case it might be advantageous to employ normal-conductor TF coils for materials test reactors or fissile breeders whose main purpose is not electrical power production — especially if the substitution of normal coils for superconducting coils results in a substantial decrease in capital cost. However, for steady-state operation a further reduction in B_{max} (at the coil windings) and corresponding increase in a_p are required: It is estimated that for a 6-m bore, steady-state copper-coil operation is feasible with $B_{max} = 6$ T and a resistive power consumption of 300 MWe.

10. CONCLUSIONS AND OUTLOOK

The observed performance of injected energetic ions in experimental tokamak plasmas is in essential agreement with theoretical expectation. Ion deceleration is classical, the expected bulk-ion heating is observed, fast-ion orbits are understood, plasma equilibrium and stability are maintained. The injected power delivered to electrons, however, has not yet been sufficiently strong compared with the Ohmic power input to have a significant effect on T_e . Present results support the plasma physics basis of at least some types of beam-driven tokamak reactors, although many performance criteria must still be demonstrated in larger devices with greatly enhanced injected power.

The chief technological problem is probably the efficient production and transport of high-power, high-voltage, quasi-steady-state neutral beams. For injection energies up to about 200 keV (D^+), there appears to be no insuperable obstacle to achieving injector efficiency of at least 40%, including transport efficiency, even with present-day ion sources. Nevertheless, a great deal of development work lies ahead in the provision of adequate beam-injector pumping systems, and in providing ion sources of increased atomic-ion content.

The most serious technical problem in the area outside the injectors themselves is satisfactory penetration of the energetic neutral beams into dense plasmas that may contain appreciable impurity content. If the ripple-trapping injection technique proves successful, however, beam energies in the range 150 keV or less would seem to be adequate for all applications involving D-T plasmas. If beams of much higher energy must be used for special applications, negative-ion sources are required in order to maintain injector

efficiency, but as yet there can be no assurance of the successful development of high-current negative-ion beams.

Finally, CIT reactors employ injector systems that are practically state-of-the-art together with tokamaks of rather modest size, but the required plasma operation places severe limits on the allowed influx of cold plasma and neutrals into the reacting region.

In sum there appear to be no fundamental plasma-physical obstacles to the development of beam-driven tokamak fusion reactors, and the requirements on the injector systems for at least some applications can apparently be met by reasonable extrapolation of known technology. The principal plasma-engineering problems would seem to be those associated with achieving large fusion power densities in quasi-stationary operation, such as providing adequate means for maintaining a large particle throughput, impurity control, and fueling.

Perhaps the most appealing feature of neutral-beam-sustained tokamak plasmas is their capability for useful performance in devices of small size (and relatively small unit cost), a feature that follows directly from the unique role of the energetic ions in enhancing fusion reactivity. It seems to the present author that this special characteristic should be used to advantage in the gradual implementation of tokamak reactors, as follows. First, the construction of small low-power devices on the CIT model ($P_{\text{fusion}} \sim 10\text{-}50$ MW, $Q \sim 1$) to produce 14-MeV neutrons for specialized applications such as radiation testing, neutronics, and blanket studies. Second, larger TCT or CIT devices ($P_{\text{fusion}} \sim 100\text{-}300$ MW, $Q = 1\text{-}3$) producing large 14-MeV neutron fluxes for materials testing or fissile or tritium breeding. Next, somewhat larger beam-driven thermonuclear devices ($P_{\text{fusion}} \sim 500\text{-}1000$ MW, $Q = 5\text{-}10$) that could be useful as pure-fusion power sources for specialized purposes where low plant efficiency is tolerable. Finally, high- Q beam-driven or even

ignited thermonuclear reactors with exceedingly large power outputs ($P_{\text{fusion}} \geq 2000 \text{ MW}$) that could be used as base-load power sources. Following the development of such D-T reactors, one could then turn confidently to the design of beam-driven D-D reactors with their greater demands on beam energy and plasma confinement.

The gradual increase in reactor size as outlined above can be undertaken only with experience in device operation, and following on new ideas for optimization of engineering design and plasma performance. Although there is now little doubt that a number of test devices of certain of the types discussed herein will eventually be constructed, in the end the deployment of large numbers of such reactors for practical application must depend on their economic competitiveness, a subject that remains to be investigated.

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Table 1

BEAM-DRIVEN TOKAMAK REACTORS

	CIT	TCT	BDTN
Injection energy (keV)	50-100	100-250	≥ 300
Plasma current (MA)	0.5-4	1.5-5	4-12
$n_e \tau_E$ (cm^{-3}s)	$\geq 3 \times 10^{12}$	$\geq 10^{13}$	$\geq 5 \times 10^{13}$
T_e (keV)	3-10	5-15	8-20
(Energetic-ion pressure)/ (electron pressure)	$\gg 1$	≈ 2	$\ll 1$
Fusion power multiplication	0.5-3	0.5-2	> 3
Fusion power density (W/cm^3)	0.5-2	1-6	1-3

Table 2
Functions of the Energetic Ions in Various Reactor Regimes

	$\Gamma \ll 1$ (Beam-Driven Thermonuclear)	$\Gamma \sim 1$ (TCT)	$\Pi_i \gg 1$ (CIT)
Heat bulk plasma	✓	✓	(electrons)
Control T(r)	✓	✓	✓
Fuel plasma	-	partly	entirely
Assist plasma current	-	partly	entirely
Fusion reactions with warm ions	✓	✓	-
Fusion reactions among energetic ions	-	✓	✓

Table 3
Stability Criteria for Counterstreaming-Ion
Velocity Distributions

Mode	Sufficient Condition for Stability	Practical CIT Parameters
Electrostatic, ω not near ω_{ci}	$T_{\parallel} > 1/2 T_e$, $u_b \lesssim 10 (T_e/m_h)^{1/2}$	$T_{\parallel} \gtrsim 1.5 T_e$ $u_b < 5 (T_e/m_h)^{1/2}$
Electromagnetic, ω not near ω_{ci}	$u_b < 1/2 v_{Alf}$	$u_b \gtrsim 1/5 v_{Alf}$
Electrostatic ion cyclotron waves	$\Delta W/W_b > 0.8 T_e/T_i$	$\Delta W/W_b > 2 T_e/T_i$
Electromagnetic ion cyclotron waves	$u_b < 1/3 v_{Alf}$ or $\Delta W/W_b \gtrsim 1$	$\Delta W/W_b > 4/3$

Table 4

NEUTRAL-BEAM INJECTOR REQUIREMENTS

	CIT Neutron Source	TCT Neutron Source	Beam-Driven Thermonuclear Power Reactor	Heating to Ignition
$\langle n_e \rangle a_p$ (cm^{-2})	$2-5 \times 10^{15}$	$\sim 10^{16}$	$\sim 2 \times 10^{16}$	$\geq 3 \times 10^{16}$ (final plasma)
W_b (keV), D°	50-100	150-250	≥ 300	≥ 300
P_{beam} (MW)	10-100	50-400	50-100	100-300
I_{beam} (A-equiv.)	150-1000	300-2000	100-300	≥ 300
η_I	0.5	0.65	0.7	0.5
Pulse length	continuous	continuous	continuous	~ 2 s
Ion source	D^+	D^+ with electrical energy recovery	D^-	D^-

Table 5

INJECTOR STATE-OF-THE-ART

	Ideal System	Berkeley [142] D ⁺	ORNL [143] D ⁺	Brookhaven [152] D ⁻	Fontenay [145] D ⁺
Ion source		Multi-filament Discharge	Duo-pigatron	Magnetron	Peri-plasma-tron
J(A/cm ²) source	0.3	0.5	0.35	0.7	0.4
I(extracted A)	≥ 50	75	30	1	40
Monatomic Fraction	1.0	0.75	0.60	0.5	0.70
Gas Effic.	≥ 0.5	0.3	0.5	0.04	0.5
Pulse Length(s)	Steady-state	0.1	0.7	0.02	0.2
Source Dissipation (kW/A)	≤ 10	3	4	35	1.5
W _b (keV)	50 to 500	40	40	18	30
I _{beam} (A-equiv.)	50	50	15	--	--
η _n	1.0	0.8	0.8	--	--
η _I (D ^o)	> 0.9	0.5 (η _t = 0.8)	0.4 (η _t = 0.8)	--	--
η _t in recent expts.	1.0	≥ 0.5	≥ 0.5	--	--

η_n = neutralization efficiency

η_t = transmission efficiency

η_I = overall injector efficiency

Table 6

PLASMA PERFORMANCE TO DEMONSTRATE FEASIBILITY

	CIT	TCT	Beam-Driven Thermonuclear
Maintain $\langle T_e \rangle \sim \langle T_i \rangle$ with $P_b \gg P_{OH}$	$\geq 3 \text{ keV}$	$\geq 4 \text{ keV}$	$\geq 7 \text{ keV}$
$n_e \tau_E \text{ (cm}^{-3} \text{ s)}$	$\geq 3 \times 10^{12}$	$\geq 1 \times 10^{13}$	$\geq 5 \times 10^{13}$
Equilibrium at $\beta_p \geq 0.5 R/a_p$	$\Pi_i \gg 1$	$\Gamma \sim 1$	$\Gamma \ll 1$
$f_h \text{ (v)}$	shifted Maxwellian- like, $n_h \gg n_i$	classical slowing-down	-
Charge-exchange loss	$\approx 0.1 P_b$	$\approx 0.05 P_b$	$\approx 0.05 P_b$
Fast-ion confinement	$\tau_h \geq \tau_s \gg \tau_{Ee}$	$\tau_h \geq \tau_s$	$\tau_h \geq \tau_s$
Beam-induced current	$J_b \geq J_{OH}$	beams control $J(r)$	-
Beam fueling	all beam-fueled; recycling minimized	significant beam fueling	-
$Z_{eff} \text{ (iron)}$	≈ 2	≈ 3	< 3

Table 7

BEST EXPERIMENTAL PARAMETERS

	Best result as of May 1976	Device
Equilibrium -parallel inj. -perpend. inj.	$\langle \Gamma \rangle \sim 1/2$ $\langle \Gamma \rangle \sim 1/2$	ATC, ORMAK TFR
Fast-ion slowing- down rate	Classical at all Γ	CLEO ATC ORMAK TFR
$P_{\text{beam}}/P_{\text{OH}}$	~ 1.2	ATC
$\Delta T_i (r = 0)$	800 eV (max $T_i = 1.7$ keV)	ORMAK TFR
$\Delta T_e (r = 0)$	300 eV (max $T_e = 2.2$ keV)	TFR
n_e enhancement factor	≤ 2	ATC ORMAK
$\langle \beta_p \rangle$	~ 1.5	ATC ORMAK
Π_i	~ 1.5	ATC ORMAK TFR

Table 8

NEW BEAM-INJECTED TOKAMAKS

	DITE (Culham)	PLI (Princeton)	PDX (Princeton)	ORMAK Upgrade (Oak Ridge)	Doublet III (General Atomic)	T-10M (Kurchatov)
R_o (m)	1.17	1.32	1.40	0.92	1.4	2.3
a_p (m)	0.26	0.45	0.42	0.30	0.45	0.75
B_t (T)	2.8	5.0	2.5	4.0	2.6	3.5
I_p (kA) at $q_a = 3$	280	1250	≤ 800	650	≤ 4000	1500
W_b (keV)	30	40	40	40	80	40
P_{beam} (MW)	1.5	3	4	4	12	10
Special feature	bundle divertor		poloidal divertor		3:1 plasma elongation	superconduct- ing TF coils
Operation with intense injection	1977	1977	1978	1978	1979	1980

Table 9

PLASMA PARAMETERS FOR THE TFTR AND CONCEPTUAL ENGINEERING TEST REACTORS

	TFTR [124]	TCT Materials Test Reactor [184]	CIT Engin. Test Reactor [187]
R_o (m)	2.5	3.05	3.0
a_p (m)	0.85	0.55	0.55
Plasma elongation	1.0	2.0	1.5
Horiz. wall radius (m)	1.1	0.75	0.75
First wall area (m ²)	110	149	131
TF coil bore (m)	2.8	5.5×7.8	3.6×6.0
Magnet type	Copper	NbTi	NbTi
B_t on axis (T)	5.2	3.8	4.0
B_t at coil (T)	9.5	8.0	9.6
I_p (MA)	2.5	2.4	1.4
q at limiter	3.0	2.5	3.0
$\langle n_e \rangle$ (cm ⁻³)	7×10^{13}	1.2×10^{14}	4×10^{13}
$\langle T_e \rangle = \langle T_i \rangle$ (keV)	6.0	5.0	4.5
$\langle n_e \rangle \tau_{Ee}$ (cm ⁻³ s)	1.5×10^{13}	8×10^{12}	2.6×10^{12}
W_b (keV)	120 (D°)	200 (D°)	60 (D°), 90 (T°)
Beam power (MW)	20	177	30
Γ	0.60	0.85	$\Pi_i = 5$
β_p	2.0	3.5	3.6
Pulse duration (s)	4	50	≥ 44
Duty factor	0.013	0.80	≥ 0.70
<u>Power Production</u>			
Q	1.0	1.0	0.70
P_f (MW/m ³)	0.80	4.3	0.68
Neutron power (MW)	16	142	17
Neutron production (n/s)	7×10^{18}	6.4×10^{19}	7.5×10^{18}
Neutron wall loading (MW/m ² , uncollided 14-MeV)	0.15	0.95	0.13

Table 10

PLASMA PARAMETERS FOR TOKAMAK FUSION-FISSION REACTORS

	TCT HYBRID [192]	CIT FISSILE BREEDER [195]
R_o (m)	3.9	4.7
a_p (m)	0.9	1.0
Plasma elongation	1.5	1.6
Horiz. wall radius (m)	1.4	1.2
First wall area (m ²)	218	343
TF coil bore (m)	5.3×9.30	5.3×8.3
Magnet type	Nb ₃ Sn	Nb ₃ Sn
B_t on axis (T)	3.65	5.6
B_t at coil (T)	9.2	12
I_p (MA)	3.8	4.8
q at limiter	2.5	2.8
$\langle n_e \rangle$ (cm ⁻³)	1.0×10^{14}	7.0×10^{13}
$\langle T_e \rangle = \langle T_i \rangle$ (keV)	6.5	10 (electrons)
$\langle n_e \rangle \tau_{Ee}$ (cm ⁻³ s)	9×10^{12}	1.4×10^{13}
W_b (keV)	200 (D ^o)	100 (D ^o), 150 (T ^o)
Beam power (MW)	325	145
Γ	0.92	$\Pi_i = 5$
β_p	3.4	4.0
Pulse duration (s)	50	≥ 100
Duty factor	0.80	≥ 0.85
<u>Power Production</u>		
Q	1.25	1.5
P_f (MW/m ³)	3.4	1.35
Neutron power (MW)	320	175
Neutron production (n/s)	1.4×10^{20}	7.7×10^{19}
Neutron wall loading (MW/m ² , uncollided 14-MeV)	1.47	0.40

PRINCIPAL SYMBOLS

a_p	Plasma limiter radius
B_t	Toroidal magnetic field
B_p	Poloidal magnetic field
C	Linear compression factor
E_ϕ	Toroidal electric field
E_f	17.6 MeV
f_h	Hot-ion velocity distribution
f_i	Warm-ion velocity distribution
$H(r)$	Trapping profile of injected ions
I_b	Injected beam equivalent-current
I_p	Plasma current
J_b	Beam-induced current density
J_{OH}	Ohmic current density
m_h	Hot-ion mass
n_e	Electron density
n_{ec}	Axial electron density
n_h	Hot-ion density
n_i	Warm-ion density
n_o	Neutral density
P_α	Fusion alpha power density
P_b	Injection power density
P_f	Fusion power density
Q	Fusion energy multiplication factor
Q_b	Fusion energy multiplication by beam-target reactions
r	Plasma minor radius
R	Plasma major radius
R_o	Major radius of magnetic axis
S	Plasma circumference/ $2\pi a_p$
T_e	Electron temperature
T_i	Ion temperature
T_h	Temperature of displaced Maxwellian
u_b	Mean streaming velocity of hot-ion distribution
v	Hot-ion velocity
$v_\parallel$	Hot-ion velocity component parallel to B
v_A	Alfven velocity

v_b	Injection velocity
v_c	Hot-ion velocity where electron drag = bulk-ion drag
v_e	Electron thermal velocity
v_i	Warm-ion thermal velocity
W_b	Injection energy
W_c	Hot-ion energy where electron drag = bulk-ion drag
W_l	Transition energy between warm ions and hot ions
$\bar{W}_h$	Average energy of hot ions
$\bar{W}_i$	Average energy of warm ions
Z_{eff}	Effective ionic charge for multiple species
Z_h	Charge of hot ions
β_p	Average plasma pressure divided by poloidal field pressure at limiter
Γ	Hot-ion energy density/bulk-plasma energy density
Γ_α	Alpha-particle pressure/bulk-plasma pressure
δ	Peak-to-peak magnetic field ripple
δ_*	Ripple well depth in presence of rotat. transform
ζ	$v_{ }/v$
η_c	Thermal-to-electric conversion efficiency
η_I	Overall beam injector efficiency
η_n	Beam neutralization efficiency
η_p	Net electrical power/gross thermal power
η_t	Beam transmission efficiency
θ	Poloidal angle
λ_t	Neutral-beam mean-free-path
ϕ_w	Neutron wall loading (uncollided 14-MeV)
Π_i	Total ion energy density/electron energy density
σ_{cx}	Charge-exchange cross section
σ_I	Impact-ionization cross section
τ_{cx}	Charge-exchange time
τ_E	Average bulk-plasma energy confinement time
τ_{Ee}	Electron energy confinement time
τ_{Ei}	Warm-ion energy confinement time
τ_p	Particle confinement time
τ_s	Thermalization time of hot ions
τ_{se}	Slowing-down time (1/e) of hot ions on electrons

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 Fig. 1. Schematic of ion velocity distribution functions parallel to the magnetic field in 3 principal types of beam-driven toroidal reactor plasmas. Ions are injected at $v_{||} = v_{inj}$.

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Fig. 2. Schematic representation of various tokamak reactor regimes for plasmas with T_e and $T_i \geq 5$ keV.

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Fig. 3. Steady-state velocity distribution functions for energetic deuterons injected at 100 keV into a deuterium background plasma with $T_e = T_i = 3$ keV. Curves are labeled by direction with respect to magnetic axis. Injection velocity (v_b) parallel to magnetic axis.

763602

Fig. 4. Steady-state velocity distributions (a) parallel and (b) perpendicular to the magnetic axis for deuterons injected at 200 keV ($v_{||} = v_b$) into a tritium background plasma with $T_e = 10$ keV and $T_i = 17$ keV. Deuteron confinement time $\propto W^{3/2}$. (c) Contours of the ion speed distribution function. (From Ref. [25])

Fig. 5. Bulk-ion velocity distribution function for various distortion parameters $\rho \sim (n_h/n_i)(v_i/v_c)$. v_i is the ion thermal velocity. (From Ref. [33])

763709

Fig. 6. Steady-state ion velocity distribution functions in a CIT (a) parallel and (b) in a direction x perpendicular to the magnetic field. The velocity scale $(\text{keV})^{1/2}$ is for deuterons only. $W_b = 60$ keV (D°), 40 keV (T°). (From Ref. [14]) (c) Ion velocity distribution function for $T_e = 12$ keV, $W_b = 60$ keV (D°), 90 keV (T°).

Fig. 7. (a) Rate coefficients for charge-exchange, ionization by plasma ions, and electron ionization for a fast deuterium atom in a hydrogenic plasma. (b) Maximum plasma opacity for acceptable neutral-beam penetration by tangential injection. Trapping length λ_t is calculated at central density n_{ec} .

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Fig. 8. Neutral-beam injection geometry. (a) Plan view. (b) Elevation view, showing the limiter-intercepting outer boundary of co-injected ion orbits. Counter-injected-ion orbits are shifted inward in R.

763708
 Fig. 9. Spatial shape factor for fast ions formed from co-injected and counter-injected 25 keV beams in ORMAK. Center of beam is tangent to magnetic surface at $R = R_c < R_0$. Each curve is labelled by the percentage of beam ions trapped in the plasma. (From Ref. [52])

763556

Fig. 10. (a) Variation of magnetic field following a flux line. In the "symmetric trapping region," there are no local magnetic wells. (b) Vertical injection of energetic deuterium atoms into a tokamak plasma with large toroidal-field ripple below the horizontal midplane.

763524

Fig. 11. (a) Required injection energy (D^0) so that deuterons drift from the bottom of the torus to the horizontal mid-plane ($z = 0$). δ_0 = peak-to-peak ripple at $z = 0$. $R_0 = 6.0$ m, $B_t = 5.0$ T. (b) Guiding-center trajectory over the slowing-down time (39 ms) of a 120-keV ion. (From Ref. [62])

DEUTERONS

$$I_p = 70 \text{ kA}$$

$$J = J_c [1 - (r/a)^2]^3$$

ATC

743975

Fig. 12. Loss regions in velocity space for the ATC device. Axes are labeled in keV. At any radius r/a , particles with velocity coordinates lying above the solid line have orbits that intersect the limiter. (From Ref. [168])

Fig. 13. (a) Fraction of beam energy and momentum given to bulk ions. (From Ref. [19]). (b) Fraction of fusion alpha energy (3.5 MeV) given to bulk ions and electrons. (From Ref. [26])

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Fig. 14. Γ_α = (alpha pressure)/(bulk-plasma pressure) in an ignited D-T plasma. Γ_b = (minimum hot-ion pressure)/(bulk-plasma pressure) for fast deuterons injected to heat the plasma to ignition. (From Ref. [81])

Fig. 15. Fusion reactivities in D-T for (a) TCT plasma (Ref. [98]); (b) D and T shifted-Maxwellian distributions, each of temperature T_b , with relative displacement velocity v_r (Ref. [35]); (c) thermonuclear plasma.

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Fig. 16. Steady-state contours of the ion speed distribution function in a CIT. Deuterons are injected at 60 keV, $\theta = 4^\circ$, and tritons at 40 keV, $\theta = 176^\circ$. Ions decelerated to $W < 2T_e$ have a lifetime $\tau_p \approx \tau_{Ee}$. (From Ref. [14])

Fig. 17. Dependence of (a) T_e and (b) ideal Q on electron energy confinement parameter for steady-state CIT operation. D^0 and T^0 are injected at same energy. 17.6 MeV per reaction. Neither fusion-alpha heating nor charge-exchange loss is included. (From Ref. [186])

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Fig. 18. Production rate of 14-MeV neutrons from D-T CIT plasmas of circular cross section. D^0 and T^0 injected with same energy. $B_t = 4.0 \text{ T}$, $R_0/a_p = 5.0$, $q_t = 2.5$, $\beta_p \leq 3.3$, $\lambda_t \geq a_p$. (From Ref. [186])

763694
Fig. 19. Reduction of Q in a CIT by inward flux of warm ions. Hot-ion flux is supplied by neutral beams.

763707

Fig. 20. (a) Fusion energy multiplication factor Q_b as a function of deuterium injection energy for various values of T_e in a cold triton-target plasma ($T_i = 0$), assuming $n_T = n_e$ (i.e., beam density is neglected). 17.6 MeV per reaction. (b) Showing enhancement of Q_b due to finite triton temperature ($T_i = T_e$) and deuterium energy diffusion. Reactions of deuterons at $W < 2T_i$ are not included.

763733
 Fig. 21. (a) Maximum attainable fusion power density for 200-keV deuterons injected into a triton-target plasma at T_e , with total plasma pressure = 0.33 J/cm^3 , corresponding to $B_t = 5.0 \text{ T}$, $\beta_p = 3.5$, $q = 2.5$. Q_b = effective fusion power multiplication. Alpha-particle effects are neglected. (b) Values of n_e , τ_s , and τ_E for maximum P_f . (Adapted from Ref. [113])

743839

Fig. 22. Effect of impurity ions on Q_b and maximum attainable P_f for 200-keV deuterons injected into a 6-keV triton-target plasma. Single impurity species of charge Z_I and concentration Z_{eff} . P_{f0} and Q_{b0} are values for $Z_{eff} = 1$. (From Ref. [113])

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 Fig. 23. (a) Optimal injection energies for clamped (W_{cl}) and unclamped (W) deuteron beams in a 100% tritium target plasma. (b) Maximum energy gains (17.6 MeV per reaction) for clamped (Q_c) and unclamped (Q) beams. (From Ref. [30])

763706

Fig. 24. Fusion power multiplication for 50:50 D-T plasmas at temperature $T_e = T_i$ maintained by injection of energetic deuterons at energy W_b and tritons at energy $3/2 W_b$. The dashed curves are contours of $n_e T_e$. n_{hot} is the total density of energetic ions. (From Ref. [108])

Fig. 25. (a) Dependence of Q and power density P_f on $n_e \tau_E$ for an 8-keV D-T plasma heated by 200-keV D beams. The triton fraction in the bulk plasma is optimized for maximum Q . (b) Q vs P_f for beam-driven tokamak plasmas of fixed total pressure 0.34 J/cm³ and various temperatures $T_e = T_i$. The dashed curves are contours of constant $n_e \tau_E$.

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 Fig. 26. A D-T plasma is heated to final temperature T_f by deuteron beams injected at W_b . f_α is the fraction of the final energy content supplied by alpha particles produced in beam-plasma fusion reactions. (From Ref. [81])

Fig. 27. Fusion energy multiplication factor for TCT operation with some "advanced fuels." In the D-D case, T and 3He reaction products are not burned. (From Ref. [136])

(a) IDEAL

(b) PRACTICAL

-148-

763692
Fig. 28. Components of (a) ideal and (b) practical neutral-beam injector systems.

763705
 Fig. 29. Equilibrium efficiency of conversion of a charged beam to a neutral beam by passage through a charge-exchange gas cell.

763603

Fig. 30. (a) Space-charge-controlled direct energy converter. (b) Variation of overall injector efficiency with deuteron beam energy. (From Ref. [151])

ELEVATION OF TFTR BEAMLIN

Fig. 31. Example of proposed beam line for TFTR device.
(From Ref. [157])

763607

Fig. 32. (a) Measured T_i in the ATC device before and after compression of D^+ plasmas with and without beam heating. Compression phase is indicated by dashed lines. (b) T_i vs time in the ATC device with beam injection of 230 kW in the period $t = 20$ -32 ms. (From Ref. [84])

Fig. 33. (a) Charge-exchange neutral spectra parallel and perpendicular to the magnetic axis in the ORMAK device for steady-state injection of 26-keV H^+ beams at 26° with respect to the axis. Circles show experimental data. Curves are computed for several Z_{eff} . (From Ref. [162]). (b) Charge-exchange spectrum parallel to the magnetic axis in the ATC device for quasi-steady-state tangential injection of 14.2-keV D^+ beams. Theoretical curve is computed for $Z_{eff} = 4$ and central $\tau_{cx} = 6$ ms. (From Ref. [23])

Fig. 34. Neutron flux resulting from injection of a pulsed 15-keV, 3-A D^0 beam into a D^+ plasma in the ATC device. (From Ref. [175])

TOKAMAK FUSION TEST REACTOR

- 1 - VACUUM VESSEL
- 2 - TOROIDAL FIELD COILS
- 3 - OHMIC HEATING FIELD COIL
- 4 - EQUILIBRIUM FIELD COIL
- 5 - DEVICE SUBSTRUCTURE
- 6 - SHIELDING
- 7 - WATER COOLING HEADERS
- 8 - VACUUM PUMP
- 9 - NEUTRAL INJECTION DUCTS

Fig. 35. Trimetric view of the TFTR device. (From
Ref. [124])

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763693
 Fig. 36. Vertical cross section of TCT actinide burner reference design. Dimensions in meters. (From Ref. [192])