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STATUS OF TOKAMAK RESEARCH

October 1979

U.S. Department of Energy
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STATUS OF TOKAMAK RESEARCH

October 1979

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FOREWORD

This statement on the status of the tokamak program may be regarded as an update of "Status and Objectives of Tokamak Systems for Fusion Research," a document of similar scope prepared in 1974 by a review panel chaired by S.O. Dean. The intervening five years witnessed an explosion of knowledge in tokamak research characterized by the increased breadth of the program, the steady flow of data from devices of increased size and toroidal field strength, and the maturation of neutral beam injection technology. The original intent of this document was to record the program status on the occasion of the Eighth International Conference on Plasma Physics and Controlled Nuclear Fusion Research held in August, 1978 in Innsbruck, Austria, but some revisions have been made to include developments of the past year.

The endeavor was begun at the urging of N. Anne Davies and was nourished by her continual support. The technical content is due primarily to the contributing authors; I am especially indebted to Paul Rutherford and Rob Goldston for their very substantial contributions.

In addition to those identified as co-authors, a host of others helped to enrich the work in detail: D. R. Baker, R. A. Blanken, F. P. Blau, R. W. Conn, R. L. Freeman, R. W. Harvey, G. L. Jahns, T. H. Jensen, R. L. Miller, N. Ohya, P. B. Parks, D. E. Post, L. H. Rovner, A. M. Sleeper, W. E. Toffolo, and R. E. Waltz.

Finally, I am delighted to acknowledge the heroic efforts of Verla Cravens, who was solely responsible for the design and production of the document.

John M. Rawls,
October, 1979

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1. INTRODUCTION

For over a decade, since the first announcement of successful results on the T-3 tokamak, the tokamak has maintained a preeminent position among magnetic confinement concepts (*Ref. 1-1*). The status and objectives of the tokamak program were reviewed in 1974 (*Ref. 1-2*) in preparation for a major expansion of the U.S. magnetic fusion program. In the succeeding five-year period, a number of experimental successes have been realized with major implications for the required parameters for tokamak reactors. Chief among these successes have been the observation of good ion and electron energy confinement at reactor level collisionalities in the Princeton Large Torus (PLT) and the cumulative data from a wealth of tokamaks indicating that confinement scales favorably with those parameters that must be significantly increased to reach the reactor regime. Given results of such significance, it is timely to reassess the status and objectives of the tokamak program.

The present document is intended as an overall review of the tokamak program, with particular emphasis upon developments over the past five years in the theoretical and experimental elements of the program. Considerably less attention is devoted to development and technology questions although the issues do arise in the context of discussions of plasma heating and the current perception of tokamak reactors.

The basic components of the tokamak are illustrated in Fig. 1-1. In its most primitive form tokamak confinement relies upon the simultaneous action of three magnetic fields: a strong toroidal magnetic field created by external coils, a weaker poloidal magnetic field created by a toroidal plasma current inductively driven by poloidal field coils, and a vertical magnetic field created by field-shaping windings used to maintain equilibrium. The toroidal plasma current also provides a source of heating for the plasma, one that is especially effective at low plasma temperatures.

Variants upon the basic tokamak scheme abound. Additional coils are utilized both to shape the plasma cross-section for studies of non-circular configurations, and as divertors to form open magnetic field lines at the surface of the plasma to remove particles and energy from the system. In addition, a number of heating techniques are used to supplement, and in some instances dominate, the heating produced by passing a current through a plasma of finite resistance. These techniques, which will be required to raise the plasma temperature to the level needed to produce a self-sustaining reaction, include the injection of energetic neutral particles (neutral beam injection) and the use of electromagnetic waves at a variety of frequencies extending from near dc to millimeter wavelengths.

The remainder of this section presents a brief synopsis of the history of the tokamak program, a summary of the operating parameters for the principal tokamaks throughout the world, and a discussion of some key outstanding issues of the tokamak program. The remaining sections will discuss in turn plasma confinement in tokamaks, from both a theoretical and an experimental viewpoint, plasma heating, and finally issues in tokamak design.

Fig. 1-1. Basic components of the tokamak.

1.1 BRIEF HISTORY OF TOKAMAKS

Tokamaks were developed during the 1960s in the USSR under the direction of the late Lev Artsimovich. Generically, they represent a subset of toroidal Z pinches in which a large toroidal magnetic field is utilized to establish gross stability. The first major successes, temperatures in the multi-hundred eV range with confinement considerably in excess of predictions based on Bohm diffusion, were announced by Artsimovich in 1968 (*Ref. 1-3*), and were confirmed by a group of British diagnosticians in 1969 (*Ref. 1-4*). In retrospect, it appears that the early successes may have had their origin in the relative economy of the tokamak in heating and containing large volumes of plasma. The larger volumes compared to other contemporary devices (e.g., the stellarator) implied a decreased surface-to-volume ratio and a reduced susceptibility to energy losses due to the influx of impurities from the wall.

The successes of the tokamaks in the late 1960s spawned a variety of similar devices and ultimately led to the widespread adoption of the tokamak as the leading magnetic confinement approach. The characteristic parameters of these devices are shown in Table 1-1. The model C Stellarator was converted to the ST Tokamak, and soon reproduced Artsimovich's results (*Ref. 1-5*). Neutral beam injection as a method for auxiliary plasma heating was first demonstrated in the Adiabatic Toroidal Compressor or ATC (*Ref. 1-6*); this was followed by a sequence of successful

TABLE 1-1
DESIGN PARAMETERS AND "TYPICAL" PLASMA PROPERTIES FOR MAJOR U.S. TOKAMAKS

Tokamak	a (cm)	R ₀ (cm)	B _T ^a (kG)	I _p ^a (kA)	T _e (0) ^b (eV)	T _i (0) ^b (eV)	n _e ^b (cm ⁻³)	τ _E ^b (ms)
ST	13	109	50 (40)	130 (70)	2500	600	4 × 10 ¹³	10
ATC ^c	17	90	20 (15)	100 (60)	1100	250	1.5 × 10 ¹³	5
ORMAK	23	80	25	340 (170)	1500 ^d	1800 ^d	3 × 10 ¹³	10
ALCATOR A	10	54	120 (80)	400 (200)	1200	800	6 × 10 ¹⁴	20
DOUBLET II-A	14 × 40	66	10 (8)	320 (150)	400	200	3 × 10 ¹³	4
PLT	45	130	45 (32)	1400 (550)	4000 ^d	6500 ^d	8 × 10 ¹³	100
ISX ^e	25	90	18 (12)	160 (110)	1100 ^d	900 ^d	6 × 10 ¹³	20
DOUBLET III ^f	45 × 150	143	40 (20)	5000 (2000)				
ALCATOR C ^f	17	64	140 (70)	1000 (400)				
PDX ^g	45	145	24 (24)	500 (500)				

^aDesign values (normal operating values in parentheses).

^bMaximum obtained – not necessarily in same discharge.

^cBefore compression.

^dWith neutral beam heating.

^eNot yet at maximum performance.

^fJust begun operation, B_T achievable power supply limited at present.

^gJust begun operation.

experiments employing beam injection on a variety of devices, the most thorough analyses being carried out on ORMAK. Heating via electromagnetic waves at a number of characteristic plasma frequencies and at modest power levels was investigated in many smaller tokamaks throughout the world. The first attempts to quantify the effects of non-circular plasma cross section were explored in Doublet II (*Ref. 1-7*); such effects are important in optimizing β, the ratio of plasma pressure to magnetic field pressure, which is a key parameter in the economics of generating power from tokamaks. The use of very high field and high energy density was investigated in experiments on Alcator A, a device which has produced density-confinement time products, nτ, in excess of 10¹³ cm⁻³ sec, more than an order of magnitude greater than those achieved in T-3 (*Ref. 1-8*).

The generally successful results of these experiments paved the way for a second generation of tokamaks designed to study tokamak physics at temperatures and collisionalities approaching those anticipated for a reactor. The first device with the necessary size and heating power was PLT. In the summer of 1978, with 2.5 MW of injected neutral beam power, PLT achieved a peak ion temperature in excess of 6 keV (*Ref. 1-9*), more than 20 times greater than that achieved in T-3 and greater than the minimum required for energy breakeven in a D-T plasma. The plasma collisionality, a key measure of the susceptibility of the plasma to microinstabilities, reached values characteristic of a reactor but the ion confinement remained consistent with classical predictions (though with fairly large error bars) and in fact the electron thermal conductivity appeared to improve. One implication is that more compact reactor designs will be possible than those based on the anomalous "trapped particle instability" scaling laws; in fact the predicted instabilities are observed experimentally but prove to be harmless.

Other major facilities in the second generation of tokamaks include the Impurity Studies Experiment (ISX-B), a small, low field tokamak which will investigate β limits, the use of very high frequency microwaves (ECRH) for heating, and the use of a bundle divertor to control impurities; the Doublet III, a large tokamak which will investigate non-circular cross section plasmas at large current and energy densities; the Poloidal Divertor Experiment (PDX), which will investigate impurity control via the use of special (poloidal divertor) coils to direct impurities out of the device; and the Alcator C device, a larger, higher field version of Alcator A which will investigate the confinement of high density plasmas and should achieve $n\tau > 10^{14} \text{ cm}^{-3} \text{ sec}$, a value characteristic of reactor grade plasmas.

A major portion of these facilities will involve auxiliary heating necessary to obtain temperatures approaching those required for reactor operation. ISX-B, Doublet III, and PDX will rely principally upon neutral beams; Alcator C and PLT will utilize electromagnetic waves at the lower hybrid frequency and at the ion cyclotron resonance frequency. Upgraded versions of several of these devices are being contemplated to further probe the reactor regime.

The second generation of facilities will be followed in the early to mid-1980s by a third generation of devices, which will exhibit a number of the characteristics of actual fusion reactors. These devices, cataloged in Table 1-2, comprise the Tokamak Fusion Test Reactor (TFTR), the Joint European Tokamak (JET), the JT-60 tokamak in Japan, and the T-15 tokamak in the USSR. TFTR and JET will employ D-T fuel and sufficient auxiliary heating to demonstrate plasma energetic breakeven. JT-60 will reproduce similar plasma conditions, but will not employ tritium; in addition, it will be capable of producing discharges which persist for many seconds. T-15 will demonstrate the operation of superconducting toroidal field coils on a large tokamak.

1.2 KEY OUTSTANDING ISSUES OF THE TOKAMAK PROGRAM

The basic appeal of the tokamak is its conceptual simplicity. In its most primitive form the tokamak is axisymmetric, circular in cross section, low beta (and therefore more easily susceptible to external control), and efficient in the sense that it utilizes the same poloidal field coils for both heating and plasma confinement. The inherent simplicity of the tokamak allows considerable room for adaptability and improvements, e.g., one or more auxiliary heating methods, an energy conversion medium surrounding the plasma, superconducting coils, moderate beta, and a non-circular cross section; additional systems may be employed for impurity control, fueling, and to sustain the plasma current indefinitely. This evolutionary development of the primitive tokamak

TABLE 1-2
THIRD GENERATION TOKAMAKS

Tokamak (Country)	a (cm)	R_0 (cm)	B_T (kG)	I_p (kA)
TFTR (USA)	85	250	52	2500
JET (Europe)	130 x 200	280	28	3000
JT-60 (Japan)	100	300	50	3000
T-15 (USSR)	70	240	35	1400

into a concept suitable for the core of a power plant can be expected to continue as the tokamak program addresses reactor relevant issues. Some of these issues include:

1. **Beta.** The plasma beta is a measure of the magnetic field utilization of the tokamak. The thermonuclear power density for D-T reactions is given by

$$P_t = n_D n_T \langle \sigma v \rangle_{D-T}(T) \cdot 17.6 \text{ MeV}$$

where n_D and n_T are the deuterium and tritium density, respectively, and $\langle \sigma v \rangle_{D-T}$ is the D-T fusion reaction rate. At temperatures of interest for ignited or self-sustaining plasmas, $\langle \sigma v \rangle_{D-T} \propto T^2$ so $P_t \propto (\beta B^2)^2$. Hence, for a given P_t or for a given wall loading, the beta will determine the necessary magnetic field strength, an important variable in overall plant cost. To date, average values of beta have not exceeded on the order of 1%, the limitation being the available auxiliary plasma heating. Theoretical predictions indicate a threshold for magnetohydrodynamic instabilities at average betas of $\sim 2\%$ in circular cross sections of moderate aspect ratio and 5–10% in carefully tailored noncircular cross sections but these conclusions must be regarded as preliminary since there is virtually no contact between theory and experiment on the question of β limits. In addition, the limiting mode in many cases (the ballooning mode) should not lead to loss of gross confinement but merely to enhanced transport. Hence, the β limits so deduced are not inviolable. Sufficient heating power will be available on ISX-B (3 MW), Doublet III (7 MW), and PDX (6 MW) in 1980 to test these theoretical predictions and to qualify reactor design β values, which are generally in the 5–10% range. An important aspect of these experiments will be the determination of what, if any, profile control measures will be needed to access the high β regime. The theoretical and experimental status of beta limits is reviewed in Section 2.1.

2. **Confinement Scaling.** The minimum plasma size needed to achieve ignition conditions depends most sensitively on the dependence of the dominant underlying transport mechanisms on plasma and machine parameters. There is now a considerable body of evidence to support the contention that the energy confinement time, τ_E , scales quadratically with plasma size, as would be expected from a white noise or random walk diffusion process, and also scales linearly with plasma density. Furthermore, there are preliminary indications that confinement improves with increasing temperature. Extrapolations based on such expressions must, however, be regarded as speculative because of the empirical nature of the results and the consequent uncertainty in the limits of validity. See Section 2.2 for further discussions of the whole range of transport questions.
3. **Impurity Control.** There are stringent requirements on the plasma purity required to produce and sustain a reactor grade plasma. While increased experience with tokamak operation, particularly in the areas of particle control and plasma position control, have resulted in plasmas which are essentially impurity free, there is concern that the longer pulses and higher wall loadings necessitated in reactor operation will result in significantly increased impurity generation. In addition, the D-T fusion process yields a helium "ash" which must be removed to avoid prematurely terminating the burn.

A number of planned programs, notably the poloidal divertor experiments on PDX and the bundle divertor experiments on ISX-B, will focus on these issues; see Section 4.2.

4. **Burn Physics.** The operating point of a tokamak reactor will be determined principally by an equilibrium between two processes with which we as yet have little experience: alpha particle heating and an undetermined loss (or burn control) mechanism. TFTR D-T experiments are expected to reveal much about the behavior of 3.5 MeV alpha particles, in particular whether they behave classically, as experience to date with the hot ions produced by neutral beam injection might suggest. Direct information about the dynamics of a burning plasma, in which alpha heating dominates the energy balance, must await the production of high Q or ignited plasmas.
5. **Pulse Length.** High duty cycle is clearly a desirable objective for any reactor system. The principal limitation on the pulse length of present-day tokamaks is the fact that the current is inductively driven and is therefore limited in time by the volt-second capacity of the transformer. Ideas are presently under investigation for driving the current by momentum transfer from either neutral beam or radio frequency waves, or by injection of relativistic electron beams. Other issues under investigation which will affect long pulse operation include maintaining the magnetic configuration for times in excess of the characteristic time for resistive diffusion, controlling impurity influx as the plasma wall equilibrium time is approached, controlling fuel depletion and ash removal in thermonuclear plasmas, and regulating the thermonuclear burning processes to avoid a thermal instability. With the addition of long pulse heating sources, several existing tokamaks (notably Doublet III and PDX) would be able to achieve sufficiently long pulses to address a number of these issues. The questions associated with long pulse operation are addressed in Section 4.3.
6. **Maintenance.** Existing tokamaks present considerable difficulties for access and repair and these problems will be greatly aggravated when the additional requirement of remote maintenance, a necessity for reactor operation, is imposed. One of the reasons for these difficulties is that tokamaks operate at small aspect ratio to achieve higher current at a given toroidal magnetic field and to optimize the probability of achieving high beta values (on the basis of theoretical expressions for beta limits). The use of increased auxiliary heating, and better understanding of beta limitations, may allow a larger aspect ratio with corresponding improved access. In addition, considerable improvement in access can be achieved by employing fewer, larger toroidal field coils. This approach, while unacceptable from a power supply point of view for copper coils, is an attractive alternative for superconducting coils, which are of course a necessity for reactor designs. An alternative solution to the question of access is to remove the hole in the middle of the tokamak and to produce the confining poloidal and toroidal fields by current drive in the plasma.

One such approach, the spheromak, has attracted considerable attention recently (*Ref. 1-10*) and, while as yet untested, may develop naturally from a program aimed at noninductively driven currents in tokamaks.

The tokamak program also faces a number of other major issues in engineering and technology areas but these are for the most part outside the scope of this document. In any case, the rapid progress being made in all phases of the program has created considerable optimism that the tokamak will prove to be a scientifically and technologically feasible power source.

2. PLASMA CONFINEMENT

Since the only plasma distribution in thermodynamic equilibrium is a spatially uniform, Maxwellian velocity distribution, a confined plasma is necessarily out of equilibrium. The most obvious relaxation mechanism, discussed in Section 2.2.1, is by Coulomb collisions; in a tokamak geometry, the resulting particle and energy flows are termed neoclassical transport. Fortunately, the cross sections for such processes are sufficiently small at high plasma temperatures that very compact, attractive tokamak reactor designs are possible if no other transport mechanisms are operative. However, the plasma can also relax toward equilibrium through various types of collective plasma instabilities. Plasma instabilities are customarily divided into two classes: macroinstabilities (derivable from a fluid description) and microinstabilities (derivable only from a kinetic description).

The most important and relevant macroinstabilities for a tokamak plasma are discussed in Section 2.1. In general, the basic theoretical description of low β , large aspect ratio, circular cross section plasmas (Section 2.1.1) is confirmed by the observed operating regime in existing devices (Section 2.1.2). While disruptive instabilities (Section 2.1.3) and noncircular cross section effects (Section 2.1.4) need further exploration before a confident extrapolation can be made to reactor-grade plasmas, present theory predicts the existence of tokamak configurations with reactor-relevant parameters that are absolutely stable to all the important macroinstabilities.

In contrast to the macroinstabilities, it is likely that not all microinstabilities can be stabilized. Because of their generally small wavelengths, and the small amounts of energy driving them, the microinstabilities are generally expected to grow rapidly into the nonlinear regime, where their further growth will be severely limited. Thus, they are not expected to be as large-scale and chaotic as the macroinstabilities. However, they can cause rapid, small-scale, but relatively ordered plasma transport, and hence can be detrimental to plasma confinement. Since the transport caused by microinstabilities is in general in excess of, and relatively unrelated to, classical transport, it is usually referred to as "anomalous" transport. This topic, the subject of Section 2.2.2, is taken to include not only drift waves and trapped particle modes, but also the effects of perturbations in the magnetic topology, an area that closely entwines magnetohydrodynamics and transport considerations.

Section 2.2.3 describes the general features of typical tokamak discharges and discusses the primary determinants of the particle and energy flow patterns. Particular attention is devoted to classes of experiments that can reveal the functional dependence of the underlying scaling laws. In general, the correlation of the observed ion energy transport with the predictions of neoclassical theory is quite good. However, the electron particle and energy flows are not yet understood on a quantitative basis; indeed, the elucidation of the fundamental physical processes responsible for such behavior remains one of the major areas of basic tokamak research. Generally, empirical rather than theoretical scaling laws are employed for device design purposes. The extraction of such scaling laws from experiment is dealt with in Section 2.2.4.

2.1 PLASMA STABILITY CONTROL

Tokamak experiments have shown that with proper attention to technical requirements, the gross behavior of toroidal discharges can be controlled rather well over a wide range of parameters. Magnetohydrodynamic (MHD) models of plasma equilibrium and stability have been studied extensively and have shown generally good correlation with existing data for large aspect ratio, low β , circular cross section tokamaks. Such models appear to offer good promise of mapping out the accessible parameter regime for tokamak operation.

Some of the basic concepts of ideal MHD theory are dealt with in Section 2.1.1 in the course of deducing rough predictions for the parameter regimes accessible to tokamak operation. For comparison, the experimental situation is described in Section 2.1.2. In comparing experimental observations with the predictions of linear stability theory, a body of work that constitutes the bulk of the literature on MHD, it is important to keep in mind that nature does not draw a sharp distinction between stable and unstable regimes. In particular, it is not essential to avoid instabilities which may either be removed by a minor modification of plasma configuration or whose nonlinear consequences are tolerable. For example, if the saturated states of the modes that limit the plasma pressure take the form of enhanced transport, a certain degree of instability may be desirable, the tradeoff being between the improvement in β and the degradation of confinement. In such a situation, nonlinear theory is essential for a meaningful comparison with experiment.

One particular form of MHD behavior whose experimental consequences seem to be rooted in essentially nonlinear phenomena is the major disruption. Because of the potentially major implications for reactor design, the topic of disruptions has been singled out for consideration in Section 2.1.3. This section deals with the observed properties, speculations on the underlying physical mechanisms, and the implications for reactor design.

Finally, Section 2.1.4 examines the theoretical and experimental status of noncircular cross section plasmas, an approach that seems to be essential to producing the power densities required for economically competitive reactor designs.

At the risk of gross oversimplification, the status of plasma stability, as detailed in the sections just outlined, can be summarized as follows. Experimental tokamak equilibria now include a variety of noncircular shapes, and these equilibria continue to be adequately described by ideal MHD theory. The available operating range of tokamak plasmas has been considerably expanded, primarily as the result of experimental techniques that maintain broad current profiles. The observed helical modes and internal disruptions are compatible with resistive tearing-mode theory. It is still not clear whether a general mechanism exists for the disruptive instability, though numerous features of these observations also indicate tearing modes. Work with shaped plasmas thus far confirms the expectation that these should be capable of higher β values. The tests for critical β values are not yet in hand; these await a next series of experimentation, with massive auxiliary heating.

2.1.1 Elementary Considerations

The fundamental considerations involved in gross plasma positional stability are exhibited in Fig. 2-1. While an axisymmetric plasma in a purely toroidal magnetic field, B_T , cannot be in equilibrium but will instead expand in the direction of the major radius, magnetohydrodynamic equilibrium can be achieved via the addition of a poloidal field component, B_p , thus producing helical magnetic field lines. In a tokamak, the poloidal field is generated primarily by a toroidal current induced in the plasma, as in the secondary winding of a transformer. Establishing equilibrium also requires an externally applied magnetic field, B_v , which couples with the plasma current to produce

Fig. 2-1. Plasma equilibrium in major radius can be provided with a copper shell (A) or by external coils only (B). In either case, the required inward force (along-R) is due to the Lorentz force from plasma current flowing across the "vertical field" B_v . In case (A), the B_v -field is produced by image currents in the copper shell; in case (B), the B_v -field is externally applied. Total poloidal field patterns are shown above; B_v -field patterns by themselves are shown below.

an inward force along the major radius (Ref. 2-1). The B_v -field in early devices was generated (transiently) by eddy currents in a copper shell (Fig. 2-1A); in recent machines, it is driven by programmed external coils (Fig. 2-1B).

The degree of twist of the field lines is an important quantity for MHD considerations. The traditional measure of this quantity is the safety factor q , defined by $q = 2\pi/\iota$, where ι is the rotational transform, i.e., the angle through which a field line passes in the poloidal or θ -direction upon completing one revolution around the toroidal or ϕ -direction. For a general plasma cross section, q on a given flux surface may be written in terms of an integral over the projection of the flux surface on a poloidal plane

$$q = \frac{RB_T}{2\pi} \oint \frac{d\ell}{B_p} R^{-2} \quad (2-1)$$

This reduces to the more familiar result $q = rB_T/RB_p$ for the case of a large aspect ratio circular flux surface.

Plasmas can be unstable to a variety of fluid or MHD instabilities. There are basically three types of ideal modes:

1. Axisymmetric modes, which arise partially from the shaping of the plasma boundary;
2. Global kink modes, for which the plasma surface is disturbed;
3. Internal modes, which leave the plasma surface essentially undisturbed.

Broadly speaking, axisymmetric modes can be eliminated or controlled by proper design of the plasma control circuitry while kink and internal modes limit the parameter regime accessible to tokamak operation. A general treatment of this subject is contained in *Ref. 2-2*, which is also a good jumping off point to the extensive literature in the field. In discussing such modes, it is useful to introduce notation to describe their approximate spatial dependence:

$$f(r) \exp[i(m\theta + n\phi)], \text{ where } m = 0, 1, 2, \dots; n = 0, 1, 2, \dots$$

Consider first the axisymmetric ($n=0$) modes. In early experiments, these were stabilized (or at least slowed down) by eddy currents in the copper shell. For most present tokamaks, the externally driven field is tailored to have an appropriate gradient or slight outward curvature as in Fig. 2-1B so that $n=0$ displacements are stable. It is not always possible to stabilize $n=0$ modes completely, but the residual displacements have been found to be controllable by relatively simple passive or active feedback circuitry. Such modes are particularly important in the design of non-circular cross section tokamaks, in which case special measures are required for stabilization (see Section 2.1.4).

There can also be gross helical perturbations or kinks of the plasma column satisfying $m/n = q$. Such modes, in which the perturbation helix conforms to the field helix, are easily driven unstable because they require such a small investment of magnetic energy. Restricting attention initially to low beta, circular cross section devices, if the safety factor q drops below unity anywhere inside or outside the plasma, the $m = 1, n = 1$ kink mode is unstable for normal current density profiles peaked at the plasma center. The stability condition $q(a) > 1$ for this mode imposes an upper limit on the plasma current for a circular tokamak: $I = rB_p/2 < r^2 B_T/2R$, where rationalized MKS units are employed except for the current, which is in megamperes. If the plasma edge is at $r = a$, then $q(a) > 1$ corresponds to the Kruskal-Shafranov condition,

$$I(\text{MA}) < a^2 B_T/2R, \quad (2-2)$$

the basic limit on plasma current in tokamaks. Even if the condition $q > 1$ is satisfied everywhere, there can be residual instabilities (with $m/n > 1$), provided the point $q(r) = m/n$ falls in the vacuum outside the plasma, or into the outer plasma region of low conductivity. In the latter case, the unstable kink is called a resistive (or "tearing") kink mode.

If q exceeds 1 throughout the plasma, the tokamak is predicted to be stable to "flute" modes, which are helical perturbations driven principally by plasma pressure rather than by magnetic energy; furthermore, they are localized within the plasma. In tokamaks, stability to flute perturbations is brought about by magnetic shear and by a geometric effect that allows "minimum-B stabilization" to occur on the average. For sufficiently high values of $\beta_p = 8\pi n(T_e + T_i)/B_p^2$, the flutes can concentrate in regions of unfavorable curvature to become unstable "ballooning modes" (*Ref. 2-5*). This effect appears at a pressure level of $\beta \approx a/Rq^2$.

An approximate representation of the constraints imposed on tokamak operation by stability to ideal kink and internal modes is illustrated in Fig. 2-2, following *Ref. 2-2*. The constraints are exhibited in terms of $q(a)$, which determines the maximum plasma current, and $q(a)/q(0)$. The latter quantity has been translated into a measure of the peaking of the plasma current profile

Fig. 2-2. Stability diagram for ideal kink and internal modes. Stability to internal modes requires $q(0) > 1$, to $m=1$ kinks requires $q(a) > 1$, and to $m > 1$ kink modes requires a minimum peaking of the current profile (from Ref. 2-2).

using the parametrization $j(r) = j(0) (1 - r^2/a^2)^\nu$, the relation being simply $q(a)/q(0) = \nu + 1$. The criterion for stability to internal modes has been taken to be $q(0) > 1$. Kink stability hinges on the form of the current profile; in fact, complete stability against kinks may be obtained for any $q(a) > 1$ solely through sufficient peaking of the current profile ($\nu > 2.5$), while $q(a) < 1$ is unstable to $m = 1$ kinks for any current profile. While the region for stable operation extends to $q(a) = 2$, it is seen that the window diminishes dramatically as $q(a)$ drops below 3. Tearing modes further restrict the operating window (Ref. 2-4) but, since the growth rates for these modes are considerably slower than those for the ideal modes, complete stability to such modes may not be essential.

The two limitations discussed above, i.e. that global kink modes provide a practical limit of $q(a) \gtrsim 3$ and that ballooning modes impose the restriction $\beta \lesssim a/Rq^2$, combine to yield the constraint $\beta \lesssim a/9R$ or $\beta \lesssim 2-4\%$ for typical designs. More careful study has revealed that 2% is in fact a more realistic estimate and such a low number results in unattractive reactor prospects. However, significant improvements are predicted for vertically elongated plasma cross sections; this is the subject of Section 2.1.4. Moreover, results from experiments suggest that the condition $q(a) > 3$ is overly restrictive, albeit reduction in $q(a)$ is generally accompanied by some deterioration in confinement.

While the broad aspects of the MHD stability problem are fairly well understood, considerable analytic and computational work remains to be done in order to arrive at a detailed explanation of the present experimental results and at reliable predictions for future tokamaks. The areas most in need of theoretical clarification are the optimization of plasma cross sectional shape,

nonlinear effects (such as the disruptive instability), finite β effects, and the stabilization of resistive modes in high-conductivity, high- β , noncircular plasmas. Some of these topics will be addressed in subsequent sections.

2.1.2 Observed Operating Characteristics

The equilibrium position of tokamak plasma columns continues to conform well with ideal MHD theory. Centering within a conducting shell is accomplished by adding the required vertical field component through use of external windings. In the absence of this shell, external windings must also generate those components previously supplied by eddy currents. This method appears to work equally well, but requires a more sophisticated electrical system typically including feedback control of the external windings. Considerable effort has been devoted to windings and control systems that simultaneously maintain plasma position and noncircular plasma cross-section shape (elliptical, dee, droplet and doublet). The same physical principles apply to control of these complex magnetic boundary conditions as to the circular cross-section. *Reference 2-6* provides a useful review of control of equilibrium position and the associated magnetic diagnostics, and includes a number of citations to the literature.

2.1.2.1 Observation of Helical Structures — Oscillatory perturbations of the plasma column have been studied principally by means of magnetic pickup coils sensitive to δB_p and collimated detectors of soft X-rays (*Refs. 2-7, 2-8*). Rotating (helical) kink structures, typically with $m = 4, 3$, or 2 on the magnetic signals and with $m = 2$ and 1 in the X-rays [$m \leq q(a)$] are found, the toroidal mode number almost invariably being $n = 1$. The observed frequencies are typically 5 to 15 kHz. Observation of a sequence of modes of decreasing m values, from as high as $m = 14$ (*Ref. 2-9*), is typical of the magnetic signals during the buildup of the discharge current. It is believed that these modes play a role in preventing formation of a pronounced skin current (*Ref. 2-9*). The $m/n = 1/1$ perturbations appear when $q(0) \lesssim 1$ and cause spatial and temporal modulations of the soft X-ray flux that are localized near r_1 , the radius at which $q = 1$. The amplitude of this mode grows in time and then quenches abruptly as the result of an "internal disruption". The associated X-ray signals are periodically modulated at the mode frequency and also in a "sawtooth" fashion, $m/n = 0/0$, with a phase inversion near r_1 (*Ref. 2-8*).

This internal disruption prevents $q(0)$ values lower than about unity by limiting the central electron temperature. The temperature profile is peaked just before internal disruption but is flattened to somewhat beyond r_1 by the disruption. The sawtooth waveform is mostly a response to these changes in temperature, the relative density fluctuations being of considerably lower magnitude. The buildup in soft X-ray intensity for $r \lesssim r_1$ reflects continuous heating between internal disruptions, and the inverted waveforms for $r > r_1$ result from propagation of a heat pulse into the outer plasma (*Refs. 2-10, 2-11*). In present ohmically-heated tokamaks, the period of the sawtooth relaxations extends to 10 – 15 msec, and the $\Delta T_e/T_e$ at the center upon internal disruption can be 10 – 20% . The phenomenon can dominate energy confinement in the central region, but generally has a negligible effect on gross energy confinement because of the small volume enclosed by the $q = 1$ surface. Period and amplitude typically increase with addition of auxiliary heating [to 50 msec and 30% in PLT discharges with > 2 MW of neutral beam injection (*Ref. 2-12*)].

As shown in Fig. 2-3, discharges tend to be bimodal with respect to these perturbations, displaying either: (1) strong $2/1$ oscillations on poloidal magnetic flux loops ("Mirnov oscillations") and soft X-ray fluxes from near r_2 , but weak or no internal disruptions; or (2) strong internal disruptions and weak, possibly higher m number, Mirnov oscillations. A third possibility, though less frequently observed, is a small amount of MHD activity.

SHOT 9892 (TYPE A)

SHOT 9893 (TYPE B)

$\dot{B}_p$ AND SOFT X-RAY SIGNALS FOR TYPES A AND B DISCHARGES
 ($I = 155$ kA, $B_T = 21.6$ kG, $q_a = 4.6$, $P = 4.1 \times 10^{-4}$ TORR).

Fig. 2-3 Characteristic $\dot{B}_p$ and X-ray signals. The X-ray detectors were collimated to view across different chords of a minor cross section, with tangent radii from $r = -2$ to $+14$ cm. The set of signals on the left are dominated by the $m/n = 2/1$ mode ($r_2 \approx 12$ cm), and those on the right by internal disruptions ($r_1 \approx 6$ cm). (ORMAK data, courtesy J. L. Dunlap, ORNL.)

These large scale oscillations, especially the sawtooth and the $m = 2$, are observed by numerous other diagnostics. Thomson scattering has detailed the T_e fluctuations involved with the sawtooth (Refs. 2-13 – 2-17) as have electron synchrotron radiation detectors (Ref. 2-18) which have also provided information on δT_e for both the $m = 1$ and $m = 2$ modes. Other oscillations appear in the form of modulations of density, of emitted light, or power to the walls of the vacuum vessel, and of hard X-rays from bombardment of the limiter by runaway electrons.

Figure 2-4 shows a remarkable representation of an $m = 2$ structure derived by inversion of X-ray data from PLT (Ref. 2-19). Similar structures in the perturbed magnetic field have been mapped through use of probes inserted directly into the discharge for low energy density plasmas (Ref. 2-20).

Fig. 2-4 Contours of equal X-ray emissivity in a minor cross section of PLT, depicting an $m = 2$ structure (Ref. 2-20).

Experimental properties of the oscillations fit well the results of the theory of tearing modes (see for example, the comparisons given in Ref. 2-21), and in this context the structures are the predicted magnetic islands. The oscillations exist for times very long compared with theoretical growth times of linear modes, and the structures are much wider than a tearing layer. It follows that they represent nonlinearly saturated states of the instability.

At least through moderate levels of perturbation, the observed oscillations are consistent with rotation of the mode structures at θ -velocities similar to the electron diamagnetic drift of the structures in the main magnetic field. Including the added velocity for plasma electrostatic potential has led to closer agreement between observed and calculated rotational velocities (Ref. 2-22). In the ATC device, the rotation velocity of the oscillations increased in agreement with the expected increase in diamagnetic speed (Ref. 2-23). The rotation of the perturbations with the characteristic electron diamagnetic speed is qualitatively consistent with theoretical expectations for resistive kink modes, particularly in the nonlinear range (Ref. 2-24). An influence of unbalanced tangential neutral beam injection has been observed in PLT; there, beam-induced toroidal rotation has been interpreted as giving rise to poloidal propagation in the ion diamagnetic direction (Ref. 2-25).

Under some conditions, but particularly with modest $m = 2$ excitation, the $m = 1$ and $m = 2$ oscillations are at different frequencies; they are then regarded as independent structures. Under other conditions, especially at high levels of $m = 2$ ($\delta B_p/B_p$ approaching 1%), the modes appear to couple so as to be locked in frequency and phase (Refs. 2-13, 2-21 – 2-28). Further growth of the

modes leads to significant reduction of frequency and pronounced distortion of the waveforms from both magnetic and X-ray diagnostics. These effects seem correlated with the location of the limiter, and details of the structure have led to the suggestion that toroidal rotation is then involved (Ref. 2-26). Recent work on the TFR concluded that the apparent coupling is due to the emergence of a new $m = n = 1$ mode that, owing to the toroidal symmetry, is itself a part of the $m = 2$ mode, with the same frequency (Ref. 2-28).

2.1.2.2 Stable Operating Regime — An information representation of the “stable” operating regime of tokamaks is shown in Fig. 2-5 (from Ref. 2-29 and formed from the data base of Ref. 2-30).

The onset of disruptive instabilities (Section 2.1.3) provides limits at both high and low plasma currents. Qualitatively, the onset of disruption at high currents results from operation at too low an edge q ($q \sim a^2/l$), with the critical value quite sensitive to the effects of impurities (Z_{eff} in Fig. 2-5) and gas influx (density). The impurities and gas lower the electron temperature at the edge of the plasma, shrinking the discharge radius and producing an even lower q at some effective discharge radius. Disruptions at the low current boundary are explained similarly. Here the shrinkage of the discharge radius more than offsets the dependence of the safety factor on current. At very low plasma densities, there are intense runaway phenomena, often accompanied by positive voltage spikes and sometimes by negative spikes — but it is not clear whether these effects are related in any way to the usual disruptive instability.

Extensions of the operating regime are obtained by experimental techniques that maintain broader profiles. It is necessary to control the radiative edge cooling associated with high levels of low Z contaminants, and to avoid the hollow profiles that may result from the large energy sink associated with the presence of high Z contaminants in the plasma core. Maintaining a properly centered discharge is also important. In order to reach the highest densities, gas injection with a positive ramp in discharge current is effective (Ref. 2-32). Neutral beam heating can also have beneficial effects, allowing operation at lower q and at higher density (Refs. 2-29, 2-33). It is possible that the injected power works in combination with that from ohmic heating to resist shrinkage of the current profile (Ref. 2-31).

Fig. 2-5 The stable operating regime of tokamaks, in a representation using normalized plasma current and normalized density as coordinates. It is interesting to note that densities beyond unity on this scale represent densities beyond the 1976 data base for ohmically-heated tokamaks operated with steady gas fill (see Refs. 2-29, 2-30, 2-31).

Plasma energy confinement begins to deteriorate before the critical conditions for disruptions are reached. As noted above, strong internal disruptions can dominate energy confinement for $r \lesssim r_1$. This region is usually small and so the effect on gross energy confinement time, τ_E , is generally not significant. However, strong correlations are found between decreasing τ_E and the increasing levels of $m = 2$ perturbation that generally accompany reducing $q(a)$. A review of these observations along with development of a consistent representation of the results has recently been presented (*Ref. 2-34*). Large $m = 2$ oscillations have also been associated with deteriorating particle confinement (*Ref. 2-35*).

The role of the conducting shell in influencing stability of the tokamak discharge is qualitatively consistent with the prediction of the ideal MHD theory: that is, a shell is not expected to be important unless close fitting and then mainly at low q ($= m/n$). With very close-fitting shells, T-6 (*Ref. 2-36*) and DIVA (*Ref. 2-17*) operated with edge q values of about 1.2–1.3. For tokamaks with loose-fitting or resistive shells, the lowest q value has been obtained by ALCATOR, with stable discharges at densities $2 \times 10^{14} \text{ cm}^{-3} \leq \bar{n}_e \leq 5 \times 10^{14} \text{ cm}^{-3}$ and $q(a)$ from 1.9 to 2.0 for approximately 20 msec (*Ref. 2-37*). Maintaining fairly low levels of helical oscillations was crucial for attaining such discharges, and the most important factor in avoiding disruptions was precise control of the rate of rise of the plasma current. PULSATOR has operated at $q(a) \approx 2.1$ with $\bar{n}_e = 7 \times 10^{13} \text{ cm}^{-3}$, but these discharges were irreproducible and had a significant population of runaway electrons; increasing $q(a)$ to ≈ 2.5 made thermal, reproducible discharges (*Ref. 2-38*). A number of other tokamaks have also operated at $q(a) = 2.5$.

The basic MHD features discussed here, i.e., helical structures, sawteeth, and disruptions, are common to all tokamaks, including those that employ an axisymmetric divertor [DIVA (*Refs. 2-17, 2-39*), T-12 (*Ref. 2-40*)] or a bundle divertor [DITE (*Ref. 2-41*)], or that utilize compression [ATC (*Refs. 2-29, 2-42*), TOSCA (*Refs. 2-43, 2-44*), TUMAN 2A (*Ref. 2-45*)]. In fact, the MHD properties of ohmically heated stellarators operated at low ι_0 [CLEO (*Ref. 2-46*), R-0 (*Ref. 2-47*), W VIIA (*Ref. 2-48*)] have much in common with those of conventional tokamaks. With increasing ι_0 the $m = 2$ modes and disruptions are suppressed and eventually disappear, with operation possible at edge q values well below 2. The stabilizing effects are consistent with tearing-mode theory (*Ref. 2-48*).

Until quite recently, the β values obtained in conventional, quasi-steady discharges lasting a large fraction of a skin time or longer were relatively modest, $\beta_p \leq 1$ and $\bar{\beta} \leq 1\%$. High β values, β_p approaching R/a , had been demonstrated only in transient cases. For example, with a triangular shape component and compression, TOSCA has obtained $\beta = 3\%$ (*Ref. 2-44*). This value is considerably in excess of those calculated for emergence of the ballooning mode in this geometry, and there was no evidence of the mode. Preliminary results from ISX-B with on the order of 1 MW of neutral beam injection into a nearly circular cross section plasma have produced plasmas at high β_p ($R/2a < \beta_p < R/a$) that are predicted to be ballooning unstable. Even at the highest β values achieved to date ($\approx 1.8\%$), there is no significant deterioration of confinement.

2.1.3 Disruptive Instabilities

A disruptive instability is a sudden, large disturbance that develops very rapidly (typically tens-to-hundreds of microseconds in ohmically heated plasmas) and is characterized by hard X-ray bursts, flattening of the current profile with expulsion of poloidal flux and prominent negative voltage spikes, loss of energy, and shift of the plasma column to smaller major radius. Small disruptions may repeat several times; large ones may terminate the discharge.

Major plasma disruptions present a formidable design problem because of the rapid release of both thermal and magnetic energy. In general, disruptions are expected to occur at the limits of operation in current and/or density but the mapping of an essentially disruption-free operating regime has been an empirical exercise for each device. It is important to understand the underlying mechanisms to provide a basis for dealing with the thermal, electrical, magnetic, and mechanical design issues associated with such events for future tokamaks.

2.1.3.1 Characteristics of Disruptions — Whatever the fundamental reason for its origin, a disruption can be regarded as beginning with a broadening of the current profile. The consequent reduction of the internal inductance is reflected in a positive current step and an associated negative voltage spike. The impact on confinement depends on the interaction of the plasma with its surroundings, an interaction having its origin in the initial expansion in minor radius and the inward shift that accompanies the decrease in internal inductance. If the plasma-wall interaction releases little gas or material, the discharge can be expected to recover and the disruption is relatively minor. If the influx is sufficiently large, however, the current may be quenched in a short time and the disruption is a major one. Maintaining equilibrium during a rapid change in either the magnitude or the profile of the plasma current requires an equally rapid reduction of the applied vertical field; the general inability to accomplish this in large tokamaks will enhance the interaction with the inner wall and accelerate the rate of current decay. It is to be noted, however, that a close fitting, flux-programmed poloidal coil system can tolerate substantial change in either total plasma current or shape of the current profile without loss of equilibrium.

A detailed explanation of the sequence of events that culminates in a disruption is complicated by the fact that apparently similar final results can be arrived at for discharges with widely differing characteristic MHD signatures. Indeed, in a detailed study of a large number of disruptions on PLT using arrays of soft X-ray detectors (*Ref. 2-20*), no universal MHD characteristics of disruption emerged, although $m = 2$ oscillations were a common feature. Nevertheless, recent theoretical work on the nonlinear evolution of tearing modes (*Ref. 2-49*) has provided powerful evidence that such phenomena are fundamental to the description of disruptions.

One route by which discharges reproducibly proceed to disruption involves growing $2/1$ and $1/1$ modes and possibly involving other nonlinearly excited modes. Several detailed studies of such discharges have been reported (*Refs. 2-15, 2-26, 2-28, 2-41, 2-48, 2-50*). Since overlap of island structures results in stochastic motion of the field lines (see Section 2.2.2.4.5), the explanation for at least these particular disruptions may lie within the present framework of tearing mode theory.

Studies on PULSATOR have included the creation of static $m = 2$ islands at the $q = 2$ surface by the imposition of a magnetic field from external $\ell = 2$, $n = 1$ helical windings (*Refs. 2-26, 2-51*). Driven disruptions occurred when these islands were the same size as those in naturally disrupting discharges, and the time histories of the various diagnostic signals were similar. The critical island size for disruption was that which touched the limiter. Smaller induced islands had a stabilizing influence on the $m = 2$ mode, apparently due to a flattening of the temperature and current profile around the $q = 2$ surface. Use of $4/2$ and $1/1$ windings produced qualitatively similar effects (*Ref. 2-52*). These windings also had strong field components in resonance at the $q = 2$ surface.

The key factor in determining the design impact of a major disruption is the time scale over which it occurs. In this respect, it is important to distinguish between the width of the negative voltage spike, which is a plasma property and present in both major and minor disruptions, and the time required for the current to decay in a major disruption, which may be expected to depend on wall conditions and on plasma position control as well as on plasma properties. The width of the negative voltage spike apparently follows the linear growth rate of the $m = 2$, $n = 1$ mode.

However, if that time becomes long compared with the time required for impurities to flood in and cool the plasma edge, or cause a loss of equilibrium, then it may no longer be relevant. Therefore, using the growth time of this mode, typically of order 20 ms for an ignition tokamak (*Ref. 2-53*), to estimate the characteristic time for design purposes is likely to be overly optimistic. The duration of the negative spike on PLT has never been directly measured, but soft X-ray measurements imply plasma expansion times possibly as short as 50 μ s. On this basis, disruption times one-to-two orders of magnitude shorter than the predicted MHD time may be expected. Fortunately, the current decay times for the vast majority of major disruptions in PLT are in excess of 2 ms (*Ref. 2-19*). Another promising note is the preliminary indication that disruptions are less rapid in vertically elongated plasmas.

2.1.3.2 Design Implications — The design impact of a sudden loss of confinement is severe. Indeed, for all large tokamaks, major disruptions play an important role in first wall design, wall cleaning requirements, plasma control requirements, vertical field coil placement, and even impact toroidal field coil design when superconducting coils are employed. To give a feeling for the magnitude of the problem, the consequences of a major disruption on a typical long pulse, D-T burning tokamak design are outlined below.

The deposition of 100–200 MJ of plasma energy onto the surface of the first wall on a time scale short compared to the thermal diffusion time will cause rapid heating and subsequent vaporization of substantial quantities (of order kilograms) of wall material. The major unanswered question in this regard is the spatial deposition of the energy on the surfaces that bound the plasma. Excessive localization may give rise to intolerable design constraints. In any case, the magnitude of the problem is a powerful argument for the adoption of a sacrificial and easily replenishable surface on the first wall.

The limitation on the time rate of change of the magnetic field in a superconductor makes it desirable to sustain a toroidal current for several seconds after loss of the plasma current. The vacuum vessel may in principle be used for this purpose but, because of its small L/R time constant, it will be necessary to solve the problem of switching in low resistance current paths.

The forces associated with the currents induced in the vessel are large and more or less predictable, at least in the uniform part of the vessel. The inward pressure acting on the vessel due to the combination of the toroidal current and the diamagnetic current is on the order of ten atmospheres and, since it is a dynamic loading, it may well determine the expected lifetime of the vessel, depending of course on the frequency of disruptions. A better choice to carry the plasma current is the field-shaping coils. Since they have lower impedance, no switching is necessary. Furthermore, the associated stresses are generally smaller in this case. But this involves the incorporation of poloidal coils inside the toroidal field coils, a major problem for assembly and maintenance, particularly for remote handling.

In TFTR, the vacuum vessel, with an L/R time ~ 2 ms, is used to protect the coils from the voltage spike that accompanies a rapid disruption. This imposes a constraint on the mechanical design of the vacuum vessel. Nonetheless, the voltages induced are large enough to be the determining factor in insulation requirements for the poloidal coils. In order to further minimize the voltage standoff requirements, the positive and negative vertical field winding turns are interleaved. Spark gaps are employed in PLT and PDX in addition to passive circuitry to limit the maximum coil voltage.

It is obviously beneficial to provide a mechanism for control of major disruptions. In addition to the engineering advantages associated with the reduced thermal loading to the first wall and to the limiter, and reduction of severe $\mathbf{J} \times \mathbf{B}$ forces on coils, the prevention of the major disruption

also allows lower q operation, which, if β_p is limited by ballooning, implies a higher β operation. While some success has been achieved in disruption control, the techniques employed also present major design problems when examined in a reactor context.

Recent theoretical investigations (*Ref. 2-54*) suggest that disruptions can be prevented by using helical feedback coils to nonlinearly stabilize modes with poloidal mode number $m = 2$ and toroidal mode number $n = 1$. The feedback signal must be in phase with the mode, which rotates with a frequency close to that of the drift frequency ω_* . Thus, $1/\omega_*$ sets a lower bound on the response time for the feedback system. The needed coil current scales with the toroidal field and minor radius but is a quite modest requirement even for reactor application.

This approach has met with some experimental success. While the situation on ATC with $\ell = 2/n = 1$ windings feedback-controlled to the 2/1 mode was unclear (*Ref. 2-55*), stabilizing effects were evident in a similar arrangement in TO-1 (*Ref. 2-56*). In the R-0 device, feedback control of the 2/1 with multiple helicity windings ($\ell = 1, 2, 3/n = 1$) improved plasma characteristics (*Ref. 2-57*). Improvements have also been noted in R-0 with imposed high frequency (1–2 MHz) helical quadrupole magnetic fields (*Ref. 2-47*) and in TR-4 with a high frequency 0/0 mode component of discharge current (*Ref. 2-58*).

2.1.3.3 General Remarks on Tearing Modes — Generalizing beyond the above considerations, it is conceivable that the saturated nonlinear states of resistive tearing modes are responsible for much of the experimentally observed behavior in tokamaks. The associated magnetic perturbations provide plausible explanations for internal disruptions, Mirnov oscillations, major disruptions, and anomalous electron thermal transport; an exposition of this point-of-view with thorough referencing to the literature is contained in *Ref. 2-21*. With careful choice of operating parameters (especially by avoiding steep current density gradients), it appears that the instabilities saturate at tolerably low levels, leading at most to modest deterioration of confinement. The analyses carried out to date, however, have not been definitive because rather severe simplifying assumptions have been necessary to make the studies tractable. Work underway should extend the results to cases with realistic toroidal curvature, noncircular cross sections, and significant values of β . Such studies are just becoming feasible with the expanded computing capabilities now available.

Considerable work remains to be done on the theory of resistive instabilities in future higher-conductivity tokamak plasmas. It is well established (*Ref. 2-59*) that, as the conductivity increases, the tearing modes become overstable, with frequencies corresponding to rotations approaching the magnitude of the electron diamagnetic drift speed. In this limit, the growth rates are greatly reduced, and the modes eventually become either stable or radially convective (*Ref. 2-60*). This occurs at conductivities somewhat higher than those in present experiments. The same criterion applied to a reactor suggests that only the lowest mode ($m = 2$ in particular) remains unstable, because of the reduction in diamagnetic drift at larger size. Calculations using a slab model to study the effect of field curvature on resistive tearing modes suggest that the average magnetic well should stabilize these modes (or reduce the growth rates greatly) at high conductivity (*Ref. 2-61*). However, consideration of toroidal geometry and noncircular cross section effects reopen the question of the ultimate stability, even in the presence of a magnetic well, so that further work is required to reach a definite conclusion.

2.1.4 Noncircular Cross Section Plasmas

The preceding discussions have dealt almost exclusively with approximately circular cross section plasmas. Much of the recent thrust of the tokamak program, both theoretical and experimental, has been directed toward noncircular cross sections, such as those depicted in Fig. 2-6. The dee shape is a natural one in the sense that intense heating of approximately circular plasmas is predicted to produce such a shape; furthermore, the shape closely follows the shape of the toroidal field coils. The doublet shape, while more difficult to produce experimentally, has large shear, which leads to attractive MHD stability properties.

It follows from Eq. (2-1) that vertical elongation permits a greater poloidal magnetic field strength for fixed values of the toroidal magnetic field and the safety factor. A higher value of B_p is expected to be highly beneficial to plasma performance. Improved confinement is expected, if for no other reason than the fact that neoclassical losses will diminish as a result of the decrease in the banana width (see Section 2.2.1.6). The associated increase in the joule heating power density makes startup easier and permits operation at higher densities, further improving confinement. Improved MHD stability properties are also anticipated; for example, any modes that are β_p limited translate into less stringent limits on β . Equation (2-1) also shows that similar advantages can be gained by reducing the aspect ratio [the limiting case $R/a \rightarrow 1$ has been the subject of recent investigation (Ref. 2-62)]. Other advantages may accrue to elongated plasmas in the areas of neutral beam heating, because of a shorter path for penetration to the plasma core, and impurity control, where the small values of B_p near the corners of the dee facilitate the implementation of a poloidal divertor.

The potential advantages of noncircular tokamaks have been recognized for a number of years (Ref. 2-63), and several experiments have been undertaken to investigate equilibrium and gross stability. During the same period, many theoretical studies have been made of equilibrium (Ref. 2-64), stability (Refs. 2-65, 2-66, 2-67), and transport (Ref. 2-68). Extensive theoretical studies have dealt with the effect of cross-sectional shape on the maximum stable beta, taking into account the localized interchange ("flute") modes and ballooning modes (Refs. 2-65, 2-66, 2-67, 2-69). More recently it has been possible to evaluate the ideal (external) kink mode stability for conventional discharges with circular, elliptical or dee-shaped plasmas. These studies have generally indicated a significant increase in the maximum stable beta for carefully tailored shapes and radial profiles of current and pressure. Stable dee-shaped equilibria have been generated (Ref. 2-66) with (volume-averaged) values of β in the range 5–10% for R/a values typical of large tokamaks; and stable doublets have been found (Ref. 2-65) with $\beta \sim 10\%$. The latter, however, have not yet been fully analyzed for external kink modes, since the existing MHD stability codes are not applicable to systems with internal magnetic separatrices. The treatment of cross-section effects on resistive internal and kink modes awaits further code development.

Fig. 2-6 Tokamaks with noncircular minor cross-section: (A) ellipse, (B) dee, and (C) doublet. The axis of symmetry is to the left.

It was recognized very early in the development of the noncircular tokamak concept that a large class of macroscopic perturbations with axial symmetry ($n = 0$) could be unstable in elongated discharges. Early analysis (*Ref. 2-70*) of these modes used very simple models of the plasma column in which, for example, the current was assumed to be constant in a cylindrical plasma column with elliptical cross-sectional shape. Subsequent numerical studies of more realistic toroidal plasma models have shown that wall stabilization of these modes is effective (*Ref. 2-71*).

For tokamaks with just a single magnetic axis (circular cross section, ellipse and dee) the axisymmetric modes are readily calculated. If one disregards effects from a finite conductivity wall surrounding the plasma, the time scale for such modes is always the Alfvén time scale (using the poloidal field strength) and the dominating mode is an $m = 1$ mode with plasma motion parallel to the axis of symmetry. For such modes a finite resistivity of the plasma has no effect. For doublets the situation is quite different (*Ref. 2-72*), because of the existence of a hyperbolic magnetic axis. The most important mode is $m = 2$ -like; the plasma motion is such that the plasma tends to become elongated or compressed in the direction of the axis of symmetry. The growth rate for such modes is proportional to the plasma resistivity near marginal stability.

The status of ideal MHD for dee and doublet plasmas has been summarized in *Refs. 2-66* and *2-65*, respectively. The problem of gross vertical stability can be solved transiently by a copper shell (or an equivalent shunt-connected "passive" coil system) or on a longer time scale by active feedback. For steady externally-applied magnetic fields, the problem can apparently be eased for flat current profiles (*Ref. 2-73*) by rectangularizing the ellipse: stable cases with $b/a \sim 4$ are found to be possible. Similar techniques apply to both the dee shape and the doublet. Extensive numerical computations have been done to obtain MHD stability predictions for arbitrary but realistic current profiles, toroidal geometry, and finite β -values. However, an experimental demonstration that stable, high beta equilibria can be maintained in the presence of resistive and diffusive processes is lacking.

Some early attempts to numerically generate the sequences of equilibria used to study the stability properties discussed above met with difficulty in maintaining $q > 1$ at plasma pressures sufficiently large to produce the β values of interest (*Ref. 2-74*). The task of producing such high β equilibria has been greatly simplified through the imposition of a dynamical constraint that high conductivity plasmas are expected to follow during rapid heating (heating time $\ll$ skin time) (*Ref. 2-75*). This so-called FCT (flux conserving tokamak) model "freezes in" the $q(r)$ profile and thus allows the generation of high β equilibria for any choice of $q(r)$.

The earliest experiments (*Refs. 2-76, 2-77*) on noncircular tokamaks used shaped conducting walls to achieve the desired cross-sectional shape. This can sustain the shape only for limited times: the duration of the discharge cannot greatly exceed the skin time for penetration of magnetic flux into the conducting shell. Moreover, experimental access is severely restricted by the walls, so a substantial research effort has been applied to achieving shape control by external magnetic coils, under the varying conditions likely to be encountered in experiments. Substantial theoretical progress has been made in optimizing external shaping coil systems to achieve shaping over a wide range of plasma parameters; moreover, the theory has been validated to some extent by experiments.

Experimental programs focusing on dee and doublet plasmas in ISX-B and Doublet III, though at an early stage, have succeeded in generating reproducible discharges. In general, the observed MHD properties of noncircular cross section discharges [Doublet (*Refs. 2-77, 2-78*), TOSCA (*Refs. 2-43, 2-44, 2-79*), DIVA (*Ref. 2-39*), TNT (*Ref. 2-80*)] exhibit most of the features

seen in circular tokamak plasmas. The modes seem dependent on shape primarily through the safety factor, offering some validation of the above analysis of the advantages of vertical elongation. In particular, the values of $q(a)$ observed in elongated plasmas are similar to circular cross section results except in the cases of doublets, where larger $q(a)$ results from the geometrical effect of the proximity of the plasma edge to the separatrix, which of course has $q = \infty$.

There are preliminary indications that the overall stability of shaped plasmas is somewhat enhanced relative to circular ones. A potentially significant example is that major disruptions typically occur over a longer time scale in elongated plasmas. Even though theory has not provided definitive insight into this problem, the present experimental trend bodes well for next generation tokamak design.

Ultimately, the push toward elongated plasmas stems from the desire to increase the fusion power density, which scales approximately as β^2 . To summarize the status of MHD theory on this vital point, present predictions are that β -values of order 5–10% (as required in typical tokamak reactor designs) should be attainable for $R/a \sim 3$ –4, provided current profiles are not excessively peaked. Present experimental results generally show good qualitative agreement with the MHD theory, but direct experiments on β limits have not been carried out.

2.2 TRANSPORT

Despite an increasingly rich data base and the availability of highly sophisticated analytical and numerical tools, the fundamental mechanisms governing the diffusion of particles and energy from the center of the tokamak discharge to the plasma boundary remain uncertain. Fortunately, as discussed in Section 2.2.4, the scaling laws deduced from experiment predict that ignition will occur in devices which represent only modest extrapolation in size and/or field beyond TFTR.

One exception to the lack of contact between theory and experiment is the ion energy balance, where collisional processes seem to offer an adequate explanation of existing data. Before beginning a discussion of classical and anomalous transport in tokamaks, the subjects of Sections 2.2.1 and 2.2.2, respectively, it is useful to consider in some detail the orbits of individual particles in the tokamak geometry, especially magnetically "trapped" particles. These are illustrated in Fig. 2-7.

All particles gyrate in small circles about a field line of radius $v_{\perp}/\Omega_i$. The center of this gyration orbit is called the guiding center. For most particles ("untrapped" particles) the guiding center simply follows the field line around and around the torus. However, the strength of the toroidal field varies like $1/R$, where $R = R_0 + r \cos \theta$ is the major radial coordinate, and the field is thus larger on the inside of the torus than on the outside. It follows that the guiding centers of some particles ("trapped" particles) are reflected, due to the constancy of the adiabatic invariant $\mu = v_{\perp}^2/2B$, before they reach the high-field region on the inside of the torus. To a first approximation, these particles will simply "bounce" back and forth on the same field line, as shown in the upper part of Fig. 2-7. Writing $B \approx B_0 (1 - \epsilon \cos \theta)$, where $\epsilon = r/R$, it follows that the trapped particles will be those with $\mu B_0(1 - \epsilon) < \mathcal{E} < \mu B_0(1 + \epsilon)$, where $\mathcal{E} = v^2/2$, roughly corresponding to $v_{\parallel}/v \sim \epsilon^{1/2}$.

Fig. 2-7 Trapped particle orbits in tokamaks.

As the trapped-particle guiding centers bounce back and forth along the field lines, they are also subject to a drift

$$\vec{v}_D = \frac{mv_{\perp}^2}{2eB} \frac{\vec{B} \times \vec{\nabla} B}{B^2} + \frac{mv_{\parallel}^2}{eB} \frac{\vec{B} \times (\vec{B} \cdot \vec{\nabla}) \vec{B}}{B^3} \quad (2-3)$$

arising from the inhomogeneity of the magnetic field. For a tokamak, in which the dominant component of the magnetic field is the toroidal vacuum field varying like $1/R$, the instantaneous drift $\vec{v}_D$ is directed vertically upwards (or downwards, depending on the sign of the charge and the sense of the toroidal field), and has a magnitude given by

$$v_D \approx \frac{m}{eBR} \left(\frac{v_{\perp}^2}{2} + v_{\parallel}^2 \right) \quad (2-4)$$

Consider the projection of the guiding-center orbit onto a cross section of the torus, as shown in the lower part of Fig. 2-7. On the part of the guiding-center orbit that lies in the upper half of the torus, the vertical drift $\vec{v}_D$ will have an outward component in the r -direction; on the part of the orbit that lies in the lower half of the torus, it will have an inward component in the r -direction. The result of this component of the drift is that the projection of the guiding-center orbit has the "banana" shape shown in the lower part of Fig. 2-7.

The width of the banana orbit Δr is an important consideration in computing the resulting transport. Noting that the time the guiding center spends in the upper half of the torus is of order $qR/v_{||}$, it follows that

$$\Delta r \sim v_D (qR/v_{||}) \sim (r/R)^{1/2} \rho_p \quad (2-5)$$

where $\rho_p = mv/eB_p$ is the gyroradius in the poloidal field. A more rigorous derivation would proceed from the three constants of motion $\mathcal{E}$, μ , and $p_\phi = R(mv_\phi - e \int r B_p dr)$ where $v_\phi \approx v_{||} = [2(\mathcal{E} - \mu B)]^{1/2}$. Using these three constants, it is easy to show that the particle with the "widest" banana is the one that is marginally trapped, i.e., the one that is reflected at $\theta = \pm\pi$, and this particle has a banana width $\Delta r = 2(r/R)^{1/2} \rho_p$.

The vertical drift has an additional component *within* the magnetic surface. The effect of this component is to cause the banana orbit to precess around the torus in the ϕ -direction, as shown in the upper part of Fig. 2-7. The magnitude of this precession velocity is of order $mv^2/eB_p R$. (This estimate can be obtained by noting that the particle is unable to drift in the θ -direction, because of trapping, so that the component of $\vec{v}_D$ perpendicular to the field line but within the magnetic surface must come about from a motion of the whole banana in the ϕ -direction.) A more detailed treatment of this precessional drift shows that trapped particles precess in both directions around the torus, depending on their pitch angle. Most trapped particles precess in the "unfavorable" sense, meaning that the direction of precession is such as to tend to drive a Rayleigh-Taylor instability. However, trapped particles that are reflected close to $\theta = \pm\pi$ precess in the "favorable" direction.

The trapped particles form a modest fraction, of order $(2r/R)^{1/2}$, of the total particle population. More than half of the particles are untrapped and simply travel along field lines, thereby tracing out an entire magnetic surface. Binary collisions will tend to scatter particles in and out of the region of velocity space corresponding to particle trapping. Since trapping corresponds to the region of velocity space where $v_{||}/v \sim \epsilon^{1/2}$, it follows that the effective collision frequency for scattering out of the trapping region is $\nu_{\text{eff}} \sim (R/r)\nu$, where ν is the 90° scattering frequency. The time for a trapped particle to complete one orbit to and fro along a field line is called the "bounce time" $\tau_b \sim qR/v_{||} \sim (R/r)^{1/2} (qR/v)$. If a particle is scattered out of the trapping region before it has completed a single "bounce orbit" of this type, i.e., if $\nu_{\text{eff}} \tau_b \gg 1$, the concept of trapping loses all relevance. The value of the quantity $\nu_{\text{eff}} \tau_b$ for a thermal particle is usually denoted

$$\nu_* = \left(\frac{R}{r}\right)^{3/2} \frac{\nu q R}{v_{\text{th}}} \quad (2-6)$$

where $v_{\text{th}} = (kT/m)^{1/2}$ and ν is the Braginskii 90° collision frequency:

$$\begin{aligned} \nu_i &= 4\pi^{1/2} n_i e^4 \ln \Lambda / 3m_i^{1/2} (kT_i)^{3/2} \\ \nu_e &= 4(2\pi)^{1/2} n_e e^4 \ln \Lambda / 3m_e^{1/2} (kT_e)^{3/2} \end{aligned} \quad (2-7)$$

The quantities ν_{*i} and ν_{*e} are important parameters in tokamak research, providing a measure of the collisionality of the plasma. If $\nu_* < 1$, trapped-particle effects may be important; if $\nu_* \gg 1$,

such effects cannot play a significant role. It must be emphasized that the plasma may be "collisionless" as far as the majority of the particles are concerned, in the sense that the mean-free path, of a thermal particle (for 90° scattering) exceeds the distance around a field line (i.e., $\nu q R / v_{th} \ll 1$), while at the same time being sufficiently collisional for particle trapping to be irrelevant [i.e., $(R/r)^{3/2} \nu q R / v_{th} \gg 1$]. For a plasma with roughly equal ion and electron temperatures, the parameters ν_{*i} and ν_{*e} will be about the same; in such a plasma, the ions and electrons enter the regime where trapping is important at about the same degree of collisionality.

2.2.1 Theoretical Status: Neoclassical Transport

Confinement in tokamak systems is limited by various forms of cross-field transport, which can make its appearance as particle diffusion, involving a net loss of particles from the confinement region, or as heat conduction, involving the loss of relatively higher energy particles which are then replaced by lower energy particles.

Classical transport represents the irreducible minimum level of "leakiness" from a confinement system. The basic physical mechanism involved in classical transport is simply binary Coulomb collisions causing the particle guiding centers to step from one orbit centered on one magnetic surface to another orbit centered on a neighboring magnetic surface. In classical diffusion in a straight magnetic field, the "step size" will be a gyration radius ρ . The frequency with which such steps occur will be of order the 90° scattering frequency ν , so that the diffusion coefficient will be of order $\nu \rho^2$.

In a toroidal geometry, however, we have seen that particle orbits involve excursions away from a magnetic surface that are considerably larger than the gyration radius. In particular, trapped particles make "banana" orbits of width Δr given by Eq. (2-5). Thus it should not be surprising that transport is enhanced over that given by the classical result. Cross-field transport of this type, in which nonlocal properties of the magnetic configuration play a role in determining the transport coefficients themselves, is called "neoclassical transport".

2.2.1.1 Particle Diffusion and Thermal Conduction — If the collision frequency is sufficiently small, a particle will trace out several banana orbits before scattering occurs. However, collisional scattering out of the trapped-particle region of velocity space will take place with an effective frequency $\nu_{eff} \sim \nu / \epsilon$ where $\epsilon = r/R$. Clearly, a cross-field diffusion process will occur, involving a fraction $\epsilon^{1/2}$ of the particles, for which the step size is Δr and the frequency with which steps occur is ν_{eff} . The diffusion coefficient will be of order $D \sim (r/R)^{1/2} \nu_{eff} (\Delta r)^2 \sim (r/R)^{1/2} \nu \rho_p^2$, a factor $(R/r)^{3/2} q^2$ larger than the classical diffusion coefficient. This lowest collision-frequency regime is called the "collisionless" or "banana" regime, and applies when several complete banana orbits are traced out before collisional scattering occurs, i.e., when $\nu_* \ll 1$.

In the regime where $1 \ll \nu_* \ll (R/r)^{3/2}$, the collisions are too frequent to permit a trapped particle to complete its banana orbit, but typical untrapped particles make many transits around the torus before colliding with another particle. In this case, the largest contribution to the transport comes from a particular group of untrapped particles that have relatively small values of $v_{||}$ chosen in such a way that their effective collision frequency is comparable to their transit frequency around the torus, i.e., $\nu (v/v_{||})^2 \sim v_{||}/qR$. The overall diffusion coefficient can be estimated as $D \sim (v_{th}/qR) q^2 \rho^2$. Since this value of D is independent of collision frequency, this regime is usually called the "plateau" regime.

Fig. 2.8 Schematic illustration of the dependence of the neoclassical diffusion coefficient D on collision frequency ν , the transitions between regimes occurring at $\nu_1 = A^{-3/2} \nu_{th}/qR$ and $\nu_2 = \nu_{th}/qR$.

At even higher collisionality, there is a regime $\nu_* \gg (R/r)^{3/2}$ in which collisions are too frequent to permit a typical untrapped particle to complete a single transit around the torus before colliding. In this case, particles will make a random-walk around the torus, drifting in one direction in minor radius while they are on the "top" of the torus, and in the other direction while they are on the "bottom" of the torus. The overall diffusion coefficient can be estimated as $D \sim q^2 \nu \rho^2$, only a factor q^2 larger than the straight-field classical diffusion coefficient. This regime is known as the "collisional" or "Pfirsch-Schlüter" regime.

A schematic illustration of the dependence of the diffusion coefficient D on collision frequency ν in the three neoclassical regimes is shown in Fig. 2-8. The exact neoclassical diffusion coefficient D will, of course, have a "rounded" dependence on ν as shown by the solid curve in Fig. 2-8.

Both electrons and ions will diffuse across the magnetic field by these neoclassical processes. However, it is easy to see that the diffusion coefficient for ions is larger than that for electrons by a factor or order $(m_i/m_e)^{1/2}$ in all three regimes. Charge neutrality must, of course, be preserved. In some diffusion processes, the charge neutrality constraint gives rise to an "ambipolar" radial electric field that acts to hold back the more rapidly diffusing species. Actually, in both classical and neoclassical diffusion, no such ambipolar field is needed. It turns out that the particle diffusion arising from ion collisions is much smaller than our heuristic estimate would suggest (due to the constraint of momentum conservation for ion-ion collisions). Thus, the overall particle diffusion coefficient is of order the electron diffusion coefficient. Electron thermal diffusivity will clearly also be of this order.

Ion thermal diffusivity will have the larger value characteristic of the ion diffusion coefficient. We can obtain the following limiting results for the ion thermal transport in the three different neoclassical regimes:

$$\frac{\partial}{\partial t} \left(\frac{3}{2} n_i k T_i \right) = \frac{1}{r} \frac{\partial}{\partial r} \left(r \chi_i \frac{\partial k T_i}{\partial r} \right)$$

where

$$\frac{\chi_i}{n_i} = \begin{cases} 0.66 (r/R)^{1/2} \nu_i \rho_{ip}^2 & \text{(banana)} \\ 1.9 \left(\frac{k T_i}{m_i} \right)^{1/2} \frac{q \rho_i^2}{R} & \text{(plateau)} \\ (1 + 1.6 q^2) \nu_i \rho_i^2 & \text{(collisional)} \end{cases} \quad (2-8)$$

with

$$\rho_i = (2 m_i k T_i)^{1/2} / e B \text{ and } \rho_{ip} = \rho_i (B/B_p)$$

Actually, of course, the expressions given in Eq. (2-8) will be only approximations to the true thermal conductivity χ_i , which will have a smooth dependence on the collisionality parameter ν_{*i} . A treatment encompassing all three regimes at once has been given by Hinton and Hazeltine (*Ref. 2-81*), a reference which contains a thorough review of the literature on neoclassical transport. Their result for χ_i as a function of the collisionality parameter ν_{*i} , taken from *Ref. 2-82*, is shown in Fig. 2-9 for the case $r/R = 0.2$, where χ_i has been normalized to the plateau-regime value given in Eq. (2-8). For comparison purposes, the three limiting expressions given in Eq. (2-8) are also shown. Note that the more accurate result is considerably smaller than the approximate form in the plateau regime.

Fig. 2-9 Neoclassical ion thermal conductivity χ_i , scaled to the plateau regime value given in Eq. (2-8), as a function of ν_{*i} for the case $r/R = 0.2$ (from *Ref. 2-81*).

The neoclassical theory of *electron* particle and thermal transport in tokamaks proceeds along similar lines, although it is much more complicated because of the many "off-diagonal" transport coefficients that occur (Ref. 2-18). However, at least in a formal sense, all electron neoclassical particle and thermal transport processes are smaller than ion thermal conduction by a factor $(m_e/m_i)^{1/2}$, assuming $T_i \sim T_e$. Moreover, electron thermal transport in tokamaks is experimentally found to be anomalous, i.e., substantially greater than the neoclassical theory would predict. The same is almost certainly true of particle diffusion, although the experimental data is not so definitive on this question. Thus, since it does not seem to be particularly relevant to the experiments, no details of the neoclassical theory of particle diffusion and electron thermal conduction will be presented.

2.2.1.2 Bootstrap Current and Trapped-Particle Pinch — Neoclassical transport theory has a number of implications beyond simple particle and heat transport. For example, the same frictional forces between trapped and untrapped particles that cause the radial particle diffusion also cause a "bootstrap" current to flow in the plasma. In the collisionless regime, i.e., $\nu_{*e} \ll 1$, this current can exceed the ohmic heating current if $\beta_p > \epsilon^{-1/2}$. The bootstrap current is proportional to the pressure gradient and hence to the plasma β . The condition that the bootstrap current be kept weak enough so that $q > 1$ (for macroscopic stability) leads to the limit $\beta < \epsilon^{3/2}$, a constraint that may be more severe than that imposed by MHD ballooning stability ($\beta < \epsilon/q^2$). The bootstrap current is small near the magnetic axis, where the plasma is necessarily in the plateau of collisional regime. However, there has been speculation that the plasma current induced by neutral injection could act as a "seed" for the bootstrap current, so that the resultant system, if at the correct β -value, might be able to be run steady-state with an ohmic-heating transformer required only for the initial breakdown and heating phase. So far, however, this aspect of neoclassical theory has not been verified experimentally; in fact, the bootstrap current has never been observed and the recent observations on ISX-B of β_p values in excess of $\epsilon^{-1/2}$ cast doubt on its existence.

Another implication of neoclassical transport theory is that the ohmic-heating induction field E_z causes enhanced plasma "pinching", i.e., inward motion in minor radius. In classical straight-field plasmas, the inward pinch velocity would be $-(c\mathbf{E} \times \mathbf{B}/B^2)_r = cE_z B_p/B^2$. In toroidal plasmas, collisionless trapped-particles are found to drift inward, by what is often called the Ware pinch effect, at a velocity of cE_z/B_p , i.e., about 10^2 times faster than in the inward pinch in the classical straight field case. For practical tokamak parameters, the collisions are sufficiently rapid to keep the particle distribution functions approximately Maxwellian, but collisional friction between trapped and other particles causes a net inward drift of the plasma, both electrons and ions. In the collisionless, or "banana", regime the magnitude of the inward drift velocity is given by

$$v_r = -2.3 \left(\frac{r}{R} \right)^{1/2} \frac{cE_z}{B_p} \quad (2-9)$$

for a pure hydrogen plasma. However, the magnitude of the pinch velocity is much reduced in the plateau and collisional regimes. The result for all three regimes (Ref. 2-81) is plotted in Fig. 2-10 as a function of the collisionality parameter ν_{*e} for the case $r/R = 0.2$, where v_r has been normalized to the banana-regime value given in Eq. (2-9). We see that, for $\nu_{*e} \approx 1$, the pinch velocity is reduced to less than a third of its banana regime value.

The trapped particle pinch effect has been invoked as a possible explanation of the inward transport of newly ionized plasma in several gas-puffing experiments. If the pinch velocity had the value given in Eq. (2-9), this would certainly be sufficient to explain the data. However, many of the gas-puffing experiments operate in a regime where $\nu_{*e} \gtrsim 1$ over most of the profile and, in these cases, the neoclassical trapped-particle pinch seems to be too small by a factor 2–3 to account

Fig. 2-10 Trapped particle pinch velocity, scaled to the banana regime value given in Eq. (2-9), as a function of ν_{*e} for the case $r/R = 0.2$ (from Ref. 2-81).

quantitatively for the observations. In addition, some enhancement over the neoclassical value is required to explain data in the largest of present-day tokamaks, or for high density plasmas in smaller tokamaks.

2.2.1.3 Electrical Resistivity — In a uniform straight magnetic field, the electrical resistivity of a hydrogen plasma along the field takes the form

$$\begin{aligned} \eta_{||} &= 0.51, \frac{m_e \nu_e}{n_e e^2} \\ &= \frac{2.8 \times 10^{-8}}{[T_e (\text{keV})]^{3/2}} \text{ ohm-cm} \end{aligned} \quad (2-10)$$

In a plasma containing both hydrogen and impurity ions, the resistivity also becomes a function of the quantity

$$Z_{\text{eff}} = \frac{1}{n_e} \sum_{\text{ions}} n_i Z_i^2 = 1 + \frac{1}{n_e} \sum_{\text{impurities}} n_I Z_I (Z_I - 1) \quad (2-11)$$

As will be exhibited below, the resistivity increases approximately linearly with increasing Z_{eff} , with a correction factor arising from the diminishing role of electron-electron collisions as Z_{eff} increases.

In a tokamak plasma in the "collisionless" regime of neoclassical theory, only the untrapped particles can carry the current. Also, there is a friction or momentum exchange between the untrapped (current-carrying) particles and the trapped particles. These factors cause the resistivity in a tokamak to be increased by a correction term proportional to the number of trapped particles. In the collisionless regime, this correction term can be quite large. However, in the plateau and collisional regime ($\nu_{*e} \gtrsim 1$), the correction becomes smaller.

When trapping corrections and impurities are considered together, the situation is quite complicated. Clearly, impurities tend to increase the resistivity, but they also decrease the trapping correction, both by increasing the effective collisionality, and by decreasing the importance of friction between the current-carrying electrons and trapped electrons, compared with the friction due to the ions. A useful formula has been given (Ref. 2-83) that takes all these effects into account:

$$\eta_{||} = \frac{m_e \nu_e}{n_e e^2} G(Z_{eff}) \frac{1}{\left(1 - \frac{f_T}{1 + \xi \nu_{*e}}\right) \left(1 - \frac{C_R f_T}{1 + \xi \nu_{*e}}\right)} \quad (2-12)$$

where

$$G(Z_{eff}) = \frac{Z_{eff} (2.67 + Z_{eff})}{3.4 (1.13 + Z_{eff})}$$

and f_T is the fraction of trapped particles

$$f_T = f_T(\epsilon) = 1 - \frac{(1 - \epsilon)^2}{(1 - \epsilon^2)^{1/2}} \frac{1}{(1 + 1.46 \epsilon^{1/2})}$$

and

$$\epsilon = r/R$$

$$\xi = \xi(Z_{eff}) = 0.58 + 0.20 Z_{eff}$$

$$C_R = C_R(Z_{eff}) = \frac{0.56 (3.0 - Z_{eff})}{Z_{eff} (3.0 + Z_{eff})}$$

2.2.1.4 Impurity Transport. — Tokamak plasmas are often plagued with high-Z impurity ions (see Section 4.2 for a discussion of impurity control). Although the relative density of such ions is low (less than 1% of heavy metal impurities such as iron, tungsten, etc., and not more than 5–10% of lighter ions such as oxygen and carbon), they can contribute significantly to Z_{eff} and to the energy loss by radiation, and they would deplete the reacting fuel in a tokamak reactor.

When more than one ion species is present, the neoclassical theory must be modified in all three regimes, because of the friction between the various different ionic components. It is well known that, in classical theory, the ionic component with the higher Z feels a frictional force that causes it to concentrate near the center of the discharge, or, more precisely, near the point where the density of the lower- Z component (hydrogen) is maximum. This tendency persists in neoclassical theory and, moreover, the rate of inward impurity transport tends to be greater. Indeed, since ion-impurity collisions are responsible for impurity transport, one might expect the relevant diffusion coefficient to be comparable to the ion thermal diffusivity, in which case the diffusion

of high-Z ions to the plasma center would be quite rapid, and should certainly be observable on present-day experimental discharge time-scales. Impurity concentration at the plasma center would, of course, be quite deleterious to the achievement of ignition in a tokamak, and would also tend to prevent removal of the helium "ash" from a tokamak reactor.

A great deal of theoretical work has been done on neoclassical impurity transport in recent years, and the picture that emerges is quite complicated, and somewhat different from what one would naively expect.

Consider a plasma composed of hydrogenic ions, denoted i , and impurity ions, denoted l . The collision frequencies are given by

$$\nu_i = 4\pi^{1/2} (n_i + n_l Z_l^2) e^4 \ln \Lambda / 3m_i^{1/2} (kT_i)^{3/2} \quad (2-13)$$

$$\nu_l = 4\pi^{1/2} n_l Z_l^4 e^4 \ln \Lambda / 3m_l^{1/2} (kT_l)^{3/2}$$

We may define collisionality parameters

$$\nu_{*i} = \left(\frac{R}{r}\right)^{3/2} \frac{\nu_i q R}{v_{th,i}} ; \quad \nu_{*l} = \left(\frac{R}{r}\right)^{3/2} \frac{\nu_l q R}{v_{th,l}} \quad (2-14)$$

where $v_{th} = (kT/m)^{1/2}$. Ion temperature equilibration is usually sufficiently rapid that $T_i = T_l$ to a good approximation. Clearly, then, in the case $n_l Z_l^2 > n_i$,

$$\nu_{*l} / \nu_{*i} \approx Z_l^2 \gg 1 \quad (2-15)$$

Thus, in general, the impurities are in a higher collisionality regime than the hydrogen ions.

This feature makes the impurity transport problem qualitatively different from the electron/ion neoclassical transport problem. Roughly speaking, in any neoclassical theory, the radial flux of a given species is related to the parallel frictional force on that species due to its parallel flow relative to the other species. The strong collisional coupling of the parallel impurity flow to the parallel hydrogen flow reduces the parallel frictional force on the hydrogen ions. The result is that the magnitude of the radial hydrogen flux due to collisions with impurities (and, by ambipolarity, the magnitude of the impurity flux) is reduced to a value comparable to that which would pertain if the hydrogen were in the higher collisionality regime characteristic of the impurities.

For example, even if the hydrogen is in the banana regime ($\nu_{*i} < 1$), plateau-like scaling applies for the impurity diffusion if $\nu_{*l} > n_l m_l^{1/2} / n_i m_i^{1/2}$ (Ref. 2-84). As a second example, if the hydrogen is in the banana regime, collisional (Pfirsch-Schlüter) scaling applies for the impurity diffusion if $\nu_{*l} > (R/r)^{3/2}$. This seems to be the most relevant case in practice, over most of the discharge profile, both in present-day and future tokamaks. For this case, the result (Ref. 2-85) for the radial impurity flux is

$$(nv_r)_l = \sqrt{2} q^2 \nu_i \rho_i^2 \left(Z_l n_l \frac{\partial n_i}{\partial r} - n_i \frac{\partial n_l}{\partial r} \right) \quad (2-16)$$

where Eq. (2-7) is to be used for ν_i , and where we have neglected ion temperature gradients.

Fig. 2-11 Schematic illustration of the dependence of the neoclassical impurity diffusion coefficient on ν_{*i} for the case $n_I Z_I^2 > n_i$. The applicable diffusion coefficient is the lower of the two curves, usually the impurity curve for cases of practical experimental interest.

The behavior of the impurity diffusion coefficient is schematically illustrated in Fig. 2-11, for the case $n_I Z_I^2 > n_i$. The applicable diffusion coefficient is the lower of the two curves. Thus, within the range of experimental interest, the time scale for inward impurity diffusion is always longer than the time scale of neoclassical ion energy confinement. (For the ion thermal conductivity, it is of course the collisionality of the hydrogenic ions that matters: if the hydrogen is in the banana regime, then the thermal conductivity has a banana-regime value, even if the impurities are in the collisional regime.)

Ion temperature gradients have a significant effect on neoclassical impurity diffusion. Indeed, in some regimes the terms in the radial impurity flux arising from ion temperature gradients correspond to an outward transport of impurities. This effect, called "temperature gradient screening", could well be important in impeding impurity flow to the center of the discharge.

2.2.1.5 Field Ripple Effects — The preceding discussion has assumed implicitly that the magnetic field lines are closed within the plasma confinement region, and that the toroidal magnetic field is homogeneous in the toroidal direction, i.e., the system is axisymmetric. In a tokamak, deviations from axisymmetry will arise from the discrete nature of the external coils producing the toroidal field, leading to a fairly short wavelength modulation or "ripple" in the field strength as a field line is followed around the torus.

Clearly, particles with very small values of $v_{||}$ can be trapped in the troughs of these "ripples". Such particles will simply drift vertically with speed v_D given by Eq. (2-4), until they are collisionally scattered to larger $v_{||}$ -values. If the magnitude of the field ripple is given by $\Delta B/B \sim \delta$, then the ripple-trapping region of velocity-space will be $v_{||}/v \sim \delta^{1/2}$. Thus, the effective collision frequency for scattering out of the ripple-trapping region will be $\nu_{\text{eff}} \sim \nu/\delta$. Noting that the fraction of ripple-trapped particles will be of order $\delta^{1/2}$, the diffusion coefficient for this process will be of order $D \sim \delta^{1/2} \nu_{\text{eff}} (v_D/\nu_{\text{eff}})^2 \sim \delta^{3/2} v_D^2/\nu$, an expression that varies inversely with collision frequency. As $\nu \rightarrow 0$, however, a regime will be reached in which a ripple-trapped particle drifts completely out of the system before scattering, i.e., $v_D/\nu_{\text{eff}} > r$. The above formula for D should not be applied in such a regime; rather, the confinement time will be simply the time for collisional scattering into the loss region. However, this very collisionless regime is not important in practice, at least for thermal particles.

Since the diffusion rate varies inversely with collision frequency, thermal losses due to ripple-trapping will be much larger for ions than for electrons. The main effect of the ripple will be to increase the neoclassical ion thermal conductivity χ_i and also the losses of high-energy ions. Detailed calculations of the ripple contribution to χ_i have been given by Stringer (*Ref. 2-86*) and others. Their result for the thermal diffusivity is

$$\frac{\chi_i}{n_i} \approx 33 \frac{\delta^{3/2} G}{\nu_i} \left(\frac{ckT_i}{eBR} \right)^2 \quad (2-17)$$

where ν_i is given in Eq. (2-7) and 2δ is the maximum peak-to-peak ripple on the magnetic surface of given r . The function G depends both on the parameter $\alpha = r/RNq\delta$, N being the number of toroidal field coils, and on the θ -dependence of the ripple magnitude (*Ref. 2-87*). Typically, the value of G is about 3×10^{-2} for an α -value of 0.5.

The ripple-trapping contribution to ion thermal conductivity given in Eq. (2-17) does not seem to pose a very severe threat to near term high-temperature tokamak experiments. For example, we might consider the precompression plasma in TFTR which has a fairly large field ripple (about 5% peak-to-peak at the plasma edge) since the plasma is situated in the outer major radius portion of the TFTR vacuum vessel. For this case, with $T_i(0) \approx 5$ keV, $n(0) \approx 7 \times 10^{13} \text{ cm}^{-3}$, $a/R \approx 0.2$, $N = 20$, we find that the ripple contribution to χ_i is, on the average, about three times the neoclassical χ_i . For somewhat smaller ripple amplitudes, e.g., TFTR plasmas centered at smaller R -values, the ripple contribution to χ_i is comparable with, or smaller than, the neoclassical χ_i . However, the trend in the design of tokamak reactors has been toward smaller N and larger δ . In this case, ripple transport may well dominate the energetics of burning plasmas. In fact, the strong temperature dependence of χ_i has led to the suggestion of ripple as a mechanism to control the operating point of burning plasmas (*Ref. 2-88*).

Particles that are trapped in the toroidal field, and are following the usual "banana" orbits, experience small radial drifts in the presence of field ripple, due to the fact that the vertical drifts v_D no longer cancel exactly when averaged over the banana orbit. This effect gives rise to "banana drift" diffusion (*Ref. 2-89*), but this is typically no larger than the diffusion due to ripple-trapping. Finally, a new effect of field ripple has recently been identified, involving small-angle collisional scattering of toroidally trapped particles near their turning points, so that they are reflected at

different phases of the ripple on successive bounces. This process results in a thermal diffusivity independent of collision frequency (*Ref. 2-90*):

$$\frac{x_i}{n_i} \approx \frac{4}{a^2} \frac{R}{N v_{th,i}} \left(\frac{ckT_i}{eBR} \right)^2 \quad (2-18)$$

In many practical cases, this mechanism seems to dominate the ripple-trapping mechanism.

2.2.2 Anomalous Transport

As summarized in Section 2.2.3, it is quite clear from a wealth of experimental data that electron transport is greatly enhanced over the neoclassical prediction. This has led to intense theoretical investigation into those "anomalous" processes that may govern electron confinement.

The main sources of free energy that drive plasma microinstabilities are

- a. **Velocity-space Anisotropy.** However, velocity distributions in tokamaks are essentially Maxwellian, except for the longitudinal current that requires some electron streaming velocity along the field. In large tokamaks, this streaming velocity is less than the sound speed, in which case current-driven instabilities are predicted to be absent, or at least very weak. Thus, we will neglect this free-energy source.
- b. **Plasma Expansion Energy.** This arises whenever there are density or temperature gradients, and is the main source of free-energy driving tokamak microinstabilities.
- c. **Magnetic Energy.** Energetically, this is available to drive only the long-wavelength macroscopic kink/tearing modes discussed in Section 2.1.

Historically, anomalous particle and energy transport in tokamaks has usually been attributed to low-frequency microinstabilities, i.e., drift-waves, trapped-electron modes, and the like (*Ref. 2-91*). For such modes, the $\vec{E} \times \vec{B}$ drifts arising from the fluctuating $\vec{E}$ -fields provide the most obvious physical mechanism of cross-field transport, although more recently it has been realized that the fluctuating magnetic fields may also play an important role (*Ref. 2-92*). In any case, the central theoretical problem will be to identify the processes of nonlinear saturation of microinstabilities of this type, so that the steady-state spectrum of fluctuations can be determined. At the present time, this whole area seems to be one of very substantial uncertainty, with little consensus on the correct theoretical approach, and with no convincing agreement between theory and the experimental results. The experimentally observed fluctuation spectra do, however, have many features in common with theoretical predictions for drift-wave-type instabilities in toroidal geometries.

2.2.2.1 Drift Waves — If an electrostatic perturbation is taken to be the most dangerous type of mode for a low- β plasma then, neglecting resistivity and electron inertia, the parallel Ohm's law may be integrated to give

$$n_{1e} = \frac{e\phi}{kT_e} n_0 \quad (2-19)$$

where n_{1e} denotes the perturbed electron density, and it has been assumed that the electron temperature is uniform along the field lines. If $\omega > k_{\parallel} v_{th,i}$ there will not be time for sound waves

to produce equilibrium along the field. Hence with their large inertia, the ions are still dominated by the $\vec{E} \times \vec{B}$ drifts and obey the equation

$$\frac{\partial n_{1i}}{\partial t} + \vec{\nabla} \cdot \left(n_0 \frac{\vec{E} \times \vec{B}_0}{B_0^2} \right) = 0, \quad (2-20)$$

with $\vec{E} = -\vec{\nabla}\phi$. For highly sheared magnetic fields, the assumption $\omega > k_{\parallel} v_{th,i}$, i.e., the neglect of sound wave propagation along the field, may not be entirely justified.

Microinstabilities are relatively localized phenomena that, roughly speaking, require only that the plasma be inhomogeneous in space. Thus, the geometrical details of the confinement configuration are commonly quite unimportant. Accordingly, microinstabilities have usually been analyzed in the simplest possible configuration, namely, a plane slab of plasma nonuniform in the x-direction with a main magnetic field in the z-direction. A small shear field $B_y(x)$ is optional. (To relate this to tokamak geometry, identify x with the minor radius r, y with the poloidal angle θ , and z with the toroidal angle ϕ .) Because perturbations take the form $f(x) \exp(-i\omega t + ik_y y + ik_z z)$, it follows from Eq. (2-20) that

$$n_{1i} = - \frac{k_y \phi}{\omega B} \frac{dn_0}{dx} \quad (2-21)$$

and quasineutrality further gives

$$\omega = \omega_{*e} = - \frac{k_y k T_e}{eB} \frac{d \ln n_0}{dx} \quad (2-22)$$

If $\omega < k_{\parallel} v_{th,i}$ then the ions would also have a response of the form of Eq. (2-19), and the mode could not exist. Noting that $k_{\parallel} \approx k_z + k_y B_y/B$, it is seen that such modes are only possible for small k_z and for small values of the magnetic shear B_y/B , i.e., the modes take the form of long cells extended along the field lines.

The quantity ω_* is known as the diamagnetic drift frequency. The phase velocity of the wave is lower than those of MHD waves, being given by $\omega_*/k_y \sim v_{th,i} \rho_i (d \ln n_0/dx) \ll v_{th,i}$, where ρ_i is the ion gyroradius. Microinstabilities with $\omega \sim \omega_*$ can occur in many variants: drift-waves, trapped-electron, and trapped-ion modes, etc. Because of the requirement that modes be highly elongated along the field lines, magnetic shear may often be an important stabilizing influence.

A more detailed treatment of drift waves requires a kinetic treatment of the electrons. If it is assumed that $\omega \ll k_{\parallel} v_{th,e}$, Eq. (2-19) is recovered. However, electrons moving slowly along the field, specifically those with $v_{\parallel} \approx \omega/k_{\parallel}$, can resonate with the wave. The contribution from these electrons cannot be neglected, since it leads to an imaginary term in the dispersion relation for ω . This resonant electron contribution leads to instability of "electron" drift waves in three cases:

1. For density gradient only (i.e., no temperature gradient) provided terms of order $k_{\perp}^2 \rho_i$ are retained (the "universal" instability).
2. In the unusual case where $\eta_e = d \ln T_e / d \ln n < 0$.
3. Where the electrons have a net streaming speed u along the field (corresponding to a parallel current) that exceeds a threshold value $u_{crit} = (2\eta_e)^{1/2} v_S$ where $v_S = (kT_e/m_i)^{1/2}$.

So far, the discussion has been based on a "local approximation" in which the dispersion relation for ω is solved at each magnetic surface x , without attempting to construct an actual eigenfunction in x . This local approximation fails if there is significant magnetic shear. In a sheared field, $k_{\parallel}$ becomes a function of x , given by $k_{\parallel}(x) = k_z + k_y B_y(x)/B$. For appropriate values of k_y and k_z , there will be some surface x_s , the familiar "singular surface", where $k_{\parallel} = 0$. In the neighborhood of the singular surface, $k_{\parallel}(x) = k_y(x - x_s)/L_s$ where L_s , given by $L_s^{-1} = B_y'/B$, is called the "shear length". For a mode with $\omega \approx \omega_*$, the assumption that $\omega \gg k_{\parallel} v_{th,i}$ requires that the mode be localized around a singular surface within a distance $\Delta x \ll \omega_* L_s / k_y v_{th,i} = \rho_i L_s / L_n$, where $L_n^{-1} = n'/n$. Unless an unstable eigenfunction can be constructed that satisfies this localization condition, there will generally be a convection of wave energy away from the singular surface, which will tend to damp the mode. A further consequence of $k_{\parallel}$ becoming a function of x is that the assumption $\omega \ll k_{\parallel} v_{th,e}$ will fail in an even narrower region around the singular surface, specifically $\Delta x \sim \rho_e L_s / L_n$.

It has been known for some time that drift waves should be stable if the shear exceeds a certain critical value (*Ref. 2-93*). However, exact eigenfunctions in x have only recently been obtained and it has become clear that the drift waves discussed above are in fact all stable even at small values of the shear, at least in this simple geometry (*Ref. 2-94*).

At shorter mean free path, there is a set of "collisional" drift waves that resemble the "collisionless" modes discussed above. However, in a slab geometry, these are also stabilized by a minimal amount of magnetic shear (*Ref. 2-95*).

In toroidal geometry, this picture is radically altered. Besides trapped particle effects, discussed later, toroidal geometry introduces the drift v_D due to the inhomogeneous magnetic field, and this has the effect of coupling various poloidal harmonics of the mode in such a way as to create a "ballooning" type of structure in which the mode amplitude is quite small on the inside major radius side of the torus (*Ref. 2-96*). In this case, the periodicity condition in θ is much less important, with the result that the stabilizing effect of shear is greatly reduced. Detailed calculations of this effect have been made by means of two-dimensional numerical computations of the eigenmodes (*Ref. 2-97*). For values of the magnetic shear typical of tokamaks, it has been found that drift-type modes remain unstable over a wide range of collisionalities.

An additional class of drift waves, referred to as "ion-drift waves", arises from the assumption $\omega \sim k_{\parallel} v_{th,i}$. The effect of magnetic shear will be much less important for these modes, since they do not need to be as sharply localized about their singular magnetic surface as do the electron drift waves. Indeed, it seems that the "local approximation" is quite adequate for these modes, even in sheared fields. However, these modes are usually stable, in fact quite strongly damped by "resonant ion" effects. The only important exception is the case where the ion temperature gradient is unusually strong (*Ref. 2-91*), specifically if $\eta_i = d \ln T_i / d \ln n > 2$.

2.2.2.2 Trapped-Electron and Trapped-Ion Modes — In obtaining the response of particles to oscillating perturbed electric fields, trapping will also be inoperative in the case where the oscillation frequency ω satisfies $\omega \tau_b \gg 1$, where τ_b is the "bounce time". It is clear that the "plane slab" model provides an adequate basis for treating microinstabilities in tokamaks if either of two conditions are satisfied: namely, if $\nu_* \gg 1$ or if $\omega \tau_b \gg 1$. This section deals with some instabilities that arise when neither of these conditions is satisfied.

Trapped-electron modes arise in the case where $\nu_{*e} \ll 1$ and $\omega \tau_{bi} \gg 1$, so that trapping is important for electrons but not for ions. These modes also require that $k_{\parallel} v_{th,i} \ll \omega \ll k_{\parallel} v_{th,e}$ so that the response of the untrapped electrons and ions to the perturbed fields is the same as in the drift-wave case.

To obtain the perturbed distribution function for trapped electrons, an average is taken over the bounce orbit of the particle. A collision term representing the scattering of trapped particles out of the trapping region must also be included; the simple model $C(f) = -\nu_{\text{eff}} f$ will suffice for this purpose. Linearizing about a Maxwellian equilibrium distribution f_0 gives the expression

$$\omega = \omega_{*e} + \eta_e \omega_{*e}^2 \int_T \frac{\left(\frac{3}{2} - \frac{v^2}{2v_{\text{th}}^2}\right) f_{0e}}{\omega - \omega_D + i\nu_{\text{eff}}} d^3\vec{v} \quad (2-23)$$

where $\omega_D = \vec{k} \cdot \vec{v}_D$ represents the effect of the precessional drift of the bananas.

The stability of the trapped-electron modes hinges upon the sign of the imaginary part of the integral term in Eq. (2-23). Two cases arise: the case $\omega, \omega_D \ll \nu_{\text{eff}}$ where the imaginary contribution arises from the collisional term; and the case $\omega, \omega_D \gg \nu_{\text{eff}}$ where the imaginary contribution arises from the resonance $\omega \approx \omega_D$ (Ref. 2-98). Noting that $\nu_{\text{eff}} \propto v^{-3}$ and that $\omega_D/\omega_* \sim (v^2/v_{\text{th}}^2) (r/R)$, it can easily be seen that electrons with velocities somewhat larger than the thermal velocity make the largest contribution in both cases. Thus, in both cases, there is a positive imaginary term.

We are now in a position to list the important features of trapped-electron modes:

- Instabilities arise mainly from electron temperature gradients. (Although not demonstrated here, very short wavelength modes, $k_{\perp} \rho_i \sim 1$, can be unstable even without a temperature gradient.)
- Growth rates vary inversely with ν_{eff} at high collisionality and are roughly independent of ν_{eff} at low collisionality.
- Unstable modes tend to have fairly short perpendicular wavelengths [$k_{\perp} \rho_i > (r/R)^{1/2} r/Rq$] in order to satisfy the condition $\omega_* \tau_{bi} \gg 1$.

In addition, it is found that shear stabilization is quite ineffective for trapped-electron modes, partly because of the relatively high growth rates that can occur, and partly because the mode amplitude is small on the inside of the torus.

Trapped-ion modes arise in the case where $\nu_{*i,e} \ll 1$ and $\omega \tau_{bi} \ll 1$, so that trapping is important for both electrons and ions. These modes also require that $\omega \ll k_{\parallel} v_{\text{th},i} \ll k_{\parallel} v_{\text{th},e}$, so that the contribution from untrapped particles is unimportant both for electrons and ions.

There are several distinct asymptotic regimes for trapped-ion modes. The most relevant regime is where $\nu_{\text{eff},i} \ll \omega \ll \nu_{\text{eff},e}$ and $\omega \gg \omega_{Di,e}$. Estimating the trapped region to be a fraction $\epsilon^{1/2} = (r/R)^{1/2}$ of velocity space and writing $\nu_{\text{eff}} = \nu/\epsilon$ gives

$$\omega = \frac{1}{2} \epsilon^{1/2} \omega_{*e} + i \left(\frac{1}{4} \frac{\epsilon^2 \omega_{*e}^2}{\nu_e} - \frac{\nu_i}{\epsilon} \right) \quad (2-24)$$

It is seen that the mode is unstable if $\omega_*^2 > 4\nu_e\nu_i/\epsilon^3$, i.e., for all sufficiently short perpendicular wavelengths. However, at very short wavelengths the assumption $\omega\tau_{bi} \ll 1$ tends to be violated. In fact, for $\omega\tau_{bi} \lesssim 1$ there are additional stabilizing effects from resonances between the wave and harmonics of the "bounce frequency" τ_b^{-1} of the trapped ions, especially the low-energy ions that have relatively long τ_b . A more complete theory including these resonance effects (*Ref. 2-99*) shows that there is a threshold in collisionality above which the trapped-ion modes are stable. Although this threshold depends on many factors, for the case $T_e = T_i$ and $\eta_e = \eta_i = 0.5$ the condition for instability can be roughly expressed (*Ref. 2-100*) as $n < n_{crit}$ where

$$n_{crit} (\text{cm}^{-3}) = 6 \cdot 10^{16} [r(\text{cm})^{3/2} T(\text{keV})^2] / [R(\text{cm})^{5/2} q].$$

However, at larger η_i the bounce resonances become destabilizing and this stability condition no longer applies. The high-temperature experiments on PLT reported in 1978 fall somewhat into the unstable trapped-ion regime, according to this criterion; furthermore, the ion temperature gradient was quite large in these experiments ($\eta_i \sim 2 - 4$), in which case instabilities are also predicted to occur at higher collisionality.

2.2.2.3 Anomalous Transport From Microinstabilities — Such microinstabilities result in the degradation of plasma confinement. In order to understand the nature of such anomalous transport, it is necessary to consider the nonlinear behavior of low frequency microinstabilities. Theoretical approaches to this problem fall into three categories: phenomenological models, weak turbulence theories, and strong turbulence theories.

The phenomenological model that has been employed most often is due to Kadomtsev and Pogutse (*Ref. 2-91*), who argued that the anomalous cross-field diffusion coefficient should be given roughly by the formula

$$D = \gamma/k_{\perp}^2 \quad (2-25)$$

where γ is the linear growth rate and $k_{\perp}$ is the wave number perpendicular to the magnetic field. The argument given in support of this model is as follows. For drift waves, such as trapped-electron modes and trapped-ion modes, it is generally true that the growth rate γ is an increasing function of perpendicular wave-number $k_{\perp}$. Thus, small-scale perturbations grow first, and their amplitudes change on a shorter time scale than the time scale for evolution of the large-scale perturbations. Furthermore, the large-scale perturbations will appear as an additional inhomogeneity in the plasma, hence acting as an additional source for driving the small-scale perturbations. Thus, there is a flow of wave energy toward larger $k_{\perp}$ -values. At the largest $k_{\perp}$ -values, usually corresponding to perpendicular wavelengths somewhat shorter than the ion gyroradius, various types of wave damping by ions become important, e.g., damping by collisional viscosity, or by resonances between the wave and the motion of ions along the field.

In the phenomenological approach, the details of the fluctuation spectrum are ignored and the small-scale perturbations are regarded merely as a damping mechanism for the larger-scale perturbations. Assuming that this damping takes the form of a diffusion, the corresponding damping decrement is balanced against the linear growth rate, leading to Eq. (2-25). In evaluating this formula, the $k_{\perp}$ -value corresponding to the largest unstable wavelength should normally be used; this usually gives the largest value of the expression $\gamma/k_{\perp}^2$ considered as a function of $k_{\perp}$. The maximum scale length of unstable perturbations will generally be determined by the radial localization width of the mode in the sheared magnetic field. In employing Eq. (2-25), it is also customary to invoke the chaotic nature of the turbulent motion to justify the assumption of an isotropic perpendicular spectrum, i.e., $k_x^2 \sim k_y^2$ in the plane slab approximation.

This approach has the advantage that it leads to fairly simple expressions for the various transport coefficients. Roughly, the diffusion coefficient D and the electron thermal diffusivity χ_e scale with collision frequency in the manner indicated in Fig. 2-12. With the possible exception of the lowest collisionality, beam heated PLT discharges, present-day tokamak experiments fall either into the regime where transport due to the trapped-electron mode should dominate, or into the more collisional regime. In the trapped-electron and trapped-ion regimes, this theory clearly indicates that the transport coefficients should increase with increasing temperature. A mitigating factor is that the underlying microinstabilities are driven by plasma inhomogeneities. Thus, at the same density, temperature, and magnetic field, the anomalous transport coefficients will decrease relative to the classical transport coefficients as the device size increases.

Fig. 2-12 Schematic illustration of the dependence of the diffusion coefficient D and electron thermal conductivity χ_e on electron collision frequency ν_e for anomalous transport due to microinstabilities.

As discussed in Section 2.2.4, which deals with the comparison between theory and experiment, empirical scaling laws for the energy confinement time exhibit more favorable dependence on density and temperature than do computations based upon these particular models. Accordingly, in order to explain the experimental observations, either improvements must be made in these theoretical models of low-frequency turbulence or it must be concluded that entirely different physical processes are responsible for the observed cross-field transport of electron energy.

More systematic models of the unstable wave spectrum can be attempted using the techniques of "weak turbulence theory" (Ref. 2-101). Here, it is assumed that $\gamma/\omega \ll 1$ and that the energy in the excited modes is small compared with the plasma thermal energy, so that linear theory provides a reasonable description of the structure of the unstable modes. The turbulent spectrum can then be decomposed into a sum of linear modes which will couple with each other in various ways, giving rise to a weakly turbulent state. Weak turbulence theories of this type have not generally been very successful when applied to drift waves, trapped-electron modes, etc. However, some success has been recently achieved in comparing the spectrum of waves so obtained with the experimental observations of density fluctuations (Ref. 2-102).

"Strong turbulence theories" in which the linear theory no longer provides even an approximate description of the excited modes, are obviously extremely difficult. One approach (Ref. 2-103) is to look at the modifications of the ion orbits due to the turbulent $\vec{E} \times \vec{B}$ drifts. At a certain level of turbulence, it is found that the ion orbits become quite stochastic, leading to a damping of the waves by a kind of turbulent viscosity. To the extent that this has been worked out, it leads back

to the formula $D \sim \gamma/k_{\perp}^2$ and may ultimately provide a more convincing justification for use of the elementary approach.

2.2.2.4 Transport Due to Magnetic Perturbations — As was shown in Section 2.1.2, $m = 2$ (and sometimes $m = 3$) tearing modes can give rise to quasi-static, large-scale magnetic islands in tokamaks (see, in particular, Fig. 2-4).

Since plasma transport is many orders of magnitude more rapid along magnetic field lines than perpendicular to them, most plasma properties are very nearly constant along magnetic field lines. Thus, for a single magnetic island or a discrete set of well separated magnetic islands, the plasma temperature and other plasma parameters should be nearly constant along helical flux contours. Microscopic plasma transport is not directly enhanced by the presence of a magnetic island [except for possible effects on the orbits of high-energy trapped ions due to the non-axisymmetry of the magnetic field (*Ref. 2-104*)]. Rather, because flux surfaces on opposite sides of an island nearly touch at the x-points of the separatrix, the island just provides a transport "short circuit" over the physical region encompassed by the island. When magnetic islands of different helicity overlap to such an extent that the magnetic field lines become stochastic, plasma transport becomes considerably more complex, as described in Section 2.2.2.4.5.

2.2.2.4.1 Large-Scale Islands and Sawtooth Oscillations — The $m = 1$ tearing mode does not exhibit a quasi-static phase, and is responsible for a substantial thermal transport from the central region of any tokamak in which the q value on axis drops below unity, through the mechanism of internal disruptions, or sawtooth oscillations. The $m = 1$ mode begins to grow whenever q drops below unity at the plasma center, thereby inducing shear at the $q = 1$ surface r_s . Single helicity numerical calculations (*Ref. 2-105*) have shown that the mode continues to grow until its magnetic island becomes so large that the magnetic axis of the island moves toward the geometric center of the plasma and thereby transforms the magnetic topology of the central core from the original one centered about the central magnetic axis to one centered about the magnetic island.

Since the $m = 1$ tearing mode growth rate is much faster than the ohmic heating rate, the electron temperature is equilibrated along all of the reconnecting magnetic field lines and hence is flattened within the induced $m = 1$ magnetic island. Ultimately, the region over which the electron temperature flattens extends, from flux conservation considerations, out to $r_0 \sim \sqrt{2} r_s$. Thus, immediately after the internal disruption, the temperature profile is flat within the radius r_0 . Thereafter, on a much slower time scale, the combination of ohmic or auxiliary heating and radial electron heat transport reheats the plasma center back toward the centrally peaked electron temperature and current density profiles. An analytic model of internal disruptions has been developed (*Ref. 2-10*), based on the cyclic process in which the plasma core resistively overheats and drives q below unity, causing the $m = 1$ tearing mode to become unstable with an accelerating growth rate and ultimately to flatten the electron temperature profile as magnetic reconnection occurs. This model has been extensively and quite favorably compared with ORMAK data on the space-time evolution of internal disruptions and their $m = 1$ precursors (*Ref. 2-10*).

2.2.2.4.2 Skin-Current Penetration — For hollow electron temperature profiles such as can be produced during the current initiation phase of tokamak discharges, two singular surfaces of the same helicity (e.g., $m = 3$, $n = 1$) can occur within the plasma. Then, if tearing modes are unstable at both of these surfaces, two sets of magnetic islands develop during the nonlinear evolution phase. The final state is found (*Ref. 2-106*) to be characterized by one of two possibilities: (a) the two modes saturate separately in much the same manner as Mirnov oscillations; (b) the

two islands interact and the field lines reconnect into a new axisymmetric equilibrium for which the q profile is flat with $q \approx m/n$. The reconnection case is very similar to the evolution of $m = 1$ modes into internal disruptions and is probably the mechanism behind the increased MHD activity occurring during current initiation and with hollow electron temperature profiles, e.g., the $m = 3$ disruption in PLT. Although the simulations (*Ref. 2-106*) have concentrated on the case of single helicity "double" tearing modes of this type, in reality skin-current distributions are generally unstable to many different m -values, including fairly high m .

Transport-code simulations of the skin phase of tokamaks have been carried out (*Ref. 2-107*) using models for the skin relaxation based on magnetic reconnection due to double tearing modes of this type. These simulations have been quite successful in reproducing the experimentally observed behavior.

2.2.2.4.3 Small-Scale "Drift-Tearing" Modes — In recent years, considerable interest has developed in a type of drift wave that involves a significant perturbed magnetic field as well as the usual perturbed electric field. In particular, there is a class of modes of this type in which the radial component of the perturbed magnetic field is nonzero at the singular magnetic surface. Such modes will give rise to a structure of magnetic islands at the singular surface very much like tearing modes. For this reason the modes are usually called "drift-tearing modes".

The principal difference between drift-tearing modes and ordinary tearing modes is the source of free energy. For ordinary tearing modes, this is the magnetic free energy associated with the parallel currents in the plasma. For drift-tearing modes, expansion free energy also plays an essential role. In order to utilize magnetic free energy, ordinary tearing modes tend to have long perpendicular wavelengths, i.e., small m -numbers. On the other hand, drift-tearing modes (like other drift instabilities) tend to have short perpendicular wavelengths, i.e., large m -numbers. Indeed, in the theory of drift-tearing modes the terms that arise from the perturbation in the magnetic energy are usually stabilizing. However, in certain cases the destabilizing terms from expansion free energy can exceed these stabilizing terms.

The topic is highly complex and still under development so a thorough treatment will not be undertaken here. Perhaps the most important contributions to the literature are a fairly comprehensive theory of these instabilities for a slab or straight cylindrical configuration (*Ref. 2-108*), and the identification of important destabilizing (*Ref. 2-109*) and stabilizing (*Ref. 2-110*) modifications to the slab considerations.

The potential importance of these modes, as compared with electrostatic drift waves, lies in the fact that the magnetic field perturbations provide a mechanism of cross-field transport (i.e., particle streaming along perturbed field lines) that is qualitatively different from $\vec{E} \times \vec{B}$ drift mechanisms.

2.2.2.4.4 Magnetic Flutter Due to Drift Waves — It has been argued that electrostatic drift waves (and, indeed, also shear Alfvén modes) can be unstable in a tokamak plasma when toroidal geometry or trapped-electron effects are included. As an alternative to mechanisms based on $\vec{E} \times \vec{B}$ drifts, the magnetic component of both drift and shear-Alfvén microinstabilities (magnetic "flutter") provides a vehicle for anomalous transport by means of microscopic magnetic turbulence.

At very low amplitudes, microscopic magnetic perturbations produce non-overlapping magnetic islands in the same way that the macroscopic modes do. At this low level, if there is some extrinsic mechanism that determines the growth and decay rate γ for the microinstabilities, the

electron heat conduction coefficient induced by the birth and death of microscopic magnetic islands is estimated "quasilinearly" to be $\chi_e^w \sim (3/16) (\gamma + \nu) w^2$, where ν is the collision frequency and where w is the island width (Ref. 2-92). By analogy with the nonlinear evolution of the macroscopic modes discussed above, we expect the most significant transport to occur when two islands begin to overlap and thereby destroy the magnetic flux surfaces locally. While the induced transport has not been calculated in detail for this case, it is presumably similar in form to the quasilinear result with the island width replaced by the initial island separation. While there does not yet exist a complete theory that correctly predicts the scaling and magnitude of the anomalous electron heat conduction coefficient, the generic consequences of such "magnetic flutter" models are common features of tokamak discharges so this approach is an encouraging avenue to follow.

2.2.2.4.5 Stochastic Magnetic Fields — A prerequisite for satisfactory particle confinement is that the magnetic field lines themselves remain confined. However, tearing modes disturb the magnetic field on a macroscopic scale and, furthermore, there is a class of microinstabilities (drift-tearing modes) which can create fine-scale perturbations of the magnetic field. This has led to considerable recent interest in the general problem of transport due to stochastic magnetic fields (Ref. 2-111).

In a typical set of "magnetic islands" resulting from a single tearing mode in cylindrical geometry, the field lines are found to remain confined. This arises from the fact that the system still possesses symmetry (helical symmetry in the tokamak context). The situation with two chains of magnetic islands of different helicity, shown schematically in Fig. 2-13, is quite different. (Note that this figure is not an exact representation, in that the chains of islands cannot both be independent of z if the symmetry is to be broken.) If the island width is w and the island separation is d , an important parameter is the quantity $s = w/d$. If $s < 1$, the only consequence of the lack of exact symmetry is to produce a small region around each separatrix where the field lines become ergodic; these regions are shown dotted in Fig. 2-13. However, if $s > 1$ it turns out that the field in this entire region becomes ergodic (Ref. 2-112), i.e., if one follows a field line a large distance z , it will have diffused radially a distance Δr given by $\langle (\Delta r)^2 \rangle \approx 2D_M z$. A rough formula for the magnetic diffusion coefficient D_M is given by $D_M = \pi R B_1^2 / B_0^2$, where $\tilde{B}_1$ is the mean stochastic radial field. It has been pointed out (Ref. 2-113)

Fig. 2-13 The effect of two magnetic island chains of different helicity. The magnetic field in the dotted region around the two separatrices becomes ergodic. When $w/d \geq 1$, the entire region becomes ergodic.

that this should lead to a radial electron thermal diffusivity of magnitude $\chi_e/n_e \approx D_M v_{th,e}$. Very small turbulent magnetic fields ($B_1/B_0 \sim 10^{-4}$) could in this way account for the anomalous electron thermal transport observed in tokamaks. Such a level is reasonably consistent with what might be expected of microinstabilities. However, it should be emphasized that the theory of stochastic fields has thus far concerned itself only with the consequences of stochasticity produced by some *fixed* asymmetric current pattern. The key problem of the *self-consistent* determination of microinstability levels in a stochastic medium has yet to be addressed with any rigor.

A phenomenological approach (Refs. 2-114, 2-115), in which the cross-field correlation length for magnetic fluctuations is taken to be of order the collisionless skin depth c/ω_{pe} (where $\omega_{pe}^2 = 4\pi n_e^2/m_e$) has resulted in an expression for the anomalous electron thermal diffusivity, namely

$$\chi_e = \frac{c^2}{\omega_{pe}^2} \frac{v_{th,e}}{Rq} \quad (2-26)$$

that has a magnitude roughly consistent with experimental data on electron energy confinement.

2.2.3 Experimental Situation

By way of introduction to the experimental results of plasma confinement, some of the general features of tokamak discharges are briefly reviewed. Figure 2-14A displays the time evolution of the plasma current (I), loop voltage (V), and the line-averaged electron density ($\bar{n}_e$) in a typical PLT discharge with deuterium gas (Ref. 2-116), the data for $\bar{n}_e$ having been obtained from microwave interferometry. Similarly, Fig. 2-14B exhibits the measurements of the central [$n_e(0)$] and volume-averaged ($\langle n_e \rangle$) electron densities, the central [$T_e(0)$] and density-averaged ($\langle T_e \rangle$) electron temperatures, all of which are determined from Thomson scattering, and the central ion temperature [$T_i(0)$], the latter based on the absolute flux measurement of thermonuclear neutrons, a value which is generally in good agreement with that from the more customary measurement of charge exchange spectra. The evolution of the radial profiles of the electron temperature [$T_e(r)$] and density [$n_e(r)$] are illustrated in Fig. 2-15. The above-mentioned diagnostics on which the radial profiles are obtained are sometimes supplemented by: multichord microwave (or infrared laser) interferometry for $n_e(r)$; synchrotron and bremsstrahlung (soft X-ray) measurements for $T_e(r)$; and spectroscopic measurements of Doppler broadening of impurity lines for $T_i(r)$. A more complete catalog of diagnostic techniques is contained in Section 4.4.

The radial profile of the current density, $j(r)$, in the "steady state" is usually deduced from the conductivity measurement. Under the assumption that both the effective ionic charge (Z_{eff}) and the toroidal electric field (E_t) are independent of r , $j(r)$ and Z_{eff} are inferred from the measured $T_e(r)$ and $n_e(r)$ profiles and the total plasma resistance, with care being taken to account for the corrections to the resistivity due to trapped-particle effects (see Section 2.2.2). Basically $j(r)$ follows the distribution of $T_e^{3/2}(r)$ in the steady-state plasma. The inferred $j(r)$ is "spot-checked" favorably by comparing between the singular surface positions [$q(r) = m/n$] of the resultant $q(r)$ and those inferred from electron temperature fluctuation measurements of soft X-ray emission (Refs. 2-8, 2-117). Furthermore, several techniques for direct measurement of $j(r)$ have demonstrated feasibility by virtue of their general agreement with the conductivity-derived $j(r)$, although they have not as yet yielded enough data to document discrepancies. These techniques include measurement by heavy-ion (thallium) beam (Ref. 2-118), generation of harmonic microwaves (Ref. 2-119), angular-scattering of injected fast ions (Ref. 2-42), ruby-laser forward scattering

Fig. 2-14 Time-variation of current, voltage, density, electron temperature and ion temperature in a typical ohmically-heated discharge in PLT ($R_0 = 1.34 \text{ m}$, $a = 0.40 \text{ m}$, $B_T = 32 \text{ kG}$).

Fig. 2-15 Time-variation of electron temperature and electron density profiles. The case shown here is the same as that in Fig. 2-14.

(*Ref. 2-120*), and Zeeman effect on an injected neutral lithium beam (*Ref. 2-121*). Absolute ultraviolet spectroscopic measurements of impurity abundance in various tokamaks have demonstrated general agreement (especially at high density in clean plasmas) of their estimate of Z_{eff} with that inferred from the conductivity measurements (*Ref. 2-122*). This is further confirmed by absolute flux measurements of soft X-ray emission (*Ref. 2-123*).

Therefore the plasma resistivity parallel to the magnetic field is, within experimental error, classical in steady state plasmas under normal operating conditions (i.e., not at low density). The enhancement of resistance beyond the classical value for a hydrogen plasma is accounted for by the presence of impurities.

Outwardly peaked T_e -profiles are typically observed during the buildup phase of the discharge current (*Ref. 2-124*). This implies the presence of a skin effect, i.e., an outwardly peaked $E_T(r)$ and a peaking of $j(r)$ even stronger than $T_e^{3/2}(r)$, an interpretation supported by the values of internal inductance deduced from the plasma position measurement (*Ref. 2-7*). Values of Z_{eff} estimated from the resistance and including classical magnetic field diffusion are considerably larger than the spectroscopic measurements, indicating the presence of an enhanced current penetration mechanism (*Ref. 2-125*).

2.2.3.1 Particle Flow Pattern — Since discharges in large tokamak experiments are commonly sustained for many particle confinement times, the process of plasma maintenance by injection and ionization of neutral atoms plays a central role in determining the nature of the discharge.

The plasma electron density reached during the initial few milliseconds of the discharge tends to be proportional to the initial filling pressure. The subsequent density behavior without further deliberate introduction of neutral atoms is related to wall-evolved neutrals and ionization of impurities in addition to the particle confinement properties. However, under clean discharge conditions [especially with divertors (*Refs. 2-126, 2-127*) or gettered walls (*Ref. 2-128*)] contributions from impurity ionization is negligible, and the recycling rate off the walls and limiters tends to be small (*Ref. 2-129*), although the latter is sensitive to the wall cleaning methods employed (*Ref. 2-130*). Therefore the initial density decays quickly with the time constant of the gross particle confinement-time and subsequent plasma "fueling" becomes essential. Fueling is discussed in some detail in Section 4.4.2 and is dealt with here only to the extent needed to characterize its role in particle balance.

Although fueling by solid-hydrogen pellet injection (*Ref. 2-131*) and substantial neutral beam injection (*Ref. 2-25*) have been employed, the principal fueling method used in all present tokamaks has been "gas puffing", i.e., an injection of cold neutral gas into an already established discharge. In contrast with the former two methods, neutral atoms are initially deposited in the plasma edge in the latter technique. Therefore the subsequent transport of neutral atoms is more important in this method; furthermore, since it is a more localized source, the modifications to plasma dynamics may be more significant.

The magnitude and radial distribution of the neutral hydrogen atom density, $n_0(r)$, in the plasma is usually determined by spectroscopic and charge-exchange measurements, although uncertainties are relatively large (perhaps a factor of 2–3). The central neutral density, $n_0(0)$, is typically several orders of magnitude smaller than the edge value, and varies from 2×10^7 to $2 \times 10^9 \text{ cm}^{-3}$, with the higher values being obtained in discharges with lower plasma densities in a small device. The neutral density is peaked in the vicinity of the limiter, where up to half of the neutral influx takes place in typical cases (*Refs. 2-124, 2-132*).

Upon application of gas puffing, the central electron density increases rapidly and no significant inversion of the electron density profile is observed, contrary to expectation from balance

between ionization of neutrals and particle diffusion using conventional scaling laws. This circumstance, although not fully understood at present, is an encouraging sign for the ability to achieve the densities required for tokamak reactor operation. Indeed, if no anomalous "pinch" is operative, it can be shown on general grounds that high β operation is precluded unless a particle source in the plasma center is present (*Ref. 2-133*).

Gas puffing tends to reduce the temperature at the plasma edge, thus affecting the current density (and thereby gross MHD stability) and the plasma-limiter interaction. The choice of level and timing of gas puffing provides a means of controlling both Mirnov oscillations and the influx of impurities. However, the resultant current channel shrinkage eventually causes a disruptive instability. The maximum density attainable without disruptions, therefore, is related to the balance between the cooling and the power input. Considerations of this type have led to the suggestion that the maximum density obtainable in ohmically heated discharges scales as B_T/R and this is in rough agreement with experiment (*Ref. 2-31*) although Z_{eff} is also an important variable in this regard (see Fig. 2-5).

2.2.3.2 Energy Flow Pattern — The energy flow pattern of tokamak discharges is illustrated in Fig. 2-16. Ohmic heating power goes into the electrons, and is transmitted to the ions via Coulomb collisions on the time-scale $\tau_{ie}^{-1} = d \ln T_i / dt \propto n T_i^{-3/2} f(T_e/T_i)$, where f is roughly constant for $1.6 < T_e/T_i < 10$ (*Ref. 2-134*). Therefore, within this range of T_e/T_i , ion temperature is approximately independent of electron temperature, provided of course that τ_{ie} is not too short compared to the ion energy confinement time τ_{Ei} . The central ion temperature in most ohmically heated tokamaks shows good agreement with Artisimovich scaling (*Ref. 2-134*),

$$T_i(0) = 6 \times 10^{-7} \left(B_T I \bar{n}_e R_0^2 \right)^{1/3} \left[\text{eV}; \text{G, A, cm}^{-3}, \text{cm} \right] \quad (2-27)$$

Fig. 2-16 Energy flow pattern in the tokamak discharge.

which was derived on the basis of Galeev-Sagdeev plateau thermal conductivity (*Ref. 2-135*) and a constant f . As the plasma density becomes so high that $\tau_{ie} \ll \tau_{Ei}$, the electron and ion energy are closely equilibrated, and the Artsimovich scaling is no longer valid. For example, Alcator discharges with $\bar{n}_e \gg 10^{14} \text{ cm}^{-3}$ exhibit $T_i \approx T_e$ (*Ref. 2-37*). The energy flow may be quite different when supplementary heating is applied, but the specifics depend on the particular heating technique employed. High power ($> 2 \text{ MW}$) neutral beam injection at modest energies applied to PLT gives rise to a plasma with $T_i(0) \approx 6.5 \text{ keV}$ and $T_e(0) \approx 4 \text{ keV}$ (*Ref. 2-25*). The electron-ion heat transfer in this case is reversed from that with ohmic heating alone.

From the point of view of energy flow, it is convenient to regard the plasma as consisting of the interior region and the plasma edge region, as in Fig. 2-16. Electron energy is transported from the hot interior region of the plasma to the plasma edge principally by heat conduction and by the energy that accompanies mass transport. (In steady state, the latter quantity is equal to the energy consumed in heating the electrons of newly ionized neutral atoms to the ambient electron temperature.) Because plasma transport is many orders of magnitude faster along the magnetic field lines than across them, even small magnetic perturbations caused by tearing modes, in particular internal disruptions and Mirnov oscillations, can have significant effects on energy transport (*Ref. 2-21*). Although internal disruptions (*Ref. 2-8*) produce rapid energy transport from the central-core region ($r > r_1$) to the region exterior to r_1 , they do not constitute the dominant power loss in the interior region as a whole. The heat pulses resulting from internal disruptions are a valuable tool for experimentally examining energy transport processes in the region $r > r_1$ (*Refs. 2-136, 2-137*). Mirnov oscillations are known to cause significant energy transport from the interior region (*Refs. 2-33, 2-34*).

In some discharges containing heavy metal impurities (such as tungsten or molybdenum from limiters), the energy lost by radiation from the plasma interior may be a significant fraction of the input power. In pathological cases, severe radiation loss coupled with thermal instability in profiles of ohmic heating power density can produce the hollow T_e profiles observed in several tokamaks (*Refs. 2-138, 2-139, 2-140*). Similarly, the ion energy loss from the plasma interior is in general dominated by heat conduction and mass flow to the plasma edge, but charge exchange is of comparable magnitude at low density as an energy loss mechanism.

At the plasma edge, electrons and ions may actually be equilibrated more closely than in the interior, provided the density profile is flatter than the temperature profile, as is observed to be the case with intense gas puffing. Electron and ion energy can go from the plasma edge to the limiter by mass flow. Electron energy can also flow to the limiter by heat conduction, provided that the limiter emits enough cold secondary electrons. The other important mechanisms to remove the energy of the plasma edge are radiation and charge-exchange directly to the wall. The power falling on the wall (mostly through radiation) varies considerably (from $\sim 25\%$ to $\sim 100\%$ of the total input power) depending on the abundance and species of impurities in the plasma. In clean plasmas, hydrogen light emission dominates radiation loss. When the dominant impurities are low Z (typically oxygen, carbon, and nitrogen), the bulk of the emitted radiation originates from the plasma edge. However, when the plasma is contaminated by high Z impurities, significant radiation power comes from the interior region.

2.2.4 Transport Scaling Laws

The energy confinement time is an important practical measure of tokamak performance, and the possibility that it might scale with other experimental quantities in a predictable manner is of considerable interest both for interpreting results of existing experiments and for the design of

future devices. In fact, one can state that the goal of confinement experiments is to resolve the energy and particle loss processes into their constituent elements in order to deduce such dependences. Some of these losses occur in atomic and molecular physics channels and these can be reduced (in principle) by more effective device cleaning, by alternative magnetic configurations (divertors), by asymmetric gas feed (flow reversal) or by better plasma control; see Section 4.2 for a discussion of these approaches. On the other hand, some channels are an irreducible by-product of magnetic confinement, i.e., there must always be points of contact between the plasma and the external surface.

When plasma confinement losses can be clearly separated from radiative and charge exchange losses, theoretical models can be unequivocally tested against them. In most present confinement experiments, however, this clear separation has been effected only at the extreme ranges of performance. Lacking clear-cut experimental evidence, numerical codes are used to embed theoretical expressions for confinement properties in models which account for non-ideal processes and which thus mimic experimental behavior (*Ref. 2-141*). Such codes vary greatly in complexity, some treating only spatial averaged quantities (zero-dimensional), others radial transport (one-dimensional) with and without time dependence, and a few deal with non-circular geometry (*Refs. 2-68, 2-142–2-144*).

The remainder of this section deals with comparisons between experimental observations and theoretical expectations for ion and electron energy losses, for particle transport, and for the diffusion of both naturally occurring and deliberately injected impurities. It is important in this regard to be precise about the terminology employed. The energy confinement time is defined by $\tau_E^* = W/P_{\text{loss}}$ where W is the sum of the electron (W_e) and ion (W_i) energy, and P_{loss} is the total power lost from the plasma. Since, in "steady state", P_{loss} is equal to the total power input (P_{in}), τ_E^* is equal to $\tau_E = W/P_{\text{in}}$. The latter quantity is often called the gross energy confinement time (or the energy replacement time), and is commonly used as a measure of tokamak performance. In order to separate electron and ion energy balance, the electron and ion energy replacement times are defined by $\tau_{Ee} = W_e/(P_{\text{OH}} + P_e - P_{ei})$ and $\tau_{Ei} = W_i/(P_i + P_{ei})$, where P_{OH} is the ohmic heating power, P_{ei} the heat transfer from electrons and ions, and $P_e(P_i)$ the power input to electrons (ions) directly from supplementary heating.

2.2.4.1 Ion Confinement — The confinement of fuel ions in a reactor core will be governed by the drift and collision processes observed in the hottest tokamaks now operating. The control of recycling and the attainment of high volume-averaged fusion energy production will, on the other hand, depend to a large extent on the properties of colder, and hence more collisional plasmas. The entire relevant range of ion collisionality has now been explored in tokamaks and the results so obtained are generally consistent with the predictions of neoclassical theory.

2.2.4.1.1 Collisional Regime ($\nu_{*i} \gtrsim 1$) — The Alcator device has produced plasmas with densities in excess of 10^{15} cm^{-3} and has yielded detailed agreement with the neoclassical confinement model (*Ref. 2-37*). The plasma energy balance is dominated by ion losses at these densities, and the ion losses show both the magnitude and the scaling with density predicted by neoclassical theory.

2.2.4.1.2 Plateau Regime ($\nu_{*i} \sim 1$) — Plasmas with $\nu_{*i} \sim 1$ have been the most extensively studied. The resulting estimates of transport rates all lie within a factor of 3–5 of the neoclassical values with agreement to within a factor of 2 more common. The uncertainty can be

ascribed to imprecise knowledge of the radial distribution of the impurity concentrations (which enters in the expression for the ion thermal conductivity), the electron-ion power transfer rate, and the charge exchange losses.

In the ISX experiment, a density scan at fixed field and current showed a clear saturation of the total confinement time with density (Ref. 2-35). Since anomalous electron confinement, $\tau_{Ee}^{(A)}$, is observed to increase with density (see Section 2.2.4.3) and neoclassical ion confinement, $\tau_{Ei}^{(NEO)}$, decreases with density, the maximum confinement time $\tau_E^{\max}$ should occur near $\tau_{Ee}^{(A)} \approx \tau_{Ei}^{(NEO)}$. This is illustrated in Fig. 2-17 (Ref. 2-145). Here, $\tau_E^{\max}$ has been observed to scale linearly with current and independent of magnetic field (Ref. 2-35). This is a straightforward consequence of neoclassical theory provided $\tau_{Ee}^{(A)}$ is not strongly dependent on I and B .

Fig. 2-17 Density dependence of total energy confinement time for deuterium discharges in ISX-A with $I = 120$ kA and $B_T = 12.8$ kG. The solid line shows full transport model predictions, fit to the experimental point at $\bar{n} = 1.5 \times 10^{13} \text{ cm}^{-3}$. The dashed line shows the corresponding results with ion neoclassical thermal conductivity neglected in the transport model.

2.2.4.1.3 Collisionless Regime ($\nu_{*i} \lesssim 1$) — Neutral beam heating in the PLT device has explored the collisionless regime, with $\nu_{*i} < 0.1$ (Ref. 2-12). Comparison with neoclassical ion transport shows consistency, again to within a factor of 3–5 (Ref. 2-144). In these experiments, much of the power was delivered directly to plasma ions. Because of the relatively low plasma target densities ($\bar{n}_e \lesssim 5 \cdot 10^{13} \text{ cm}^{-3}$), the ion energy balance was effectively decoupled from electron processes, and $\tau_{Ei} > \tau_{Ee}$.

2.2.4.1.4 Summary — In the extreme ranges of collisionality, the ion confinement properties are in reasonable agreement with neoclassical estimates. For example, in ultra-high density Alcator experiments with $\nu_{*i} \sim 8$ the gross energy confinement is dominated by the ion channel and the ion conductivity is close to neoclassical. Experiments with 2 MW of neutral injection on PLT have produced highly collisionless plasmas in which the ion channel can be isolated from the electron processes. Here again, ion confinement differs from neoclassical by at most a factor of 3–5. Similar error bars prevail at intermediate collisionality, where the uncertainties may be an artifact of the uncertainties in the measurement of power input and impurity level.

2.2.4.2 Electron Confinement — Neoclassical theory predicts that the loss through the electron channel is much less significant than ion losses, a prediction which is in disagreement with every tokamak experiment to date. Several mechanisms are thought to account for the electron losses: substantial line radiation losses from both heavy (F_e , W , M_o) and light (C , O) impurities, MHD effects such as sawtooth ($m = 1$) and saturated resistive tearing ($m \geq 2$, $n = 1, 2$) instabilities (or Mirnov oscillations), and high m microinstabilities.

Under some circumstances, the impurity level can be minimized. For example, the level of impurity radiation from the bulk or central plasma is thought to play a negligible role in the high density Alcator discharges (*Ref. 2-146*) as well as in some PLT (*Ref. 2-12*) and ISX (*Ref. 2-35*) discharges.

The sawtooth oscillations are almost universal but they should generally have only a small effect on global confinement for moderate $q(r)$ discharges since only a small central region where $q(r) \leq 1$ is short circuited. However, they may be significant in reducing the central temperature. The Mirnov oscillations, on the other hand, can reduce the confinement by a factor of two or more in that they provide an effective short circuit to confinement over a substantial portion of the plasma volume (*Ref. 2-34*). Experiments in TO-1 (*Ref. 2-56*) have shown a rise in temperature and confinement time when the $m = 2$ amplitude is decreased by feedback control. The ISX confinement experiments were particularly successful in suppressing the $m = 2$ modes by careful current control (*Ref. 2-35*).

In contrast to the radiative and MHD effects, transport losses due to the microinstabilities (or collective plasma oscillations) are believed to dominate transport (see Section 2.2.2) except at high density where, as discussed above, they may be superceded by the ion neoclassical losses. At present there is no generally accepted theory as to which instability is responsible for the fluctuating electric (δE) and magnetic (δB) fields known to exist in tokamaks. It is not even clear whether the transport results from the fluctuating $\delta E \times \hat{B}$ (*Ref. 2-10*) motion across the magnetic fields $\hat{B}$, or via classical motion of the electrons along the fluctuating magnetic field $\hat{B} + \delta \hat{B}$ (*Refs. 2-113, 2-147, 2-149*). In this context many workers have sought to describe transport with simple empirical or semi-empirical scaling models. It is hoped that these will provide some guide to the theory as well as to expectations for future experiments.

Based on ohmically heated discharges ($P_e = P_i = 0$), various "empirical" scaling laws for τ_E , reflecting primarily anomalous electron transport, have been proposed (*Refs. 2-34, 2-150-2-158*). These scalings reflect the response of τ_E as a particular scaling variable is changed. Since the ranges of parameter variation available on individual tokamaks are limited, scaling laws have also been deduced by means of statistical comparisons of the total data base. The most commonly adopted scalings are listed in Table 2-1. General features of the data are the improvement noted in τ_E with increases in plasma density ($\tau_E \sim \bar{n}_e$) (*Refs. 2-15, 2-55, 2-153, 2-158*) and machine size ($\tau_E \propto a^2$ being characteristic of a random walk process). Further, τ_E tends to increase initially with lowering $q(a)$ (or with rising current at a fixed toroidal field), but to fall off at low q (or at large current) (*Ref. 2-34*). These two trends are reflected in scaling laws of the form $\tau_E \sim a$ (*Refs. 2-34, 2-150*), and $\tau_E \sim \bar{n} a^2 \sqrt{q}$ (*Refs. 2-154, 2-15*) respectively. Deterioration of τ_E at low q is believed to be primarily due to Mirnov oscillations (*Refs. 2-34, 2-15*), but empirical relationships involving the amplitude and period of internal disruptions also lead to a similar scaling law ($\tau_E \sim R a^2 n \sqrt{q}$) (*Ref. 2-15*). Note that while some of the entries in Table 2-1 employ the line density $\bar{n}$ and some employ the volume-averaged density $\langle n \rangle$, it is probably safe to say that the accuracy of the identification of τ is insufficient to discriminate between these alternative forms.

The scaling laws represented in Table 2-1 have been deduced by a variety of techniques. Some parametric dependences have emerged from the study of sequences of discharges in a single tokamak with systematically varied parameters and others have their origin in comparison of shots under similar conditions in devices of different size and Z_{eff} . Still another approach, one that has been adopted by several authors (*Refs. 2-155, 2-156*) is statistical in character, the scaling laws originating from a best fit to a large body of documented discharges from many different tokamaks.

Under the conditions that the scalings of Table 2-1 were obtained, the electrons were the major loss channel ($P_{ei} \ll P_{OH}$), so $\tau_E^* = \tau_E \sim \tau_{Ee}$ and the scalings are dominated by electron

TABLE 2-1
SCALING LAWS FOR ENERGY CONFINEMENT
(MKS UNITS)

Gorbunov, Mirnov and Strelkov (Ref. 2-150)	$\tau_E = 4 B_p a^2$
"Alcator" (Ref. 2-158)	$\tau_E = 6 \times 10^{-21} \bar{n}_e a^2$
Daughney (Ref. 2-153)	$\tau_{Ee} = 3 \times 10^{-17} \langle n_e \rangle \langle T_e \rangle^{1/2} a^2 / I$
Jassby, Cohn, and Parker (Ref. 2-154)	$\tau_E = 3 \times 10^{-21} \bar{n}_e a^2 \sqrt{q(a)}$
Hugill and Shettfield (Ref. 2-155)	$\tau_E = 9.1 \times 10^{-16} \langle n_e \rangle^{0.61} a^{1.57} B_T^{0.88}$
Pfeiffer and Waltz (Ref. 2-156)	$\tau_{Ee} = 9.5 \times 10^{-20} \langle n_e \rangle^{0.90} a^{0.98} R^{1.63} Z_{eff}^{0.23}$

confinement. Further, with ohmic heating, confinement and heating are coupled, since the input power depends on confinement through T_e . As a result, the scalings derived, particularly the dependence on I and T_e , may not solely represent *confinement*. It is important to determine which functional dependences survive when scaling laws are derived for different situations.

One such situation is that with P_{ei}/P_{OH} being a significant fraction. Since

$$\tau_E = \tau_{Ee} (1 - P_{ei}/P_{OH}) + \tau_{Ei} P_{ei}/P_{OH}$$

this could result in ion- rather than electron-dominated confinement, i.e., $\tau_E^* = \tau_E \sim \tau_{Ei}$ for extreme conditions. Because of the emphasis on the high density operation to maximize $n\tau$, some recent experiments have reached the ion-dominated regime. This is revealed both in the ISX-A data shown in Fig. 2-17 and in the Alcator data in Fig. 2-18 (Ref. 2-37). In the latter case, the experimental values of τ_E at high densities approach within a factor of 1.5 the calculated values of τ_E which are deduced from the assumption of neoclassical ion thermal conduction. Provided that ion confinement is approximately neoclassical, an assumption which leads to an experimentally verified prediction for $T_i(0)$, τ_E reaches a maximum at $\bar{n}_e = (\bar{n}_e)_{crit}$ corresponding to $\tau_{Ee} = \tau_{Ei}$, even though τ_{Ee} continues to increase for $\bar{n}_e > (\bar{n}_e)_{crit}$. In these experiments, effects of Mirnov oscillations and impurities were not believed to play an important role in the energy balance, so the observed $\tau_{Ee} \sim \bar{n}_e$ dependence, termed Alcator scaling, is considered to be the scaling of the energy confinement time related to transport losses (i.e., thermal conduction and convection).

Auxiliary heating dominated discharges provide the major source of data for extracting the temperature dependence of the scaling laws. Neutral beam injection is a better prospect than rf heating to study scaling laws in detail since it currently provides higher heating power and the radial profile of input power can be calculated from classical processes. The ion energy confinement times observed are approximately consistent with neoclassical predictions. In general, electron energy loss mechanisms are similar with and without injection, so injection power seems equivalent to ohmic heating power (Ref. 2-159). Of particular importance here is that the scaling $\tau_{Ee} \sim \bar{n}_e$ is also observed in injection-heated plasmas (Refs. 2-25, 2-159).

Fig. 2-18 The variation of τ_E with average density at constant B_T and I_p (solid circles) is compared with the neoclassical prediction. Also shown are values of τ_E for high current, low q discharges (open circles); (Alcator data, Ref. 2-37).

The extraction of the temperature dependence of the energy confinement time is of prime interest in extrapolating present experimental results to reactor parameters. Comparisons between high-Z impurity contaminated discharges and clean ones provide some information on the scalings related to transport losses. The shorter electron energy confinement times observed in the former discharges seem to indicate that τ_{Ee} does not deteriorate with increasing temperature (Ref. 2-160), although the extent of effects of profile changes is uncertain. Again the most definitive test is high power neutral beam injection in which significant (up to 50%) increases in T_e have been observed in several tokamaks (Refs. 2-25, 2-159, 2-161, 2-162). This data is discussed in some detail in Section 3.2.4.

The evidence remains inconclusive, but there are some preliminary results from the PLT injection experiments (Ref. 2-25) in which τ_{Ee} in the core of the plasma was observed to improve with injection. Calculations of the deduced electron thermal conductivity ($\equiv n_e \chi_e$ = electron density x electron thermal diffusivity) as a function of radius show a strong increase from the center of the plasma to the outside, scaling approximately proportional to $1/T_e$. With beam heating, the electron thermal conduction in the inner region of the plasma ($r/a < 0.5$) fell further, again approximately proportional to $1/T_e$. In the outer regions where T_e also increased, however, there was no decrease in the thermal conduction coefficient. This improvement in central confinement, while it appears to suggest a favorable scaling of χ_e with electron temperature, may derive from other sources, such

as direct effects of the beam ions on microturbulence, or indirect effects, such as small profile changes. One result which is very clear from these measurements, however, is that the electron thermal conductivity is not flat across the plasma profile in PLT, as would be the case for the commonly assumed $\chi_e \propto 1/n_e$, but rather $n_e\chi_e$ increases strongly from the plasma core to the outer region. This re-emphasizes the need to obtain good beam penetration for adequate electron heating. More importantly, these results indicate that thermal transport does not possess the strong unfavorable temperature dependence predicted by trapped particle modes, a conclusion that takes on considerable importance when coupled with the fact that PLT has operated at collisionalities characteristic of ignited plasmas.

The models discussed above for electron energy confinement or, equivalently, the heat diffusivity χ_e , are generally accompanied by similar models for the plasma diffusivity D . Although the ratio between the two is poorly established, $D \simeq (1/5 - 2/5) \chi_e$ is generally found to be consistent with the measured density of recycled neutrals for small ion density machines. However for large and/or dense machines (and particularly with strong gas puffing), cases in which neutrals cannot penetrate to the plasma center, the simple Fick's Law model for plasma flow ($\Gamma = D\partial n/\partial r$) tends to produce either flat (or in the case of gas puffing, inverted) density profiles. An additional inward plasma flow appears to be required. The Ware pinch (*Ref. 2-163*) and an ion-mixing mode (*Ref. 2-164*) have been suggested as possible mechanisms.

2.2.4.3 Noncircular Cross Section Effects — In recent years, considerable theoretical evidence has accumulated that vertically-elongated cross section plasmas offer more promise than circular cross section plasmas for the achievement of MHD stable equilibria with β -values that are of reactor interest (see Section 2.1.4).

Vertical elongation of the plasma cross section can also have a favorable effect on neoclassical confinement. Consider an elliptical cross section with minor semi-axis a and major semi-axis b . For the case of uniform toroidal current density, the magnetic surfaces are similar ellipses, and it can be shown (*Refs. 2-165, 2-166*) that the neoclassical confinement time of the ellipse is enhanced by a factor $(b/a)^2$ compared to a circle having the same values of a and q .

Although there is some evidence from Doublet IIA experiments that confinement improves with increasing vertical elongation (*Ref. 2-167*), the elongation scaling of the anomalous heat diffusivity has not been empirically established. However, if the diffusivity is a constant, a modest improvement in confinement time follows simply from the increase in the volume-to-surface ratio as b/a increases. The resulting enhancement, compared to a circle of the same a , amounts to $2b^2/(a^2 + b^2)$, and is just the dependence on b of the time required to random walk with fixed step size from the center to the edge of an ellipse.

2.2.4.4 Impurity Transport — The impurity transport rates seem to be consistent with the neoclassical picture in experiments in which a trace amount of an impurity is deliberately injected to study this process. In an ATC experiment (*Ref. 2-168*) Al was injected and diffused inward at the neoclassical rate until it became fully stripped and thus could not be monitored spectroscopically. A similar experiment with Ar injection on T-4 (*Ref. 2-169*) had somewhat clearer results because Ar quickly ionizes to Ar^{16+} at a central electron temperature of about 1 keV but requires considerably higher temperatures for further ionization. The influx rate on T-4 agrees with neoclassical rates, but the continual $m = 1$ sawtooth oscillations occurring in the core provide an expulsion mechanism and thus results in a finite lifetime for the impurities.

For impurity-rich discharges [$Z_{\text{eff}}(r) \gtrsim 3$], the impurities should be considered the dominant species and should be subject to the usual outward diffusion.

A critical question for reactor grade plasmas is whether heavy metal impurities created at the wall can enter the hot plasma core and thereby provide an unacceptably large energy sink. A direct simulation of this effect was carried out on ISX-A, and the tungsten atoms injected by laser blow-off did in fact travel to the core at about the neoclassical rate (*Ref. 2-170*) and produced hollow T_e profiles typical of radiation dominated discharges. A more hopeful note was provided by the same experimental device with preliminary indications of success of a scheme (based on neoclassical theory) for reversing the flow of impurities by poloidally asymmetric gas injection and gettering (*Ref. 2-170*). The apparent poloidal asymmetry of impurity radiation observed in Alcator and PLT is also consistent with neoclassical impurity transport (*Ref. 2-171*).

2.2.4.5 Reactor Sizing — If the condition is imposed that reactor tokamaks will be self-sustaining (ignited), the applicable transport scaling laws provide information about the required size and toroidal magnetic field of the device. The global ignition condition states that the fusion energy deposited in the plasma balances the sum of all plasma losses. For sufficiently large plasma current, the charged fusion reaction products will deposit the bulk of their energy in the plasma, at least according to classical collision theory. In the case of a D-T reaction, i.e.,

each fusion event will give rise to 3.5 MeV of plasma heating. Thus the ignition condition reads (in a volume-averaged sense)

$$\frac{1}{4} n^2 \langle \sigma v \rangle_{D-T} \cdot 3.5 \text{ MeV} = \frac{3 n T}{\tau_E} + n^2 Z_{\text{eff}} R(T) \quad (2-29)$$

where $R(T)$ is the radiative power loss per electron per impurity atom and it has been assumed that $n_D = n_T \approx n/2$, $\tau_{Ee} = \tau_{Ei}$, and $T_e \approx T_i$. The family of solutions to Eq. (2-29) is (*Ref. 2-172*)

$$n\tau = \frac{3 T}{\frac{1}{4} \langle \sigma v \rangle_{D-T} - Z_{\text{eff}} R(T)} \quad (2-30)$$

provided, of course, that the denominator is positive, i.e., $T > 4.4 \text{ keV}$ for $Z_{\text{eff}} = 1$. In the following, the impurities are taken to be low enough Z so that they are fully stripped at the (high) temperatures required for ignition; hence only bremsstrahlung radiation need be retained. In this case, the minimum value of $n\tau$ for ignition is about $6 \times 10^{13} \text{ cm}^{-3} \text{ sec}$, and occurs at a temperature of about 40 keV. Actually, ignition can occur at $n\tau < 3 \times 10^{13} \text{ cm}^{-3} \text{ sec}$ provided $\tau_{Ei} \gg \tau_{Ee}$ and T_i is significantly greater than T_e , conditions that may prevail in low density, intense auxiliary heating scenarios (*Ref. 2-173*). Most reactor designs, however, have average temperatures of order 10 keV because this leads to maximum fusion power for a given plasma beta. In such cases, the ignition $n\tau$ is in excess of $10^{14} \text{ cm}^{-3} \text{ sec}$.

To obtain a quantitative estimate of the device parameters needed for ignition, a specific scaling law must be adopted. The choice adopted here is Alcator scaling; in MKS units,

$$\bar{n}\tau_E = 6 \times 10^{-21} (\bar{n} a_{\text{eff}})^2 \quad (2-31)$$

where a_{eff} is a measure of the effective plasma size, taking into account whatever modification is appropriate for deviation from a circular cross section. Clearly the minimum reactor size depends upon the achievable density.

One key to reactor sizing is the constraint that β cannot exceed a given value determined by MHD stability considerations. The minimum β ignition point corresponding to the scaling law in Eq. (2-31) can be determined analytically provided the approximation $\langle \sigma v \rangle_{D-T} \propto T^2$ (valid to 20% accuracy for $8 < T < 25$ keV) is made and that assumptions are made about the plasma profiles (*Ref. 2-174*). The result for parabolic $n(r)$ and $T(r)$ is

$$\beta_{\min} = \frac{1.34 Z_{\text{eff}}^{1/3}}{a_{\text{eff}} B^2 f_p^{5/3}} \quad (2-32)$$

where f_p is the ratio of primary species ions to electrons. For $\beta = 6\%$ and $B = 5$ T, numbers typical of recent reactor designs, and in addition if $Z_{\text{eff}} = 1$, $a_{\text{eff}} = 0.9$ m, roughly the size of TFTR. Larger devices are required, however, to contend with the presence of a modest concentration of low- Z impurities, to give some margin for ignition, and to provide a design margin in β to allow for the thermal excursion following ignition.

An energetically stable solution to Eq. (2-29) requires an energy loss term which increases with T more rapidly than does $\langle \sigma v \rangle$. This is not the case for Alcator scaling so it cannot be the scaling law that governs the burn point of a reactor. It is possible that the enhanced transport that will accompany the violation of the β "limits" will provide a natural burn control mechanism. Another possibility is the contribution to ion energy transport arising from the magnetic ripple, a contribution that increases dramatically with increasing temperature. Fortunately for the reactor designer, density and temperature profile information are not essential for component design. The key item of information for such considerations is the wall loading (or equivalently, the fusion power) and this depends only on β , not separately on n and T , at least to the extent that the approximation $\langle \sigma v \rangle_{D-T} \propto T^2$ is valid.

3. PLASMA HEATING

Because the plasma resistivity is proportional to $T_e^{-3/2}$ and the toroidal current is limited by MHD instability, it becomes increasingly difficult to reach higher temperatures by means of ohmic heating alone. In the face of the anticipated loss mechanisms, some additional heating is required to produce reactor-like plasmas. This chapter describes both the limits of ohmic heating and the various types of auxiliary heating which may be employed to push beyond the ohmic heating regime.

Conceptually, the simplest heating technique is to inject an energetic beam of particles into the plasma. As the particles thermalize, their directed energy is converted into heat, thus raising the plasma temperature. A beam composed of neutral atoms is able to cross the magnetic fields surrounding and permeating the plasma; ultimately, the beam particles become trapped on field lines within the discharge by means of ionizing collisions with plasma ions and electrons. In the case of the injection of a reacting beam (e.g., deuterium injection into a tritium plasma) the energetic ions also undergo fusion reactions as they slow down, a process that may generate substantial amounts of fusion energy. Such beam fusion events will provide the bulk of the fusion power in the TFTR experiment.

To date, neutral beam heating has been the primary means of delivering large amounts of auxiliary power to tokamaks. An alternative approach that has particular technological appeal for reactor application is heating via electromagnetic waves. Because a number of resonances exist for a plasma in a magnetic field, several options exist for rf heating. Although the data base in this area has greatly expanded in recent years, no option has yet been explored sufficiently thoroughly to permit a confident extrapolation of the underlying physics to the reactor regime.

The chapter has been divided into sections dealing with ohmic heating, neutral beam injection, and rf heating, and a final section concerned with a number of the alternatives to these main-line approaches, specifically, plasma compression, turbulent heating, relativistic electron beams, and cluster injection. Compression is of particular interest since it is complementary to other approaches and may be used to give the extra boost needed to reach ignition.

Auxiliary heating has served the role of producing the plasma conditions that most closely resemble reactor operation. The most notable success to date has been the use of neutral beam heating on PLT, where peak ion temperatures of 6.5 keV have been obtained. The minimum ion temperature needed for breakeven in the presence of the unavoidable bremsstrahlung losses is about 4.5 keV for a D-T system, while the volume-averaged operating temperature for a D-T fusion reactor is expected to be in the range of 10–15 keV. Thus the achievement in PLT puts tokamaks within a factor of two of being able to study transport at true reactor temperatures.

Of equal importance to the achievement of high T_i is the ability to reproduce the conditions that permit a confrontation with theoretical models of transport under reactor conditions. For many models, the critical parameter characterizing confinement is ν_{*i} , the ratio of the ion banana detrapping frequency due to Coulomb collisions to the ion bounce frequency. In PLT a minimum ν_{*i} of 2×10^{-2} was obtained, and ν_{*i} was less than 0.1 over most of the radial profile, values that are quite comparable to those anticipated for an ignited plasma. Trapped ion modes and ion temperature gradient modes are predicted to be unstable under these conditions of low ν_{*i} . Since such

modes lead to highly unfavorable confinement scaling with increasing temperature the observation that the energy confinement time did not deteriorate offers strong encouragement that such modes will not be an important factor in the energy balance of reactor tokamaks.

Caution must always be exercised, however, in extrapolating based on results obtained in experiments that have advanced only a subset of the relevant parameters into the reactor regime. But, barring the evolution of an all-encompassing model of tokamak physics, there are no alternatives to such extrapolations, since the production of a genuine reactor-like plasma will not precede the construction of the reactor itself. The thrust of the program is to utilize increasingly ambitious devices to probe further into the reactor regime and, to a great extent, the key to this effort over the next few years lies in the multi-megawatt injection experiments planned for ISX-B, PDX, Doublet III, and TFTR, as well as high power rf experiments planned on PLT and Alcator C.

3.1 OHMIC HEATING

The simplest, most efficient and most economical method of heating a tokamak is joule heating by means of the confining current. This method is indigenous to the tokamak concept, and it can be argued that much of the early success of the tokamak was due to the elevated temperature regimes which were made easily accessible solely through use of the intrinsic ohmic heating capability. Regrettably, the parameters attainable by ohmic heating are limited by basic considerations. Foremost among these is the fact that at fixed current density, the ohmic power density decreases as $T^{-3/2}$. Of course, the power density can be increased further by raising the current density, but in so doing, the confinement properties tend to degrade, for example, by the spreading of the $q = 1$ surface and the associated increase in MHD activity, or by increase of the drift parameter, with the result that no further increase in temperature can be achieved.

This section deals with the capability of ohmic heating with emphasis on the limits of the regimes attainable by this technique. A combination of theoretical and empirical considerations will be employed in this assessment, and care must be used in forming conclusions where the results are based on an incomplete understanding of the underlying physics. In what follows MKS units are implied, except in the case of temperatures, which are in electron-volts.

3.1.1 Ohmic Heating Power Density

The power density associated with ohmic heating is

$$P_{OH} = \eta j^2 \quad (3-1)$$

where, to a good approximation, j is the parallel current density and η is the parallel resistivity. For a hydrogenic plasma, $\eta = \eta_s$ where

$$\eta_s = 1.1 \times 10^{-3} \tilde{\Lambda} T_n^{-3/2} \quad (3-2)$$

and $\tilde{\Lambda} = \ln \Lambda / 20$. For $Z_i \neq 1$, the expression for η_s should be multiplied by $G(Z_i)$, defined in Eq. (2-12), but only hydrogenic plasmas are considered in the remainder of this section.

In actual fact, the resistivity may differ from the expression given in Eq. (3-2) in at least two respects: first, the neoclassical treatment reveals an enhancement due to trapped electrons which carry no current but nevertheless provide friction to the current-carrying particles (see Section 2.2.1.3). The primary effect is of order $\epsilon^{1/2}$ and is therefore small near the magnetic axis where the

current density is large; consequently, the net increase in power density due to neoclassical effects is small and will henceforth be neglected. Equation (3-2) may undergo significant modification in the case of low particle density, high current density operation, specifically, when the parameter $\xi \equiv v_D/v_{th,i}$ becomes non-negligible. In hydrogen plasmas, $\xi > 0.05$ results in strong nonthermal distribution functions with a concomitant reduction in resistivity (*Ref. 3-1*). In fact, values of $\eta < 0.5 \eta_s$ have been observed (*Ref. 3-2*). For reactor-grade plasmas characterized by relatively high density and temperature and low current density, $\xi \ll 0.05$ so departures from Maxwellian distribution functions due to significant population of the tail should not be encountered.

The current density is limited by considerations of stability. When the safety factor q approaches unity near the magnetic axis, an $m = 1$ mode develops which degrades confinement within the $q = 1$ region. This results in a loss of energy in the central region, a rapid cooling and a redistribution of current with q slightly larger than 1. The region reheats until $q \sim 1$ and the process repeats. The phenomenon is known as the thermal or "sawtooth" instability and is discussed in greater detail in Section 2.1.

Based on these considerations, the maximum useful central current density for a circular cross section tokamak is

$$j_{\max} = \frac{2}{\mu_0} \frac{B_T}{R} \quad (3-3)$$

For example, in a reactor with $B_T = 5$ T and $R = 5$ m, $j_{\max} = 160$ A/cm². Among existing tokamaks, the Alcator devices have the highest current densities. In Alcator C, $j_{\max} = 3.5$ kA/cm² is permitted, assuming the peak field of 14 T is achieved. Under these conditions, a central electron temperature of about 2 keV is expected, corresponding to an ohmic heating power density of 14 W/cm³. Power densities of 20–30 W/cm³ have already been achieved in Alcator A.

3.1.2 Regimes Accessible by Ohmic Heating

Owing to the intrinsic limits on power density associated with ohmic heating, the regimes accessible for parameters such as density, temperature, and β depend in addition to machine parameters only on the appropriate confinement law. This is to be contrasted with auxiliary heated plasmas for which higher temperatures can be achieved by increasing the applied power. In the following, the regimes permitted by ohmic heating are discussed for three important cases:

1. Empirical energy confinement ($\tau_E \sim n a^2$).
2. Radiation dominated energy loss.
3. Neoclassical energy confinement.

The results will be used not only to estimate the performance of contemporary tokamaks, but also to assess the prospects for reaching ignition utilizing ohmic heating alone.

3.1.2.1 Empirical Energy Confinement — Electron energy losses in tokamaks have been found in a variety of devices to obey so-called "Alcator" scaling (*Ref. 3-3*) in which

$$\tau_E = 6 \times 10^{-21} \bar{n} a^2 \quad (3-4)$$

While individual tokamaks differ from this law by as much as a factor of 2 and there is some evidence of a weak dependence on $q(a)$ and on the plasma elongation (see Section 2.2.4), these second-order effects will be ignored in estimating the ohmic heating performance.

While it is possible to examine the global power balance based on Eq. (3-4), the results depend on unknown temperature profiles, and a simpler calculation using a central power balance can be

more easily carried out. To this end, it is assumed that $\tau_{E0} = \tau_E$, i.e., that the central confinement time is equal to the global time given in Eq. (3-4).

It then follows from $\eta_0 j_0^2 = 3 n_0 T_0 / \tau_E$ that

$$T_0 = 705 \left(\frac{a}{R} \frac{B_T}{q_0} \right)^{4/5} \quad (3-5)$$

where the assumption has been made that $n_0 = 2\bar{n}$. Clearly, the highest temperatures for ohmically heated tokamaks are to be obtained in high-field, low aspect ratio devices. Notice that although the temperature does not depend on the size of device, large tokamaks have lower ohmic heating power density and are more susceptible to being limited by radiative losses, which have been ignored in the above.

Quite extreme parameter regimes must be achieved in order to push the central temperature even to the minimum level required for ignition ($T > 4.5$ keV). For example, the choice $B_T = 15$ T, $R/a = 3$ and $q_0 = 0.8$ yields $T_0 \sim 3$ keV. Thus the prospect for achieving ignition with ohmic heating as the only input power is unlikely as long as the "Alcator" scaling is obeyed. Modest improvement in confinement (under ohmically-heated conditions) with increasing temperature may permit temperatures in the 4 keV range, and in this event the conclusion on the question of ohmic-heating to ignition would need further evaluation because alpha heating begins to be a significant factor in the energy balance. Incorporation of the neoclassical enhancement of resistivity (see Section 2.2.1.3) also results in somewhat higher predicted temperatures.

Additional information on confinement under ohmic heating conditions should become available from Alcator C. It should be noted that even if favorable transport scaling does permit ohmic heating to ignition, the outlook for reactor application remains doubtful because the required fields and aspect ratios appear unachievable when superconducting toroidal field coils are employed and the energetics is unacceptable in the case of normal toroidal field coils.

3.1.2.2 Radiation Dominated Regime — Since confinement time generally improves with density (see Section 2.2.3), the Lawson number $n\tau$ improves with increasing n and the optimum value of this parameter in ohmically-heated discharges therefore depends on the maximum achievable density. As illustrated in Fig. 3-1, taken from *Ref. 3-4*, the observed maximum density depends linearly on the current density and hence scales as B_T/R .

The density limit can be qualitatively understood by postulating a density-dependent loss mechanism, for example bremsstrahlung or impurity radiation. In the former case,

$$P_B = 1.7 \times 10^{-38} n^2 T^{1/2} \quad (3-6)$$

whereas, in the case of impurities, the power loss depends on $nn_I f(Z_I, T_I)$ where f is a function of the charge Z_I of the impurity, and n_I is the impurity density. Equating P_B to the ohmic power and using Eq. (3-3) yields an absolute upper limit on energy density (corresponding to infinite τ_E and $Z_{eff} = 1$),

$$(nT)_{\max} = 3.9 \times 10^{23} \frac{B_T}{R} \quad (3-7)$$

For most present-day tokamaks, which operate with ohmic heating in the vicinity of 1 keV, this yields a density limit of order 10^{15} cm^{-3} . The addition of impurities can cause the density limit to be substantially lower and this is probably a factor in presently observed density limits.

3.1.2.3 Neoclassical Limits — Since the empirically determined thermal diffusivity decreases with density, whereas the neoclassical value increases with density, the possibility of a crossover occurs, in which case the discharge energy balance changes in character from one dominated by "anomalous" electron loss to one in which neoclassical ion thermal transport loss is dominant. Such a transition has been observed in both Alcator A and ISX-A (see Section 2.2.4, Fig. 2-17). As shown below, new limits on ohmic heating capability are applicable in the neoclassical dominated case.

At sufficiently high density in an ohmically-heated tokamak the ions and electrons are close to thermal equilibrium since the energy-exchange time is much shorter than either the electron or ion energy confinement time. An appropriate model is then one in which the ohmic heating power deposited inside a given plasma surface is balanced by the neoclassical ion thermal transport across the boundary of that surface. The result is a limit on the plasma pressure which takes the form $\beta_p \lesssim 1$ for typical plasma profiles (Ref. 3-5). No separate constraints on density or temperature emerge.

Assuming that the anomalous transport is governed by Eq. (3-4), the confinement time is maximized by choosing the density such that the empirical confinement time equals the neoclassical value. The density for which this occurs, $\bar{n}_c$, is given approximately by

$$\bar{n}_c = 2 \times 10^{20} \frac{B_p}{q(a)} \quad (3-8)$$

for typical plasma profiles. This estimate is about 50% higher than the transition density observed in ISX-A (Ref. 3-6) but is in good agreement with the corresponding density found in Alcator A (Ref. 3-7). It can be concluded, therefore, that operation of ohmically-heated tokamaks at sufficiently high density can result in confinement properties approaching neoclassical estimates.

Fig. 3-1 Maximum plasma density (black dots) reported for various tokamak experiments as a function of $B_T/R_0 \propto j(0)q(0)$, which is a measure of the current density at the plasma center, assuming $q(0) = 1$. The black triangles indicate that the densities were obtained by pulsing in cold neutral gas after a discharge equilibrium was reached. The open circles are the densities obtained in the ORMAK experiment in which initial filling pressure and $q(a)$ were kept constant.

3.2 NEUTRAL BEAM INJECTION

Neutral beam injection is now a tried and proven method of heating tokamak plasmas and is the mechanism by which the most reactor-like plasmas have been produced to date. Much of the theory of neutral injection has been tested experimentally and has been found, within reasonable error limits, to be quite satisfactory. The development of the technology required for the next generation of large tokamak devices, involving multi-megawatt individual beamline systems at energies of 60–80 keV/amu, is proceeding well. However the question of how much further neutral beam technology will need to be extrapolated for application to reactor-level plasmas has not been fully answered, and if many tens of megawatts of multi-second beam power in the range of ≥ 100 keV/nucleon are required, then a very substantial research and development effort in high voltage neutral beam technology will be necessary. Fortunately, recent simulations of heating to ignition require less dramatic technological advances.

The present generation of beam-injected experiments have produced very encouraging results for neutral beam heating specifically and for tokamak physics generally. The scaling laws deduced from beam heated plasmas form the basis of current reactor designs and the behavior of the hot ions so produced gives credence to classical models of alpha heating. The mandate of the auxiliary heating program has been to investigate ion and electron thermal transport in high temperature plasmas. In this context, it is important to recognize that even a very modest dependence of energy confinement on the relevant variables can have a dramatic impact on the observed confinement in the transition from collisional plasmas with temperatures of order 1 keV on axis, as found in present ohmically-heated tokamaks, to collisionless plasmas with temperatures in excess of 10 keV for reactor-grade machines. The results of injection experiments to date in PLT together with those planned on ISX-B, PDX, D-III and TFTR should provide the data base for the needed extrapolation to high-temperature, finite beta reactor systems.

3.2.1 Technological Base

The basic elements of present neutral beam injection systems are exhibited in Fig. 3-2, which illustrates an ORNL constructed system employed on PLT. The principal ingredients are an arc chamber, an acceleration structure, a neutralizer cell, and finally a beam-line to transport the neutral beam into the tokamak. In the arc, a cold dense hydrogen or deuterium plasma is formed by drawing some hundreds to thousands of amperes of plasma current through a gas-filled chamber. Various cusp (*Ref. 3-9*), multipole (*Ref. 3-10*), and reflex arc (*Ref. 3-11*) magnetic confinement geometries have been employed to provide good efficiency in the generation of this plasma. Alternatively, since the arc power consumption is relatively modest compared to the power consumed in accelerating the beam it produces, it is simplest to use no magnetic confinement at all, as in filament sources such as those produced at LBL (*Ref. 3-12*). There is some indication, however, that improved plasma confinement in the arc chamber results in an improved ratio of H^+ to H_2^+ and H_3^+ in the source plasma (*Ref. 3-13*). This is desirable since molecular ions in the source plasma give rise to fractional energy components in the neutral beam, and these components may not penetrate the plasma satisfactorily, particularly in the case of a large reactor-grade tokamak. To date, the most dramatic example of improved species yield is the combination duoPIGatron/cusp-geometry source developed at ORNL, which gives a beam composed of 85% H^+ ions, to be compared with an H^+ species yield of $\sim 50\%$ found in the earlier generation of smaller, less well-confined sources (*Ref. 3-11*).

Fig. 3-2 A schematic of the ORNL/PLT beamline (Ref. 3-8).

In order to extract ions from the source plasma, and to accelerate them to energies many times above the plasma thermal energy, a three or four grid accelerating structure is employed. In a three grid system, such as the ORNL/PLT beams, the first gap, nearest the arc chamber, extracts ions from the arc, and accelerates them to slightly above the injection energy. The second gap then decelerates the ions by typically 1–2 keV, and provides an electrostatic shield against electrons generated downstream, which would otherwise be accelerated in the reverse direction, back into the ion sources. With careful design of the electrode shape and spacing, beam divergences of $\sim 1^\circ$ and current densities of 200 mA/cm² have been obtained in these systems (*Ref. 3-12*). At voltages greater than 40–60 keV, it is advantageous to employ a four-grid system, where the tasks of ion extraction and acceleration are split between the first two of the three gaps. In these systems, due to the increased flexibility of the design of the accelerator structure, and the increased ratio of the beam energy to the temperature of the arc plasma, the beam divergence can be reduced even further, although at high beam voltages the current density must be lowered somewhat.

After ions from the arc plasma are accelerated to form a high-energy, high-current positive ion beam, they must then be converted to atoms in order to penetrate the tokamak magnetic fields. This is accomplished by passing the ion beam through a gas cell, where charge-exchange reactions take place, converting some of the ambient gas to ions, and much of the beam to neutrals. In general, the gas cell is closely coupled to the accelerator structure in order to take advantage of the gas streaming from the arc chamber, and to minimize space charge blowup of the un-neutralized beam. Unfortunately, the efficiency of the neutralization process is intrinsically limited by the fact that impact ionization collisions between the beam atoms and the neutralizer gas molecules compete with the charge-exchange reaction. At low velocities charge-exchange dominates, but at velocities corresponding to > 50 keV/nucleon, impact ionization becomes more probable than charge-exchange, in which case neutrals are stripped as rapidly as they are created. The fall-off in the sum of $\sigma_{cx} + \sigma_{ii}$ with increasing energy also dictates that the neutralizer thickness, nl , must increase as the beam energy increases.

In reactor-grade tokamaks, energies in excess of 100-keV per nucleon may be required for adequate penetration of the neutral beam into the plasma (see Section 3.2.2). Under these circumstances an alternative technique of beam generation employing negative ions, which in principle can be quite efficient in this high energy range, may be necessary. In this approach, H^- ions are produced in a special negative ion source, or are created by passing a low-energy positive ion beam through an alkali metal vapor cell. These ions are afterwards accelerated to very high energies. The negative ions can then be stripped to H^0 atoms, and since this stripping process is efficient at high velocities, one can envision high-voltage neutral beams with high overall efficiencies. The status of negative ion technology, still in its infancy and yet to be employed in a tokamak application, is reviewed in *Ref. 3-14*.

After the neutral beam has been created, a beamline system is required to provide the interface to the tokamak. Important features of such a system (taking the ORNL/PLT beamline shown in Fig. 3-2 as the prototype) are large liquid helium cryopumps, which limit the gas flow into the torus, the bending magnet, which deflects the un-neutralized fraction of the beam onto water-cooled copper beam dumps, and the retractable calorimeters, which provide a beam-stop and power measurement when the ion source is run without the tokamak. Beam dumps and calorimeters present difficult design problems, since the maximum power density reaches a level sufficient to melt any material surface presented directly to the beam. Thus the dumps must accept the beam at grazing incidence, and thick copper plates that store heat inertially, or thin walled, rapidly water-cooled swirl tubes (*Ref. 3-15*) must make up the surface facing the beam. An attractive solution to the problems associated with deflecting the un-neutralized portion of the beam onto ion dumps is to instead electrostatically decelerate the ion beam and recover the initial energy investment. This approach, termed

direct conversion, may extend the energy range for positive ion beams with acceptable efficiency. Initial results in direct energy recovery have been encouraging (*Refs. 3-16, 3-17*), but so far this feature has not been incorporated in a tokamak neutral beam heating experiment.

The drift duct connecting the beam line to the tokamak vacuum vessel represents another difficult problem. No matter how effective the pumping system in the beam line is, some gas streams from the neutralizer directly into the drift duct, and also, in most geometries, some part of the low power halo of the beam impinges on the drift duct walls and liberates gas. Thus a neutral gas density builds up in the drift duct, and then because of the tokamak's magnetic fields, any beam which is ionized by collisions with the gas molecules in this region is swept to the duct walls. In the most pessimistic scenario, the deflected ion beam liberates more gas, and the duct is eventually completely clogged by a nonlinear buildup of neutral gas (*Ref. 3-18*). Measurements on the ATC (*Ref. 3-19*), and PLT (*Ref. 3-8*) experiments indicate, however, that this nonlinear effect has not in fact materialized, but that the simple beam reionization on gas generated in the drift duct does reduce the throughput beam power by $\lesssim 10\%$. Even this problem may be avoidable by providing greater pumping and conductance in the drift duct region.

3.2.2 Beam Penetration

The neutral beam leaving the drift duct easily crosses the magnetic fields and the tenuous plasma surrounding the discharge, and then passes into the plasma column itself. Within the plasma, the neutral atoms become ions through impact ionization collisions with electrons, ions and impurities, and by charge-transfer with plasma ions and impurities. Since the beam ions are typically traveling much faster than the plasma ions, but much slower than the electrons, the total cross section for beam stopping is approximately given by

$$n_e \sigma_{\text{tot}} = n_i \sigma_{ii} + n_i \sigma_{cx} + \frac{n_e \langle \sigma v \rangle_{ei}}{v_b} + \sum_j n_j (\sigma_{ji} + \sigma_{jcx}) \quad (3-9)$$

where $\langle \rangle$ indicates the average over a Maxwellian electron distribution, and the sum over impurity species, j , takes into account impact ionization and charge-exchange with all the impurities within the plasma. *Ref. 3-20* is a useful compendium of the relevant atomic physics data.

A concern for beam penetration has been the possibility of an impurity buildup instability associated with high energy injection into large plasmas. In this instability the impurity level would build up, stopping the beam at the plasma surface. The resulting surface heating would cause increased impurity generation, and higher Z_{eff} , until none of the beam penetrated to the center of the plasma. Recent calculations of the relevant cross sections (*Ref. 3-21*), which give good agreement with the available measurements, are well approximated by the formula (*Ref. 3-22*):

$$\sigma_{jcx} + \sigma_{ji} = 4.6 Z_j \times 10^{-16} \times \left\{ (32 Z_j A/E) \left[1 - \exp(-E/32 Z_j A) \right] \right\} \quad (3-10)$$

where E/A is the neutral beam energy in keV/amu. For fully ionized impurities, Z_j is the nuclear charge, while for partially ionized impurities, an "effective charge" is required. However, for

impurity charge states with ionization potentials greater than a few hundred eV, as found in tokamak plasmas, the effective Z_j is simply the shielded nuclear potential. Since each impurity ion replaces Z_j hydrogen ions, at energies below $E/A = 32 S_j$ these cross sections indicate only a small effect on beam penetration due to an increase in impurity concentration. Thus the concerns about the impurity buildup instability have vanished except for the case of very high energy beams (≥ 200 keV/amu).

Even if trapping by impurities is not a factor, beam penetration into reactor grade tokamak plasmas remains problematic. Since the results of a number of tokamak experiments indicate $\tau_{Ee} \propto n_e a^2$ (see Section 2.2.4) and the quality of beam penetration certainly is characterized by the quantity (hot ion mean free path/plasma size) $\propto (n_e a)^{-1}$, the direction of increasing $n\tau \propto n_e^2 a^2$ is precisely the direction of decreasing beam penetration. Measurements of electron thermal diffusivity versus radius (*Ref. 3-23*) indicate a very strong increase of χ_e with increasing r , which results in a strong dependence of electron heating on beam penetration. In the PLT neutral beam heating experiments, with tangential injection, at densities such that $\bar{n}_e \sigma_{tot} a > 1$, the efficiency of electron heating drops rapidly with increasing plasma density (*Ref. 3-24*). For a $Z_{eff} = 1$ plasma, σ_{tot} is approximately proportional to E_b^{-1} , so beam penetration can be improved by moving to higher energies. In addition, by going to perpendicular injection, penetration can be improved through shortening the path length to the plasma center, but only by a factor of perhaps 1.5. Thus for a reactor-grade plasma, rather high beam voltages (200–500 keV for deuterium beams), and direct recovery and/or negative ion technology will be needed to achieve penetration to the plasma center. Fortunately, scenarios for heating and fueling concurrently can produce an ignited plasma with D^0 beam energies of 150 keV or less when the effects of alpha heating are taken into account (*Ref. 3-25*). With this optimistic viewpoint, higher energy beams may only be needed for driven operation.

Recently an alternate scheme has been proposed to permit penetration of modest energy neutral beams into reactor-grade plasmas. In this approach, called ripple injection (*Ref. 3-26*), neutral beams are fired vertically into the plasma. Special ripple coils adjust the toroidal variation of the main magnetic field so that the injected ions are mirror trapped with the local magnetic wells created by the field "ripple". The deposited beam ions $B \times \nabla B$ drift vertically into the core of the plasma, and then, since the ripple is adjusted to fall off on the opposite side of the torus midplane, the beam ions are ejected from the ripple wells in the center of the discharge. Unfortunately, a number of problems may be associated with this apparently ideal scheme to improve beam penetration. Simple calculations suggest that there may be significantly enhanced ion energy transport (beam and thermal) out of the discharge due to nonaxisymmetry (*Ref. 3-27*). In addition there may be an enhancement of electron energy transport because the magnetic perturbation leads to a breakup of flux surfaces (*Ref. 3-28*). Finally, the microstability of the ion beam drifting into the center of the discharge is questionable due to the strong anisotropy of the energetic ion distribution function (*Ref. 3-29*). On the other hand, the ripple fields, if they are made adjustable during the discharge, could be very useful for burn control in an ignited plasma, since the field perturbation exerts its strongest influence over the most collisionless ions in the plasma; thus it could be used to control the confinement of the 3.5 MeV alpha particles which sustain the thermonuclear burn (*Ref. 3-30*).

In light of the potential problems with ripple-injection, and the known difficulty of negative ion technology, the present designs for large tokamaks beyond TFTR all call for near perpendicular horizontal injection of positive ion based high-voltage neutral beams. Under certain assumptions about anomalous transport and impurity generation, the modest penetration provided by 150–200

keV D^o beams is adequate to bring these large plasmas to ignition, but the margin of safety would be greatly enhanced by success in developing direct recovery systems, in the negative ion R&D program, or in the ripple-injection demonstration experiment planned for ISX-B.

3.2.3 Beam Ion Thermalization

The first problem associated with neutral beam injection in tokamaks that was addressed, both theoretically and experimentally, was the question of how beam ions transfer their energy to the background plasma. The predictions emerging from a straightforward analysis based on Coulomb collisions have, almost without exception, been validated by a wealth of experimental data on a variety of devices. The status of theory and experiment are reviewed in Sections 3.2.3.1 and 3.2.3.2, respectively.

3.2.3.1 Theory — Early calculations based on the assumption that beam ions lose their energy solely through Debye-shielded binary Coulomb collisions showed that at high beam voltages most of the injected energy was transferred to the plasma electrons, while at lower voltages, transfer to the ions was more efficient (*Ref. 3-31*). The drag rates on ions and electrons are given by

$$\begin{aligned}\left. \frac{dE}{dt} \right|_i &= \frac{E_c^{3/2}}{E^{1/2} \tau_s} \\ \left. \frac{dE}{dt} \right|_e &= \frac{E}{\tau_s}\end{aligned}\tag{3-10}$$

where τ_s is the Spitzer slowing-down time on electrons, and E_c is the critical energy at which the energy transfer rates to the ions and electrons are equal:

$$\begin{aligned}\tau_s &= 6.32 \times 10^8 A_b T_e^{3/2} / (n_e Z_b^2 \ln \Lambda_e) \\ E_c &= A_b ([Z]/\bar{A}_i)^{2/3} 14.8 T_e \\ [Z]/\bar{A}_i &= \sum_i n_i A_i^2 \ln \Lambda_i / (n_e A_i \ln \Lambda_e)\end{aligned}\tag{3-11}$$

The units in these expressions are: temperatures and energies, T_e , E , and E_c (eV); atomic masses, A_b and A_i (amu); and densities, n_e and n_i (cm^{-3}). The stopping time, or the mean time required for a beam ion to slow down from the injection energy, E_o , to the average thermal ion energy, $(3/2)T_i$, is given by

$$\tau_{st} = \frac{\tau_s}{3} \ln \left[\frac{E_o^{3/2} + E_c^{3/2}}{(3 T_i/2)^{3/2} + E_c^{3/2}} \right]\tag{3-12}$$

Early on it was recognized that two additional processes could compete with ion and electron drag, and thereby reduce the efficiency of neutral beam heating. In the first case, charge-exchange between energetic beam ions and the thermal neutrals naturally in the plasma, or introduced with the beam itself, can give rise to a "second generation" of beam neutrals. These have a reasonable probability of escaping from the discharge, and therefore bombarding the vacuum vessel walls with energy which they otherwise would have transferred to the plasma. Under low plasma density and high temperature conditions, the neutral atom density is maximized, and the plasma drag is minimized, so charge-exchange can become an important loss process. For typical PLT discharges, charge-exchange losses range between 5% and 30% of the input beam power, depending on plasma density and temperature.

The second process which can decrease the efficiency of energy transfer to the plasma is loss of fast ions onto orbits which escape from the discharge column. For beam ions with largely parallel or anti-parallel velocities, i.e., those satisfying $1 - |v_{||}/v| \ll (2\epsilon)^{1/2}$, the particle trajectories describe nested toroidal orbit surfaces shifted in major radius from the magnetic flux surfaces by a quantity

$$\Delta = \frac{g v_b}{\Omega_{ci}} \left(\frac{1 + \xi^2}{2\xi} \right) \quad (3-13)$$

where $\Omega_{ci} = 9578 B_t(\text{gauss})/A_b$ is the ion cyclotron frequency and $\xi = v_{||}/v$ is the cosine of the pitch angle. Note that for ions traveling parallel to the plasma current (co-injection) the shift is outwards relative to the magnetic surfaces, while for anti-parallel or counter-injection the shift is inwards.

Since the neutral beam must necessarily be injected from the outside (in major radius) of the plasma, the orbit shift results in the loss from the discharge column of all the counter-injected ions which are produced in the outermost region of the plasma. For typical cases of tangential injection, $\Delta/a \approx 0.1$ and the plasma density is relatively low in the outer periphery of the discharge so only a modest fraction (10–20%) of the counter-injected beam is depleted by prompt orbit loss. Even this modest loss fraction can be avoided by using co-injection exclusively; however, there is concern that high-power unidirectional injection could give rise to rapid toroidal rotation, with accompanying instabilities and deleterious effects on confinement. The results from PLT, on the other hand, show no significant deterioration at $v_\phi = 10^7$ cm/sec, or $v_\phi/v_{th} \approx 1/6$, and suggest that toroidal momentum is not so well contained that v_ϕ/v_{th} will increase greatly in future experiments.

A significant advantage can be gained by aiming the neutral beam slightly off perpendicular in the co-direction. Under these circumstances many of the beam ions have $|v_{||}/v| < (2\epsilon)^{1/2}$, so upon ionization they immediately execute banana orbits. Since the outer half of a trapped particle orbit is always in the co-direction, co-injected near-perpendicular ions always deposit their energy inwards of the birth point. The experimental experience to date, however, has been predominantly with more nearly tangential injection.

In addition to the prompt loss process due to initial deposition onto bad orbits, beam ions can also be lost as they slow down by pitch-angle scattering onto unconfined orbits (*Ref. 3-32*). For counter-injection cases where the banana width is comparable to the plasma minor radius, this can be a dominant loss process, since passing counter particles move outwards when they scatter around to being trapped particles. The pitch angle scattering rate is given by the drag rate on the ions, multiplied by Z_{eff} . Thus at high energies ($E \gg E_c$) fast ions dominantly drag on the electrons without pitch angle scattering, while at low energies, and for $Z_{eff} > 1$, the ions scatter around in angle more rapidly than they thermalize, so for counter-injection in a case with large

banana orbits much of the beam energy can be lost from the plasma. By contrast, because co-injected ions travel inwards when they execute banana orbits, they are not lost by simple pitch angle scattering, even when their nominal banana width is as large as the minor radius. While some losses can occur from "banana diffusion" of beam ions which have scattered into the trapped particle region of velocity space (*Ref. 3-33*), few co-injected ions are lost from the plasma, and the loss rate is always much more severe for counter-injection.

A final loss process that must also be considered has its origin in the discreteness of the toroidal field coil set, a feature that gives rise to local magnetic mirrors or ripple wells between the field coils. If beam ions are injected perpendicularly into these mirrors, or if the ions collisionally diffuse into the mirror-trapped region of phase space, they experience a unidirectional vertical $\mathbf{B} \times \nabla B$ drift which carries them directly out of the plasma. In TFR, with near-perpendicular injection, as much as 30% of the beam energy is calculated to be lost by this process (*Ref. 3-34*).

If devices are designed with small toroidal field ripple δ ($\leq 1\%$ peak-to-peak at the plasma edge) this loss process can be virtually eliminated. Ripple wells exist only if the gradient in B along the field line due to ripple exceeds that due to the rotational transform; i.e. only if $\alpha|\sin \theta| < 1$, where $\alpha = e/Nq\delta$, N being the number of toroidal field coils (*Ref. 3-35*). For cases with large values of α , ripple trapping can only occur near $\theta = 0$ and π , where the vertical $\mathbf{B} \times \nabla B$ drift only carries the trapped ions into a nearby region of the plasma where the wells disappear, and the ions are then detrapped back onto normal banana orbits.

3.2.3.2 Experiment — The process of beam-ion thermalization has been investigated in a large number of tokamak injection experiments — CLEO, ATC, ORMAK, TFR, DITE, T-11, and PLT. In general the results have shown very good agreement with the theory outlined above. Three separate diagnostic techniques have been used to measure the beam-plasma interaction: charge-exchange, neutron emission, and plasma diamagnetism, the latter in conjunction with the equilibrium position dynamics.

Among these diagnostics, the charge-exchange measurements give the most detailed information. They are made by aiming one or more neutral particle energy analyzers into the plasma at a variety of angles, in order to determine the time evolution of the flux of charge-exchanged beam neutrals as a function of time, energy and angle. With sufficiently detailed measurements it is possible to find the fast ion slow-down rate, the pitch-angle scattering rate, and the charge-exchange loss rate. In fact, in some experiments the beam-ion thermalization has been used to provide information about Z_{eff} and n_0 as well (*Ref. 3-36*). Further, the energy diffusion of beam ions, resulting from "heating" of the beam ions by the plasma, may be investigated by studying the high energy tail of beam ions. The tail arises from the fact that some few ions at first are heated by the plasma more than they are slowed down by it, and thus for a short period can experience a net acceleration, until the inexorable drag process pulls them down toward the ion thermal energy.

The first charge-exchange measurements of beam-ion thermalization were made on the CLEO tokamak (*Ref. 3-37*) in 1973. The fast neutral particle flux was measured as a function of time and energy at an angle close to the tangential injection angle. Results from the rise times of the signals as a function of energy below the injection voltage showed that the beam ions were slowing down at close to the classically expected rate. Some aspects of the results indicated deviations from classical predictions but such conclusions were not confirmed in later, more detailed investigations.

Soon after the CLEO measurements were made, a thorough study of beam-ion thermalization was carried out in the ATC experiment (*Ref. 3-38*). Short pulses of particles were injected into the plasma, and their slowing down and spreading in energy were monitored with a tangential charge-exchange analyzer. The steady-state spectrum near the beam energy and angle, and the

rise to equilibrium as a function of energy were also measured. Using a Fokker-Planck calculation of the fast ion distribution function that takes into account the radial variation of plasma parameters (rather than a homogeneous plasma model with averaged parameters as on CLEO), detailed agreement between experiment and theory was obtained. Co- and counter-tangential injection, as well as near-perpendicular injection were tested (Ref. 3-39), and the relative efficiencies of ion heating in the different geometries were in good agreement with beam-orbit Monte-Carlo calculations for the fraction of the total injected beam power delivered to the ions (Ref. 3-33).

On the basis of these favorable results, a series of fast ion diagnostic experiments were performed on ATC, using a tightly collimated neutral beam aimed outside of the magnetic axis, in order to deposit a local concentration of beam ions at a chosen minor radius. A multi-angle neutral particle analyzer was used to measure the initial orbit shifts and pitch-angle scattering rates of injected fast ions, and so to provide local measurements of q (Ref. 3-40) and Z_{eff} (Ref. 3-36). In addition these measurements gave very detailed evidence that the fast ion orbits and the thermalization process, at least near the injection energy and angle, were in excellent agreement with classical theory.

Measurements were also made of the fast neutral spectrum over a wide range of angles and energies in ORMAK. Perpendicular and parallel steady-state spectra were obtained over the entire energy range from the injection voltage down to the thermal ion spectrum. Comparison with numerical calculations that included toroidal effects in an approximate manner (Ref. 3-41) showed good agreement between experiment and theory if a Z_{eff} of ~ 4 was assumed, in agreement with spectroscopic measurements but much lower than the Z_{eff} deduced from the plasma resistance. Tangential steady-state fast neutral spectra on DITE also showed good agreement with theory, exhibiting no enhancement of the high-energy tail (Ref. 3-42). On TFR, the only experiment to date with extensive perpendicular injection, the measured perpendicular fast neutral spectrum shows qualitative agreement with theory (Ref. 3-34), but no attempt has as yet been made to compare it with a complete radially dependent Fokker-Planck calculation. This is largely because such a comparison is complicated by the distorting effect of the strong toroidal field ripple both on the measurement, and on the fast ion distribution itself.

On PLT, charge-exchange measurements were made over a wide range of angles, in this case providing a continuous scan from tangential to perpendicular (Fig. 3-3); these data give good agreement with the calculated spectrum assuming $Z_{\text{eff}} = 2$, as deduced from the plasma resistance. For the PLT case, where the banana width is much less than the plasma minor radius, and the toroidal field ripple is negligible, a formal expansion of the drift-kinetic equation could be made (Ref. 3-43), and the resulting bounce-averaged Fokker-Planck equation was solved numerically to calculate the fast neutral spectrum shown in Fig. 3-3. Fall times of the charge-exchange flux at the injection angle

Fig. 3-3 Comparison of theoretical predictions and experimental results for the fast-neutral angle and energy spectrum in PLT (Ref. 3-24).

after the neutral beam was shut off were also compared with theory as function of energy. The variation of the fall time with energy is a direct measure of the slowing-down rate, and it again exhibits agreement with theory as shown in Fig. 3-4. Finally on PLT, the high-energy tail, even at the level of 2.1 MW injection, and in the presence of enhanced density fluctuations, showed surprisingly good agreement with theory. This agreement can be viewed as a particularly good confirmation of the classical slowing-down theory since the high-energy tail should be set up by a delicate balance between beam-ions slowing down on the plasma, and plasma heating of the beam. In effect it gives a "temperature" of the heat sink absorbing the beam energy, and one might expect non-thermal turbulence in the plasma to exhibit a different effective "temperature" than the plasma itself.

The second technique which has been used to diagnose the beam-ion thermalization process is measurement of the neutron flux from D^0 injection into D^+ plasmas. Since the beam energy is much greater than the thermal plasma energy, and since the $D-D$ fusion cross section is a very strong function of energy, almost all of the neutron production is due to collisions between beam ions and bulk plasma ions, or, for co- plus counter-injection, from collisions of beam ions with each other. Measurements on ATC and ORMAK showed that the neutron production rate, and its rise and fall times, were consistent with simple calculations based on classical theory; these experiments provided the basis for the design of and commitment to the TFTR experiment. Since that time much more thorough calculations and measurements on the PLT tokamak have fortunately also shown favorable agreement. As shown in Fig. 3-5, theory and experiment tracked well together as beam and plasma conditions were varied to give neutron production rates varying over two orders of magnitude (Ref. 3-44). In addition, neutron emission rise and fall times for selected cases were compared with theory, and again gave good agreement. At the highest beam power deposited in PLT to date, emission levels of 1.5×10^{14} n/sec were measured, giving a total dose of 2×10^{13} neutrons. The "D-T equivalent" Q for this case is $\sim 2\%$.

Fig. 3-4 Comparison of theoretical predictions and experimental results for the $1/e$ fall time of tangential fast neutral flux after beam turn-off as a function of energy in PLT.

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Fig. 3-5 Comparison of theoretical predictions and experimental results for PLT neutron production (Ref. 3-44).

A final diagnostic technique for studying beam ion thermalization is the least precise, but the most straightforward. For this technique, one measures the plasma diamagnetic effect, which is proportional to $\beta_{\perp}$, and also the plasma position in a tokamak with a conducting shell, or the vertical field required for equilibrium in a machine without a shell. By combining the two measurements with measurements of T_i and T_e versus time and radius, one can determine the amount of $\beta_{\perp}$ and $\beta_{\parallel}$ unaccounted for by the plasma thermal energy, and therefore stored in circulating beam particles. For parallel (perpendicular) injection, one should find more unaccounted β in the parallel (perpendicular) direction than expected in an isotropic plasma, where $\beta_{\parallel} = \beta_{\perp}/2$. The anisotropy in the beam β , and its absolute magnitude, can in principle be compared with theory to give a good measurement of how well the beam energy is being confined over its entire history of thermalization. Measurements on T-11 (Ref. 3-45), where this technique was pursued most vigorously, and also on TFR, show gross agreement with theory, but uncertainties in profile information prohibit detailed comparison. Calculations for PLT with tangential injection indicate that the beam $\beta_{\parallel}$ and $\beta_{\perp}$ should be approximately equal, and should constitute a substantial fraction of the total β in low density, high-power injection discharges. In these cases the diamagnetic loop and vertical field data are again in qualitative agreement with theory, but detailed comparisons here are also problematic. Thus the magnetic measurements provide rather modest support for the neutron and charge-exchange diagnostics, but together these diagnostics have given extensive evidence in a number of experiments to show that beam-ion thermalization is proceeding very much as described by the classical theory of Debye-shielded collisions.

3.2.4 Plasma Heating

The first generation of neutral beam heating experiments on CLEO, ATC and ORMAK demonstrated modest percentage increases in ion temperature (up to $\sim 50\%$), by using beam powers somewhat smaller than the ohmic input power. While investigating the physics of beam-ion thermalization and demonstrating the effectiveness of neutral beam heating, these experiments laid the ground work for much more ambitious projects. The ORMAK, TFR, and DITE tokamaks then pushed their beam powers up to the few hundred kilowatt level, somewhat above the ohmic heating power, and obtained increases in ion temperature of 100 to 200%. On ORMAK and TFR, ion temperatures of ~ 2 keV were achieved. In addition, at this power level the first indications of electron heating were observed. In the recent experiments on PLT (late 1978) the beam power has been further increased to 2.5 MW, about a factor of four greater than the ohmic heating power without injection. Ion temperatures were increased by up to 600% to 6.5 keV, and electron temperatures rose by typically 50%, giving a maximum T_e of 4 keV with injection. Such plasmas represent the closest approach to reactor conditions of any fusion device to date and hold open great promise for the higher power neutral beam injection experiments planned on a variety of tokamaks.

The significance of these results lies partly in their demonstration that neutral beams effectively heat tokamak plasmas, but also in what they reveal about the complex questions of ion and electron thermal transport. In typical small tokamaks such as ORMAK and ATC, during ohmic heating the ions are held at $T_i \simeq T_e/3$ by collisional energy transfer from the electrons, with this energy transfer usually being only a small fraction of the total input ohmic heating power. Artsimovich showed that under these conditions the ion power balance could be decoupled from the electrons, since the energy transfer rate is insensitive to T_e , and scales as $T_i^{-1/2}$ (Ref. 3-46). As a result, as T_i is increased due to neutral beam power transfer to the ions (P_{bi}) the electron-ion coupling power (P_{ei}) falls. In a perturbation analysis (Ref. 3-47) appropriate for modest beam

powers, and taking the ion energy confinement time to be independent of T_i , the rise in ion temperature, $\Delta T_i(t)$, when the beam is turned on is given by

$$\Delta T_i(t) = P_{bi} \tau_{Ei}^* \left[1 - \exp(-t/\tau_{Ei}^*) \right] \quad (3-14)$$

where $\tau_{Ei}^* = (2/3) \tau_{Ei}$. In steady state, this implies $\Delta T_i/T_i = (2/3) P_{bi}/P_{ei}$. Experimentally, this relationship was found to hold at steady state on ORMAK, ATC, and TFR, implying that τ_{Ei} is approximately independent of T_i , as expected if charge-exchange and neoclassical thermal conduction are the dominant energy loss mechanisms. Verifying the time dependence of Eq. (3-14) is somewhat more difficult, however, because the finite beam ion slowing-down time must also be taken into account. On ATC (*Ref. 3-48*), good agreement with theory was found; however on TFR (*Ref. 3-34*), using a simplified model of the beam ion thermalization rate, the rise and fall times could not be made fully consistent with the calculated energy containment time.

It is more difficult to raise T_e by beam injection for a number of reasons. First, the neutral beam injected power must compete with P_{OH} , the ohmic heating power, which is much greater than P_{ei} (the relevant quantity for ions) under most circumstances. Secondly, for injection near E_C , most of the beam power is transferred directly to the ions. (Although if $T_i > T_e$, some or all of this can be subsequently transferred to the electrons from the bulk ions.) Finally, for fixed current operation, which is dictated by long magnetic diffusion times and MHD stability limits, P_{OH} is proportional to $T_e^{-3/2}$, and so falls rapidly as the plasma heats up. Using the same perturbation analysis, and ignoring ion-electron coupling yields

$$\Delta T_e(t) = P_{be} \tau_{Ee}^* \left[1 - \exp(-t/\tau_{Ee}^*) \right] \quad (3-15)$$

where $\tau_{Ee}^* = (2/5) \tau_{Ee}$, and in steady state this implies $\Delta T_e/T_e = (2/5) P_{be}/P_{OH}$. Thus it is not surprising that early injection experiments with $P_{be} \ll P_b \lesssim P_{OH}$ showed little or no electron heating due to neutral beam injection. Small density changes or changes in impurity content easily masked the direct injection effects.

At the high power levels achieved in the recent PLT and ISX-B experiments clear electron heating has been observed. However, while the perturbation analysis gives some insight into the physics involved, to extract the most information from the experiments, the full transport equations must either be solved to simulate the injection heating scenarios, or the experimental data must be unfolded to determine the transport coefficients. So far, these two approaches [with one notable exception (*Ref. 3-34*)] have given similar results (*Refs. 3-24, 3-48*). In the PLT experiment, where the most dramatic heating has been observed, the ion temperature increases were found to be consistent with code calculations which took into account neoclassical ion thermal conduction, charge-exchange, thermal convection, and electron-ion coupling as the ion energy loss mechanisms. It would have been desirable to deduce the actual ion thermal conduction coefficient from these experiments, but at high plasma densities electron-ion coupling dominates the power balance, and at low densities charge-exchange and thermal convection (which were not directly measured) appear to dominate; and as a result, the ion thermal conduction could only be shown to be within a factor of $\sim 3-5$ of the estimated neoclassical value, which is itself subject to some uncertainty in a $Z_{eff} \neq 1$ plasma.

The change in T_i was found to scale approximately as $P_b/\bar{n}_e$ on PLT (Fig. 3-6), as had been observed on TFR. At high beam power and low plasma density, resulting in $T_i \geq 4$ keV, strong density fluctuations were observed in PLT with microwave scattering techniques. These fluctuations appeared to be 5-10 times larger in amplitude than the usual turbulence found with the same scattering techniques. The frequency spectrum of the fluctuations exhibited a number of

evenly spaced distinct peaks, as opposed to the monotonic falloff with frequency observed normally. It is encouraging to note however that the scaling $\Delta T_i \propto P_b / \bar{n}_e$ observed in the experiments without fluctuations held up well into the regime where these as yet unidentified instabilities appeared. This indicates that the turbulence, while it may result in some enhancement of the ion thermal conduction, certainly does not radically alter the total ion power balance.

Electron heating in PLT was found to be very sensitive to the choice of limiter material, and to wall conditioning. With water-cooled large area carbon limiters, and with extensive titanium gettering, high-Z impurity radiation could be kept low in PLT with both co- and counter-injection. Under these circumstances, reliable electron heating was obtained. On ISX, using solely co-injection, glow discharge cleaning coupled with stainless-steel limiters, but without titanium gettering, also was successful in limiting impurity radiation (*Ref. 3-50*). In both experiments it was found that the observed electron temperature increases of up to a factor of two were consistent with the same electron energy containment time as found in the ohmic heated case. Detailed profile studies revealed, however, that τ_{Ee} must increase in the plasma core to explain this, since the beam input power is much less peaked at the plasma center than is the ohmic power input. See Section 2.2.4.2 for a discussion of the form of the electron thermal transport coefficients deduced from the PLT experiments.

3.2.5 Other Aspects of Neutral Injection

In addition to providing heat to the plasma, neutral beams also serve as sources of particles and of momentum. Furthermore, energetic beam ions and charge-exchange neutrals escaping from the discharge column can also constitute an important increase in the rate of bombardment of the vacuum vessel wall, with a resulting rise in the impurity influx. Each of these effects has been observed with neutral injection, and some limited insight has been gained thereby into the underlying processes of particle confinement, toroidal momentum damping, and impurity production. In practice, however, these subjects have been less thoroughly investigated than the questions of ion and electron energy transport, and — perhaps as a result — they are less well understood.

In the earliest neutral beam experiments the plasma density was seen to increase with injection by an amount much larger than the direct fast particle influx. On ATC, and in the early PLT experiments without titanium gettering, the enhancement was a factor of 2 or greater. The extra particle influx was attributed to cold gas streaming from the beamlines, and to increased recycling at the plasma boundary due to enhanced wall bombardment. In more recent experiments on PLT, using extensive titanium gettering on the vacuum vessel wall to reduce the recycling of neutrals at the plasma edge, the plasma density was observed to fall monotonically during a normal tokamak

Fig. 3-6 Ion temperature increase in PLT versus beam power per particle (*Ref. 3-23*).

pulse, and with the addition of neutral injection, the density increase could be almost entirely explained by the fast particle influx. Indeed the density increase on axis appeared to be even somewhat smaller than expected from the local fast-ion deposition. With injection into plasmas whose density was largely sustained by gas puffing ($\bar{n}_e \geq 2.5 \times 10^{13}$), it was possible in PLT to hold the density constant during injection, simply by reducing the gas feed. To make a flat, reproducible density trace before, during, and after injection, it proved very helpful to connect the piezo-electric gas feed valve and the fringe-counting 2 mm microwave interferometer in a feedback loop. Interestingly, the density rise is observed to cease above a given injection power level (~ 600 kW) in ISX-B unless certain characteristic MHD activity occurs, in which case the density rise may be sustained. When $\bar{n}_e$ rises during injection into low density PLT plasmas, a strong central peaking of the electron density increase is observed, corresponding qualitatively to a peaked beam deposition profile. While little insight into the mechanisms of particle transport has as yet been gained from this experience, it has at least been demonstrated that the techniques which have proven useful for controlling plasma density and neutral gas recycling in ohmically-heated discharges are also very effective during neutral injection.

In a toroidally symmetric system such as a tokamak, conservation of toroidal angular momentum plays a central role in neoclassical transport processes. Tangential neutral injection, because it is uniquely capable of providing an input of momentum, makes it possible to directly investigate how well this system (which ideally has an ignorable coordinate) conserves angular momentum. Three processes are expected to contribute to the damping of toroidal rotation: neoclassical perpendicular viscosity, friction against ripple in the toroidal field, and charge-exchange. In practice the first two processes are calculated to be quite weak, damping rotation on a timescale longer than the PLT discharge duration. On the other hand, the damping time from charge-exchange in low density experiments is often comparable to the ion energy containment time, although the toroidal momentum containment time due to charge-exchange, τ_ϕ^{CX} , must always be greater than τ_{Ei} , since charge-exchange contributes to ion energy loss as well as to the loss of angular momentum, but other processes also contribute significantly to ion energy loss.

In the early experiments on ATC, only very modest toroidal rotation (5×10^5 cm/sec) was deduced from the doppler shift of $m = 1/n = 1$ MHD oscillations in the soft X-ray flux emitted from the plasma center (*Ref. 3-51*). This corresponded to a momentum containment time much less than τ_{Ei} . On ORMAK, a rotation speed of 10^7 cm/sec was deduced from the doppler shift of the far wings of the tangentially viewed H_α radiation, but this apparent rotation persisted in high Z_{eff} discharges even in the absence of neutral injection (*Ref. 3-52*). On PLT, a rotation speed of $\sim 10^7$ cm/sec was also observed, in this case by three techniques: soft X-ray measurements of the doppler shift of $m = 1/n = 1$ fluctuations, comparison of the parallel and anti-parallel charge-exchange spectra, and the doppler shift of impurity line radiation originating from the plasma center (*Ref. 3-53*). This rotation speed corresponded to a momentum confinement time in PLT somewhat less than the ion energy containment time. Thus it is evident that toroidal angular momentum is lost by some anomalous process, since classically $\tau_\phi > \tau_{Ei}$, while experimentally $\tau_\phi < \tau_{Ei}$ is found. These measurements are in their infancy, however, and much more detailed investigations are needed to deduce the magnitude and scaling of the anomalous viscosity that appears to damp toroidal rotation in tokamaks.

With regard to impurity contamination, it has been found on both ORMAK and PLT that impurity radiation is dramatically enhanced by counter-injection, but only modestly increased (even decreased in some cases) by co-injection. This difference is consistent with theoretical considerations concerning fast ion orbits (see Section 3.2.3.1) which show that 5–20% of the counter-injected power can be lost to the limiter on bad orbits in typical experiments, while only a very small fraction of co-injected power is expected to strike the limiter in the form of escaping beam

ions. A surprising result, however, observed on both ORMAK and PLT, is that co- plus counter-injection brings in fewer impurities than counter-injection alone. This might suggest that the momentum transfer from the beams to the impurities is affecting impurity transport in the plasma, as would be expected from neoclassical calculations taking into account the presence of an unspecified anomalous damping mechanism for toroidal rotation (*Ref. 3-54*). On the other hand, a clear correlation has been observed on PLT between the edge ion temperature (as measured by the doppler broadening of low ionization potential carbon and oxygen lines) and the high-Z impurity radiation from the plasma core. This correlation persists with injection, showing a higher edge ion temperature with counter-injection than with co- plus counter-injection or with co-injection alone. The persistence of the correlation, however, cannot be taken as a proof of causation; a rearrangement of the current profile due to cooling of the plasma core by high-Z radiation could, for example, result in an increased edge ion temperature. Thus the difference between the impurity influx with co- and with counter-injection must be regarded as an outstanding question.

3.3 RADIO FREQUENCY WAVE HEATING

Plasma heating through the absorption of energy from radio frequency (rf) plasma waves encompasses a wide range of options, corresponding to a range of possible wave frequencies from around 1 MHz or below ($\omega \sim v_A/a$) to 100 GHz and above ($\omega \sim \omega_{UH} = \sqrt{\omega_{pe}^2 + \Omega_e^2}$). The options available within this range offer some important potential advantages for fusion reactor applications. From the technological standpoint, large amounts of rf power can be provided at high efficiencies in the frequency range up to a few GHz. In fact, the hardware required for a number of proposed rf heating schemes for reactor grade tokamaks is already state of the art. Moreover, several possible wave heating approaches, if successful, would permit deposition of this power deep in the plasma interior under circumstances where neutral beam penetration would necessitate very high beam energies. And particular wave-heating schemes can be used to heat selectively either the plasma ions or electrons, either the perpendicular or parallel degrees of freedom, and either the bulk or the tail of the particle distribution. In fact, the associated wave momentum deposition seems well suited to driving the plasma current required for confinement in tokamaks, making possible a steady-state tokamak free of the encumbrances needed for inducing the conventional toroidal electric field. Finally, wave heating offers alternative approaches to supplying auxiliary power to the plasma through coupling structures designed to avoid the neutron streaming problem inherent to the use of the large neutral-beam ducts which give direct line-of-sight connection between the plasma and the source/accelerator structures. Waves with frequencies above roughly 1 GHz can be launched using waveguide arrays of convenient size. Lower frequency waves must be driven by rf coils for which voltage handling requirements could pose insulator problems in reactor environments.

In practice, wave heating in tokamaks has been demonstrated for most of the recognized approaches. These will be described below in greater detail. Despite these demonstrations, however, none of these approaches has been applied routinely to a major tokamak facility. Furthermore, the identification of the underlying physical mechanisms and the detailed comparison of theory and experiment have not reached a stage that permits a confident extrapolation to the reactor regime for any of the techniques under consideration. Rather, wave-heating experiments are still primarily devoted to the resolution of problems in one or more of the following general areas:

1. **Launching:** Investigations of those processes which affect the launching of the desired wave or the excitation of the desired cavity eigenmode, and especially the excitation of unwanted waves.

2. **Propagation:** Identification of parameter ranges within which the coupling of the desired plasma wave to other parasitic modes is negligible and propagation of the pump wave into the plasma core is possible.
3. **Absorption:** Analysis of the physical processes by which the waves are absorbed, either in the interior or at the plasma surface, and the resulting changes in the plasma properties.

The present status of wave heating for tokamak plasmas is most easily assessed by considering each of the major approaches separately. For this purpose we identify the following distinct methods, ordered by increasing frequency:

1. Alfvén-wave heating
2. Magnetosonic wave heating
3. Ion cyclotron harmonic heating
4. Lower hybrid wave heating
5. Electron cyclotron harmonic heating.

For each of these approaches we shall summarize briefly the present situation with respect to the three critical areas identified earlier; namely, those processes critical to the launching, the propagation, and the absorption of the plasma waves. Technological considerations will be discussed if they appear to pose a serious constraint on application of a particular method to reactor-relevant tokamaks.

3.3.1 Alfvén-Wave Heating (*Ref. 3-55*)

1. Excitation of torsional (shear) Alfvén waves with $k_{\perp} \gg k_{\parallel}$ is achieved through modulation of the poloidal magnetic field and has been experimentally demonstrated on linear and stellarator-like devices. Relatively large antenna structures are required, in close proximity to the plasma; and some design problems may be encountered in minimizing image currents in the vacuum chamber and avoiding excitation of magnetic islands in the plasma.
2. Because the wave spectrum is continuous no discrete toroidal eigenmodes are excited, and suitable wave propagation is generally available for plasma parameters of interest. Unwanted compressional magnetosonic waves can be suppressed by proper choice of parallel wavelengths: $\lambda_{\parallel} > 2a$.
3. Absorption of the wave energy occurs as the result of a localized resonance, $\omega_0 = k_{\parallel} v_A(r)$, at which the wave is converted to kinetic Alfvén waves with perpendicular wavelengths comparable to the ion gyroradius. The wave energy is ultimately absorbed by Landau damping of these secondary kinetic Alfvén waves. Extensive theoretical studies of these absorption processes have been carried out but little direct experimental evidence is yet available.

3.3.2 (Low-Frequency) Magnetosonic Wave Heating ($\omega \ll \Omega_i$) (*Refs. 3-56, 3-57*)

1. Excitation of compressional magnetosonic waves is achieved through modulation of the toroidal magnetic field using simple antennas coupled to the lowest toroidal eigenmodes.

2. A lower limit on the density and size of a suitable tokamak plasma is expressed through the condition for existence of suitable toroidal eigenmodes:

$$n_e a^2 \geq 8 \times 10^{15} \text{ cm}^{-1} \left(\frac{\Omega_D}{\omega} \right)^2 \quad (3-16)$$

for a 50–50 DT plasma. If the driving frequency is below the gyrofrequency of singly ionized oxygen ($\omega/\Omega_D \lesssim 8$), the acceleration of surface impurities is minimized. The corresponding size/density condition becomes

$$n_e a^2 \geq 5 \times 10^{17} \text{ cm}^{-1} \quad (3-17)$$

Thus, this approach is best suited to large, high-density plasmas such as TFTR, JET, etc.

3. The absorption mechanism expected to be important in this type of wave heating is collisionless absorption through the combined mechanisms of electron transit time and Landau damping. Thus, the parallel phase velocity of the wave must be maintained near the electron thermal speed.

3.3.3 Ion Cyclotron Harmonic Heating

1. Magnetosonic waves at higher frequencies can be launched with loop antennas located near the plasma surface. Coupling to toroidal eigenmodes can significantly reduce the currents or voltages which the antennas must sustain for a given power absorption within the plasma, and thereby permit the use of higher resistivity antenna materials. Systematic investigations of launching structures are currently in progress with promising results (*Refs. 3-57, 3-58*).
2. The cavity modes become closely spaced at these frequencies in moderate to large plasmas; and excitation of a single selected mode becomes technologically more difficult. Even so, satisfactory mode tracking approaches have been successfully used on TFR (*Ref. 3-58*). The ideal quality factor for these cavity modes, Q , can be very large; but problems have been encountered in attempts to heat at the second harmonic of the deuterium gyrofrequency. A small admixture of hydrogen leads to a remarkable reduction in Q , apparently as the result of a linear mode conversion. Nonetheless, recent experiments in the PLT device (*Ref. 3-59*) have apparently demonstrated efficient heating of a deuterium plasma as the result of direct heating of a minority concentration of protons at their fundamental gyrofrequency, followed by collisional heating of the majority deuterons by these energetic protons. In these experiments, the protons, estimated to comprise only 7% of the plasma density, are heated to temperatures of several keV by 70 kW of rf power at 25 MHz. The deuteron temperature is increased by more than 100 eV by thermalization of the fast protons. Further experiments will attempt to extend this method to higher rf power levels and to determine the dominant wave properties necessary for successful heating. Early problems experienced in the ST tokamak (*Ref. 3-60*) in which the discharge was caused to disrupt through the influx of cold gas from the wall have been controlled in the more recent experiments.

3. The dominant absorption mechanism in this frequency range is collisionless cyclotron damping. This process does not depend on Landau-type wave-particle resonances and efficiently heats the bulk of the plasma ions. As seen above, heating of a minority species can yield very high energy groups of ions analogous in many respects to neutral beam heating.

3.3.4 Lower Hybrid Wave Heating (Ref. 3-61)

1. Lower hybrid waves can be launched in a tokamak plasma using various types of slow-wave structures which provide external fields that satisfy an accessibility condition,

$$n_{\parallel}^2 > 1 + 2 \frac{\omega_{pe}^2}{\Omega_e^2} - \frac{\omega_{pi}^2}{\omega^2} + 2 \frac{\omega_{pe}}{\Omega_e} \left(1 + \frac{\omega_{pe}^2}{\Omega_e^2} - \frac{\omega_{pi}^2}{\omega^2} \right)^{1/2} \quad (3-18)$$

Lower hybrid waves launched in this way have been detected directly in Doublet IIA and in Alcator A. From the point of view of eventual reactor applications, phased arrays of waveguides offer a promising approach to launching these waves with dynamic control of the $n_{\parallel}$ spectrum.

2. The lower hybrid waves propagate throughout the plasma if the wave frequency lies above the local lower hybrid resonance

$$\omega_{LH}^2 = \omega_{pi}^2 \left(1 + \frac{\omega_{pe}^2}{\Omega_e^2} \right)^{-1} \quad (3-19)$$

and if the plasma density is below a critical value at which linear mode conversion can occur:

$$n_{crit} = \frac{\epsilon_0 B^2}{m_e} \left[\sqrt{3} n_{\parallel} \frac{v_{th,e}}{c} \left(1 + \frac{4}{\hat{\omega}^4} \frac{T_i}{T_e} \right)^{1/2} + \frac{1}{\hat{\omega}^2} - 1 \right]^{-1} \quad (3-20)$$

where

$$\hat{\omega}^2 = \frac{\omega^2}{\Omega_i \Omega_e} \quad (3-21)$$

The efficiency with which external power can be coupled to the lower hybrid wave appears to be high but the experimental results needed to evaluate this quantitatively are not yet available for high power experiments. A crude estimate of coupling efficiency derived from the relative amount of reflected power can approach unity.

3. The wave energy can be absorbed either by electrons, through electron Landau damping, or by ions, through mode conversion processes near the lower hybrid resonance, or by parametrically driven modes. Both ion and electron heating have been observed experimentally; and in the electron heating cases, the

experiments are in reasonably good agreement with predictions of a transport analysis based on quasilinear Landau damping.

3.3.5 Electron Cyclotron Harmonic Heating (*Ref. 3-62*)

1. The waves in this frequency range may be conveniently launched using waveguide/radiators of various types. Most tokamak experiments to date have used waves launched from the low-field side of the device at an angle to the magnetic field.
2. The wave propagation in tokamak plasmas has been analyzed extensively using cold-plasma models. Effective propagation is available provided the density remains below a value corresponding to cutoff:

$$n_{\max} \equiv 1.24 \times 10^{12} \text{ cm}^{-3} \left(\frac{\nu_{\mu}}{10 \text{ GHz}} \right)^2 \times \begin{cases} 1 + \frac{1}{m} & , \text{ extraordinary} \\ 1 & , \text{ ordinary} \end{cases} \quad (3-22)$$

where ν_{μ} is the wave frequency and m is the harmonic number. The waves can propagate in such a way that linear mode conversion occurs at the upper hybrid resonance

$$\omega^2 = \omega_{pe}^2 + \Omega_e^2 \quad (3-23)$$

3. The wave energy can be absorbed through cyclotron damping of the incident wave or cyclotron and Landau damping of secondary waves resulting from mode conversion. Efficient absorption has been demonstrated experimentally. For heating plasmas with densities above 10^{14} cm^{-3} frequencies above 100 GHz are required, and high-power sources are not yet available in this regime. The gyrotrons developed recently in the Soviet Union and now available in the U.S. have resulted in a substantial increase in the pace of technological development. One drawback to ECH is the large cost per unit of power delivered; this has its origin in the advanced technology required and the low efficiency ($\sim 30\%$) of power generation.

3.3.6 Scaling to Reactor Size Tokamaks

The neutron environment, the greater plasma size, density and temperature, and the reliability and maintenance problems associated with reactor operation present challenges for rf heating technology that go beyond those found in present day tokamaks. Several systems studies of rf heating of reactor tokamaks have been performed, including studies employing ion cyclotron, lower hybrid, and electron cyclotron heating.

Fast magnetosonic wave heating at $\omega = 2 \omega_{cD}$ has been examined (*Ref. 3-63*) within the context of the NUWMAK (*Ref. 3-64*) conceptual reactor design. By use of low azimuthal mode number $|m| \leq 3$, and a narrow but moderate $k_{\parallel}$ spectrum, $k_{\parallel} \approx 20 \text{ m}^{-1}$, mode conversion to Bernstein modes is avoided and only a small number of eigenmodes are excited. As the plasma density changes, the alternative eigenmodes are excited but no mode tracking is required. Due to

the moderate wave Q and low wave impedance in a reactor, neutron-resistant antenna structures are required to provide an impedance transformation from the high impedance transmission lines to the plasma. Applying 80 MW of rf power for 3 seconds is sufficient to bring the plasma to ignition. Overall electrical efficiency of the heating system is estimated to be $\approx 60\%$.

Lower hybrid heating via electron Landau damping has also been examined as a possible means of heating to ignition (*Ref. 3-65*). In this approach, a waveguide grill structure (*Ref. 3-66*) is employed for launching the lower hybrid waves, and the overall electrical efficiency is again about 60%. As the plasma approaches ignition, penetration of the lower hybrid waves to the plasma center is prevented by a combination of limits due to (quasilinear) electron Landau damping and due to mode conversion from the slow lower hybrid mode to fast waves. Utilizing a programmed density rise such as may be provided by gas puffing and/or pellet injection, the central region of the plasma is directly heated in the early phase of the approach to ignition, after which fusion alpha heating begins to dominate the energetics in the plasma center and the rf provides bulk heating.

For heating to ignition via ECRF (*Ref. 3-67*), important constraints are imposed by the ratio of plasma to cyclotron frequencies and the maximum frequency of rf power that can be produced at acceptable efficiencies. Ordinary mode injection from the outboard side of the plasma chamber at a frequency equal to the cyclotron resonance frequency in the plasma center is considered to be the most favorable approach. However, the condition $\omega_{pe} < \omega_{ce}$ must be satisfied to avoid cutoff of this mode before it reaches the plasma center. This condition is difficult to satisfy in typical heating-to-ignition scenarios if Alcator scaling is applicable. Thus, it may be necessary to either launch from the inboard side or to use the second harmonic. The inboard side of the first wall presents severe access problems and the latter places a substantial burden on technological development.

3.4 OTHER HEATING TECHNIQUES

In addition to neutral-beam injection and the various rf techniques discussed in the preceding sections, a number of other methods for heating tokamak plasmas have been proposed. Of the seven methods listed in Table 3-1, four have actually been attempted in various tokamak plasmas, while the others have been applied to alternative types of confinement devices. This section presents a brief synopsis of these alternative heating methods including both the performance achieved to date in experimental tests and the likely future prospects.

3.4.1 Adiabatic Compression Techniques

A tokamak plasma can be compressed either by increasing B_T in time with the plasma at fixed R (minor-radius compression), or by increasing the vertical field at fixed $B_T(r)$, thus displacing the plasma inward in R to a new equilibrium at higher B_T (major-radius compression) (*Refs. 3-68, 3-69*). Defining the compression factor C as B_{new}/B_{old} , conservation of magnetic moment results in $W_{\perp} \propto C$, while conservation of angular momentum results in $W_{\parallel} \propto C^2$. Ohmic heating of the precompression plasma can be supplemented by neutral beam injection or rf heating, or indeed by any method that proves effective. Because the compression process can in principle be carried out at high efficiency, the efficiency of the supplementary heating system becomes a less critical issue.

Compression must be carried out in a time τ_c significantly smaller than the plasma energy confinement time, τ_E . Major-radius compression has the great advantage over minor-radius compression that the former process involves the pulsing of magnetic energy comparable with that in

TABLE 3-1
METHODS OF PLASMA HEATING OTHER THAN NEUTRAL BEAMS OR RF

Heating Method	Tokamak Experiments Performed	Major Results	Reactor Physics Outlook	Reactor Technology Prospects
1. Major-Radius Compression	ATC	$T_i \propto C^{4/3}$, $T_e \propto C$	Excellent	Very good
2. Minor-Radius Compression	TUMAN TOSCA	$T \propto C^{4/3}$, enhanced τ_E (both machines)	Very good	Poor
3. Turbulent and Shock	TTT Various small tokamaks	$\Delta T_e \sim 0$, energetic ion tails $\Delta T_e \sim 0.5$ keV Energetic ion tails τ_E maintained	Poor	?
4. Relativistic Electron Beam	SPAC V	Injecting and trapping $\Delta I_p \sim 30$ kA $T_i \sim 300$ eV $q \ll 1$	?	Good
5. Cluster Injection	None		?	Fair
6. Plasma-Gun Injection	None		?	Good
7. Laser-Pellet Hot Plasma Formation	None		?	Fair

the poloidal field, while the latter process demands the pulsing of magnetic energy comparable with that in the toroidal field, an energy content some 100 times larger. Compression in R is actually an effective method of increasing B_T at the plasma without requiring an increase in toroidal magnetic energy.

3.4.1.1 Major-Radius Compression — Compression in R has been carried out on the ATC tokamak with $C \leq 2.3$ (Ref. 3-70). Both T_i and n_e increased as expected ($T_i \propto C^{4/3}$ and $n_e \propto C^2$), but T_e increased only as C. This behavior is consistent with $\tau_{Ei} > \tau_c \gtrsim \tau_{Ee}$, the smaller τ_{Ee} probably being due to radiation loss (Ref. 3-71).

On the basis of the present scaling laws for τ_E , ohmic heating plus compression gives little advantage over the case of ohmic heating alone in the position of the compressed plasma (Ref. 3-72). The advantage of R-compression is realized when it is used to multiply the heating delivered by neutral beams or rf to the precompression plasma. The lower line density of the precompression plasma is beneficial to penetration of neutral beams and, combined with the lower magnetic field, to several forms of rf as well.

Good heating results were obtained on ATC with 230 kW of 15 keV D^0 beam injected into the precompression plasma to produce $T_i = 400$ eV, followed by an R-compression of $C = 2.2$ that increased this temperature to 1200 eV (Ref. 3-73). This "compression boosting" method can be applied to produce an ignited plasma formed by compression of a beam-heated plasma whose opacity $\bar{n}_a p$ is determined by neutral-beam (or rf) penetration into the center (Ref. 3-74). Upon compression $\bar{n}_a p$ increases by a factor $C^{3/2}$ to the value required for ignition. A value $C = 1.6$

is adequate to bring E_b into the range of TFTR-type injectors (120 keV) (*Ref. 3-74*). In order to exploit the available magnetic real estate, the cold plasma surrounding the hot ignited core can be heated by a propagating thermonuclear burn driven by alpha power deposition in the ignited core (*Ref. 3-75*). Thermal stability control of ignited or high-Q plasmas can also be achieved by regulating the plasma temperature with alternating compression and decompression (*Ref. 3-76*).

Further experimental investigation of R-compression will be carried out in TFTR and in JET. A significant technological requirement for a reactor is the large pulsed power (~ 500 MW) that must be provided for the poloidal field system. Nevertheless, because R-compression is a proven and well understood heating method that can multiply the energy input of neutral beams or rf, and because of the prospect of using it for thermal stability control, this technique may be of particular value in an ignition test reactor (*Ref. 3-77*).

3.4.1.2 Minor-Radius Compression — The effectiveness of minor-radius compression has been demonstrated in the TUMAN devices ($C \leq 3$) (*Ref. 3-78*) and in TOSCA ($C \leq 2$) (*Ref. 3-79*). In these experiments the expected temperature and density increases ($T \propto C^{4/3}$, $n \propto C^2$) were observed and the plasma energy confinement time was enhanced. The latter effect has been attributed to the increased physical separation of the plasma from the walls, thus reducing the evolution of impurities.

For future application, minor-radius compression appears to have three serious disadvantages: First, a large fraction of the toroidal field energy must be changed in a time smaller than τ_E , which is impractical in the largest present-day normal-coil tokamaks, and certainly out of the question for reactors with superconducting coils. Second, the aspect ratio of the compressed plasma is rather large, which is unfavorable for obtaining high values of beta (see Section 2.1.1). Third, the final volume of the compressed plasma is considerably smaller than that of the vacuum vessel, and worse, fails to exploit the available high- B_T region. In the case of an ignited compressed plasma, however, the surrounding cold plasma could be heated by a thermonuclear burn.

3.4.2 Turbulent Heating

This method, which was one of the most popular conceptual approaches around 1970, involves superimposing a very strong pulsed electric field on the usual toroidal field of the tokamak. To initiate strong turbulence, the pulsed field must be sufficiently large to produce a relative drift velocity between ions and electrons that exceeds the electron thermal velocity (*Ref. 3-80*). The resulting electrostatic instabilities are supposed to heat the electrons until their thermal velocity exceeds the relative drift velocity. The ions are subsequently heated both by instabilities and by Coulomb coupling to the electrons. Rapid deposition and thermalization of energies up to 1 kJ have been demonstrated by this method in small ($a \lesssim 10$ cm) moderate-density ($\lesssim 3 \cdot 10^{13} \text{ cm}^{-3}$) toroidal plasmas, where keV electron and ion energies have been produced but with confinement times much less than 1 ms (*Refs. 3-81, 3-82*). Energetic ion tails near the plasma surface have commonly been observed. The nature of the heating mechanism is only poorly understood in practice.

Two important questions for the application of turbulent heating to reactor-sized plasma are whether the rapidly rising electric field can penetrate the plasma, particularly as T_e increases from the ohmic-heating value, and whether the normally good plasma energy confinement will be degraded by the turbulence. In a large plasma the bulk of the heating seems likely to occur at the edge (*Ref. 3-83*), with subsequent energy loss to the wall and inefficient heating of the central plasma region. Because of the limited amount of energy in each pulse, repetitively applied pulses of turbulent heating would be necessary.

Perhaps the largest tokamak on which the turbulent heating method has been attempted was the TTT device at the University of Texas, where an 80-kA current was induced in less than 1 μ sec (*Ref. 3-84*). Little electron heating was observed, although energetic ion tails could be produced. As a result of this unfavorable experience, no further experiments are planned in the U.S.; however, investigations continue in Europe and in Japan on a number of small toruses.

3.4.3 Relativistic Electron Beam (REB) Injection

In principle, tokamak plasmas can be heated by the injection and trapping of relativistic electron beams (*Ref. 3-85*), which may also be used to help drive the plasma current (see Section 4.3). A beam of 5 to 10 MeV electrons is injected in the toroidal direction by a transmission line that penetrates to the plasma edge, and the beam diffuses radially inward on successive transits of the torus. Part of the beam energy is transformed into the magnetic energy of the electron "return current," and the resistive decay of this return current heats the bulk plasma. The bulk plasma is also heated by slow momentum transfer from the injected fast electrons.

In experiments in linear devices it has been found that good energy transfer from the REB is effected by electrostatic instabilities driven by the relative velocity between the relativistic electrons and the electrons making up the return current (*Ref. 3-86*). It is not known whether such instabilities would have a significant effect on energy confinement if they were to arise during REB injection into a tokamak.

The most important REB experiment to date in a toroidal device is the injection of 26 kJ of a 0.5-MeV REB into the SPAC V device at Nagoya ($a_p \sim 10$ cm, $B_T \sim 5$ kG) (*Ref. 3-87*). Penetration to the center was achieved, and $T_i = 300$ eV at densities approaching 10^{14} cm $^{-3}$ was obtained with plasma stability apparently retained for the experimental time of 1 ms. Critical issues that remain to be addressed experimentally are whether the REB can penetrate into the center of a large, dense tokamak plasma, whether the fast electrons give up all their energy before they escape the plasma, and whether instabilities produced by REB injection will degrade the energy confinement time.

The technology for REB generation and transmission is well advanced. At least 2 MJ per pulse has been produced at efficiencies in excess of 65% (*Ref. 3-88*). But the feasibility of using the transmission line injection method in a reactor environment is uncertain.

It seems likely that REB injection will be tested in medium-sized tokamaks during the next several years. If successful, this method may be considered seriously for use in reactor-sized tokamaks, at least for its ability to aid in driving the plasma current.

3.4.4 Cluster Injection

Large neutral condensates (10^3 to 10^7 molecules) can be formed by the expansion of hydrogenic gases at low temperature. These condensates are then broken up into clusters containing 50 to 500 molecules, ionized, and accelerated to 1 MeV or more (*Ref. 3-89*). The gyroradius of a cluster must be sufficiently large to obviate the neutralization process that allows particles to cross the magnetic field lines of the tokamak. The clusters are broken up by ionizing collisions while traversing the target plasma.

The energy to which the clusters must be accelerated is determined by the velocity needed for at least part of the cluster to penetrate to the plasma center, and also by the fraction of plasma fueling that is performed by the cluster. For a reactor-sized plasma, at least 60 keV per atom is needed, which implies cluster energies of more than 10 MeV (*Ref. 3-90*). Systems at Fontenay-aux-Roses and at Karlsruhe have reliably generated clusters of acceptable size and accelerated them to

1 MeV, which is adequate for present-day machines. Injection into a tokamak has not yet been attempted, but is planned for ASDEX (*Ref. 3-91*) and possibly for TFR. An important potential problem to be investigated is premature breakup of the clusters before reaching the target plasma.

Cluster injection has the advantages over neutral beam injection that at high cluster energies no neutralization should be required, so that production efficiency can be very high, and steady-state operation is readily achieved. But the apparatus is extraordinarily complicated, and it is difficult to see how the formidable energy requirement for heating a reactor-sized plasma can be satisfied.

3.4.5 Plasma-Gun Injection

Plasma guns of the Marshall or deflagration type provide reasonably pure plasmas with energies per ion up to about 10 keV (*Ref. 3-92*). Energy efficiency can be as large as 90%, and up to 10^{21} particles per shot have been generated, so that only 10 or 20 shots would be required for initial heating of a reactor plasma. To date plasma guns have been used to produce relatively cold target plasma for toroidal and other devices ($T_i \lesssim 100$ eV) (*Ref. 3-93*), but no experiments to generate hot tokamak plasmas have been attempted.

It is uncertain theoretically and unknown experimentally whether the injected hot plasma blob can penetrate a tokamak plasma of reactor level na. An important technological problem is to insure that essentially no high-Z impurities from the gun electrodes will enter the gun-injected plasma.

3.4.6 Laser-Pellet Hot Plasma Formation

If cryogenic pellets of H, D, or T are suspended in the torus and bombarded by sufficiently large pulses of laser energy, then the resulting hot plasma would rapidly spread out to fill the entire torus (*Ref. 3-94*). The toroidal plasma current would have to be initiated extremely rapidly to confine the hot particles before they escape to the wall. The method thus seems more suitable for toroidal devices where the rotational transform is applied externally, such as stellarators or multipoles. This method has in fact been attempted on multipoles, but only very low temperatures have been attained (~ 10 eV), partly due to the small laser energy deposited (*Ref. 3-95*).

For reactor application the chief technological drawback to this method is the need to obtain laser pulses of at least 1 MJ at efficiencies of several tens of percent. However, there is no requirement for the rapid pulsing needed in the laser-pellet inertial confinement fusion schemes.

3.4.7 Summary

The critical issues for the various alternative heating approaches are presented in Table 3-2. Of the techniques discussed here, only adiabatic compression has been demonstrated to work in a tokamak of significant size with significant energy confinement time. And compression appears to have practical benefit only when employed in conjunction with some other non-ohmic heating method. At this time it seems likely that REB injection will be tested in one or more small tokamaks, but mainly to investigate its potential for driving the plasma current. Turbulent heating experiments will proceed in small toruses in Europe and in Japan. But major-radius compression is the only "other heating method" that is presently under serious consideration for use in the next generation of large tokamaks, or incorporated into the design of conceptual tokamak reactors.

TABLE 3-2
CRITICAL ISSUES FOR ALTERNATIVE HEATING TECHNIQUES

Technique	Issues
1. Major-Radius Compression	Large poloidal field system power consumption Compression of noncircular plasmas
2. Minor-Radius Compression	Feasibility of rapid increase in toroidal field energy
3. Turbulent Heating	Penetration time of applied fields Effect of turbulence on plasma confinement
4. REB Injection	Penetration of REB to the central plasma region Extent of thermalization of REB energy
5. Cluster Injection	Acceleration to at least 10 MeV Premature disintegration of clusters
6. Plasma-Gun Injection	Penetration of the hot plasma blob Impurity content of injected plasma Instabilities induced in the target plasma
7. Laser-Pellet Hot Plasma Formation	Laser energy per pulse Laser efficiency Poor confinement without ohmic current

4. TOKAMAK DESIGN

The present generation of tokamak experiments and those tokamaks presently under construction are expected to increase the data base in those areas deemed especially critical for assessing the prospects of the tokamak as a reactor candidate. In some cases, this requirement takes the form of a specified level of performance, such as the ability to access the collisionless regime or to achieve energy breakeven, and in other cases, it may be to perform the experiments required to advance our understanding to the level required for design of a subsequent device in the program, examples being the exploration of divertor operation or the effects of plasma shaping. In either instance, the demand that a given device address such critical issues, a list of which is to be found in Section 1.2, has fostered a new level of sophistication in tokamak design. A number of the associated design issues are dealt with in this chapter.

According to present evidence, which is predominately theoretical in character, the key to reaching beta values of reactor interest lies in the ability to accurately control plasma position and shape. This is achieved through careful design of the magnetics system. Section 4.1 addresses a number of the design issues associated with this system, which also plays a vital role in startup and shutdown and perhaps in impurity control as well. It is important to recognize that energetics considerations demand that a reactor magnetics system consist largely of superconducting coils, in which case a whole new set of design constraints enters. This, together with the problems associated with operation of any required normal coils in a high neutron flux environment, is a dominate consideration for reactor design.

Another key design issue, that of impurity control, is discussed in Section 4.2. A small admixture of high- Z impurities can dominate the energy balance, and relatively modest concentrations of lighter ions can contribute significantly to Z_{eff} and to the depletion of the reacting fuel in a tokamak reactor. Further, there is the question of the helium ash produced by the fusion reaction itself. The most straightforward solution to this problem is to divert the field lines near the plasma edge to an auxiliary chamber. A number of possibilities for divertor design exist and the disturbingly large extent to which most of these options complicate the overall device design dictates a thorough investigation of these and other alternatives to the impurity control problem.

Acceptable reactor economics demands discharges which are, at the minimum, tens of seconds in duration. Section 4.3 deals with some of the issues involved in long pulse operation, both through a characterization of the time scales for the relevant physical properties and through discussions of fueling and non-ohmic current drive, the latter being of particular interest in promoting the possibility of steady state operation.

The means by which information is deduced about the plasma is the subject of Section 4.4. In designing to meet a specified objective (e.g. ignition) that requires significant extrapolation beyond current experience, it is important to reduce as much as possible the imprecision of existing data and this translates into demands on diagnostic systems.

With continued success, the tokamak program will culminate in power generating devices. The current status of the concept of tokamak reactors, reviewed in Section 4.5, reflects both the increased thoroughness of recent reactor design studies and a more intensive effort to incorporate state of the art understanding of tokamak physics into a feasible technological framework.

4.1 MAGNETICS OF THE DISCHARGE CYCLE

The temporal and spatial dependence of the poloidal magnetic field in a tokamak determines the plasma current, the plasma position, and the shape of the minor cross-section. This poloidal field is controlled by two or more poloidal field coil systems. One of these coils, called the ohmic heating or induction coil, produces the majority of the volt-seconds, or magnetic flux, required to induce and sustain the plasma current. This induction coil, the subject of Section 4.1.1, consists of solenoidal windings between the plasma and the axis of symmetry of the torus and additional non-solenoidal windings to eliminate stray fields in the plasma region which would inhibit plasma initiation and complicate positional control. Ideally, the individual turns of the induction coil should lie on a single flux surface. This can be achieved either by careful choice of the number and location of individual turns (in a series wound coil) or, for the case of a superconducting induction coil, as will likely be required for reactor operation, by simply connecting the turns in parallel.

The remaining coil systems, termed here the VF (vertical field) coil systems, dealt with in Section 4.1.2, provide additional volt-seconds. These systems provide the vertical field necessary for equilibrium, i.e., the magnetic field necessary to balance the forces arising from magnetic and thermal pressures. The shape of the vertical magnetic field determines the gross radial and vertical stability of the plasma as well as the plasma shape. Vertical fields with a simple curvature can be used to produce vertical or horizontal elongations of the plasma. More complicated vertical field shapes are required to produce doublets or poloidal field divertors. In contrast to the induction coil, the VF-coil current distribution is indeterminate, since any current distribution that produces no field in the plasma region (such as that characterizing the induction coil) can be added to the VF-coil with no effect on plasma shape or position. This gives the poloidal field designer the flexibility to choose the fields outside the plasma region and to optimize power supply design. In some instances, it is useful to take advantage of this flexibility to choose the sum of the VF-coil currents to be equal and opposite to that of the plasma current. Hence, if the VF-coils are inside the TF-coils, the TF-coils are effectively shielded from flux changes in the plasma such as those accompanying startup or disruptions. This is an important consideration in the case of superconducting TF-coils, a necessity for reactor design. A further attractive design feature is a VF-coil set which has zero effective mutual inductance with the induction coil; this may be achieved by connecting the coils in parallel and by appropriate choice of the mutual inductances.

The induction coil may be either air core or employ a high permeability material such as iron to limit the energy necessary to produce the flux through these regions and to reduce the stray fields in the facility. These advantages of using iron must be balanced against several disadvantages. First, the variation of the permeability of the iron with field strength complicates the electric behavior of the system. In more advanced high field devices the iron will be driven into saturation, which not only complicates the electrical behavior, but also produces field error effects. If the iron is driven into saturation, its energy saving characteristics are also degraded. Secondly, the interaction of the plasma induced magnetic field with the iron reduces the stability of the system to radial motion, an observation confirmed by experience on TFR 600 (*Ref. 4-1*); on the other hand, ISX has shown that proper poloidal field design can produce a stable configuration with an iron core but no copper shell (*Ref. 4-2*). Third, the use of iron reduces the space available for tokamak systems and tokamak access. Different choices may be optimal for different applications; e.g., TFTR and JET employ air core and iron core induction coils, respectively.

4.1.1 Induction Coil and Plasma Startup

In order to produce the flux swing necessary to induce and sustain the plasma current, a time changing current is passed through the induction coil. If the key design consideration is the limiting stress on the coil, operation proceeds via a full swing by reversing the current from a pre-biased negative current condition. If, on the other hand, stored energy is the limiting design feature, driving the coil from a negative bias to zero may be more appropriate. For reactor applications, the induction coil will be superconducting so the former mode of operation will apply.

Among the considerations that dictate power supply design are:

1. The voltage requirement determined by the voltage necessary to initiate the plasma discharge and induce the plasma current (*Ref. 4-3*).
2. The maximum current requirement determined by the maximum volt-second capability of the induction coil system.

In general, a capacitor bank or inductive interrupter (*Ref. 4-4*) are used to provide the high voltage capability for startup. For some applications, it is more convenient to apply this plasma initiation pulse to a set of VF-coils. Either a low voltage capacitor bank or controlled SCR system (*Ref. 4-5*) may be used to produce the voltage necessary to sustain the plasma discharge during current flat-top. Some relief to the volt-second requirements during the flat-top may be achieved if proposed current drive schemes prove feasible; see Section 4.3.3 for a discussion of these possibilities.

As the size of a tokamak increases, the constraints imposed on the induction coil system to initiate the discharge become quite severe. The constraints are of two types: those limiting discharge parameters and those imposed on the coil and power supply design. In the quasi-steady state, the energy input per unit volume from ohmic heating scales roughly as R^{-2} at fixed B_T and q . The energy density required to support ionization and line radiation losses is size-independent but scales unfavorably with increasing density. If the power input is limited to ohmic heating during the startup phase, this results in a maximum attainable density for a given applied loop voltage (*Ref. 4-3*).

The startup of a large tokamak typically goes through three phases:

1. Breakdown of the background gas.
2. The period during which the power loss by ionization of the background gas and radiation of low Z impurities dominates the gross energy balance.
3. The period during which the current profile evolves into its quasi-steady state.

Initial breakdown of the background gas is typically accomplished by inducing a high frequency electric field or by use of an ultraviolet source. In the latter case, avalanche breakdown of the background gas is accomplished by the induced voltage of the induction coil system. As pointed out by several authors (*Refs. 4-6—4-8*), the particle confinement time determined by particle drifts during this stage must be larger than the ionization time for discharge initiation to occur. The resulting constraints on the one-turn voltage are typically less than those required to fully ionize the background gas.

The energy balance during the ionization phase of the discharge has been incorporated into both zero- and one-dimensional transport codes (*Refs. 4-3, 4-9, 4-10, 4-11*). Experiments on PLT (*Refs. 4-12, 4-13*), as well as other machines (*Refs. 4-14, 4-15*), have shown that the energy replacement time during the startup phase is typically much shorter than that in the quasi-steady state. The short energy replacement time may be in part due to rapid plasma transport. However, the power loss by charge exchange, ionization of the background plasma, and impurity radiation is

typically very important. If the filling pressure is too high for a given applied one-turn voltage, the discharge will not initiate (*Refs. 4-3, 4-12*). The technique used to achieve high densities on nearly all tokamaks is to operate at modest initial filling pressure where the discharge will initiate readily and then program the flow rate of gas. This technique has been further extended on PLT by using a gas feedback system together with titanium gettering between shots to control the rate of density buildup. Furthermore, the maximum density achievable depends upon the low Z impurity concentration in the discharge. Because the electron temperature is low during the startup phase, low Z impurities radiate very effectively and are very important in the power balance. Recent advances in discharge cleaning, especially those proposed in *Ref. 4-16*, have resulted in a reduction of the low Z impurity content under typical operating conditions. Nonetheless, low Z radiation may be important in the startup of future devices, particularly if low Z limiters and wall coatings are extensively used; the latter may be a necessity to achieve reactor level temperatures in high density devices since high Z radiation is an enormous energy sink at temperatures of several keV.

The energetics of a PLT startup are shown in Fig. 4-1, from *Ref. 4-13*, where the electron and oxygen densities and the one-turn voltage are chosen to correspond to the maximum operating conditions for PLT. The lines of fixed E^2/η and ηj^2 represent the constraints imposed by ohmic heating, the former signifying the maximum achievable one-turn voltage and the latter representing a q limit, in this case $q = 3$.

On the basis of both extrapolation from present experiments, and the use of zero and one dimensional transport codes, the maximum density range attainable during startup using ohmic heating can be determined. In order to minimize the one-turn voltage requirements, and hence the coil insulation and power supply requirements, supplementary heating during the startup phase using neutral beam injection (*Ref. 4-7*) and rf heating (*Ref. 4-17*) has been proposed. The difficulty with neutral beam injection is that in order to confine the beam injected ions, a sufficient plasma current must exist. Since the one-turn voltage requirements are determined during this stage in the discharge, there is very little improvement with the addition of neutral beam injection. Furthermore, owing to the low electron temperature, ohmic heating during the startup phase is typically less expensive than either rf or neutral beam heating; this has represented an obstacle for experimentation to demonstrate the feasibility of the alternative methods.

Current penetration during discharge initiation has been studied both experimentally and theoretically on a number of devices (*Refs. 4-11, 4-12, 4-18, 4-19*). The current penetration in PLT was modeled by a one-dimensional magnetic field diffusion code using neoclassical resistivity. The code used the measured temperature and

Fig. 4-1 The power loss by charge exchange and oxygen line radiation is estimated for PLT and compared with the constraints imposed by ohmic heating.

density profiles as well as current and voltage measurements in order to evaluate Z_{eff} . The resulting calculations of Z_{eff} during the startup phase have, under certain circumstances, given results which exceed those of the spectroscopic measurements by a factor ~ 4 , indicating the presence of an enhanced current penetration mechanism. Probe measurements on the LT-3 tokamak have shown that a rapid relaxation ($\sim 200 \mu\text{sec}$) of the current channel occurs after the onset of $m = 4$ and $m = 1$ helical instabilities (Ref. 4-18). The startup phase of T-10 has been modeled with some success by introducing MHD modes to relax the current channel (Ref. 4-11). It has also been suggested (Ref. 4-19) that MHD activity during the startup phase may explain the observed scaling of the maximum density as B_T/R (Ref. 4-20) owing to the empirical observation that the ohmic input power, when MHD activity begins, is proportional to the electron density.

A rapid rise to full plasma current limits the resistive volt-second investment in startup and thus permits either a longer discharge duration or operation of the induction coil at reduced stress. But such operation may place unreasonable demands on the power supplies since a fully energized reactor-size induction coil stores on the order of 1 gigajoule of energy. A further consideration applicable to reactor design is the existence of an upper bound on the rate of change of field in a superconductor. Recent measurements indicate that a $\dot{B}$ of order 11 tesla/sec is acceptable for a NbTi solenoid operating at a peak field of 4.5 T (Ref. 4-21). Finally, as discussed in Section 2.1.2.3, there is strong evidence that a gradual rise in plasma current permits operation at lower $q(a)$, which may in turn lead to higher achievable β .

4.1.2 Vertical Field and Shaping

The magnitude of the vertical field required to balance the forces resulting from magnetic and thermal plasma pressures was determined by Shafranov (Ref. 4-22) and Yoshikawa (Ref. 4-23). The vertical field shape determines both the shape of the plasma minor cross-section and the stability of the plasma to radial and vertical perturbations (Ref. 4-23). Horizontal and vertical elongations of the plasma can be produced by the addition of a quadrupole component to the vertical field (Ref. 4-24) but such an addition generally reduces the stability of the plasma to displacements. Highly conductive boundaries in the form of copper shells have been used to improve the plasma stability as well as to provide equilibrium fields. Following ATC's successful operation without a copper shell (Ref. 4-25), most advanced tokamak designs have eliminated this feature. A hexapole component of the vertical field produces plasma triangulation, or a dee-shaped plasma. Computer modeling of the plasma stability indicates that triangulation can be used to decrease the susceptibility of the discharge to ballooning modes (Ref. 4-24).

The doublet configuration is the result of applying a complex poloidal field geometry to the plasma equilibrium (Ref. 4-26). The doublet configuration has been shown to allow stable discharges with high total plasma current in devices with modest toroidal field (Ref. 4-27). Theoretical investigations indicate that the doublet configuration may be capable of supporting discharges with higher total plasma betas than simpler geometries.

Divertor geometries are also the product of a complex magnetic configuration. Such an approach shows promise for controlling the plasma outflux, separating the neutralized gas from the discharge, and reducing the impurity influx. The two most widely adopted designs, the bundle divertor and the poloidal field divertor, are discussed briefly below; they are dealt with in greater detail in Section 4.2.

The bundle divertor is a toroidally asymmetric magnetic configuration, with the advantage that it can be accommodated without significant increases in the area inside the toroidal field coils and the vacuum vessel (Ref. 4-28). The bundle divertor has the disadvantages that it increases the

complexity and field errors associated with the magnetic topology and also requires a large scrape-off layer. The major problems in coil design revolve about the ohmic losses, the constraints on field ripple and the increased stress level in the toroidal field coil system.

In contrast to the bundle divertor, the poloidal field divertor is axisymmetric but requires a large volume inside the toroidal field coil and vacuum system (*Ref. 4-29*). Poloidal field divertors connect naturally to the separatrix formed by a dee-shaped plasma with one divertor above the upper null and the second below the lower null but, with closely fitting divertor coils, are compatible with more complex geometries as well (*Ref. 4-30*).

The proximity of the VF coils to the plasma is a matter of considerable importance. Shaping and stability control are clearly enhanced by a close fitting coil set, but the energy economics in a reactor is unacceptable if the coils lie between the plasma and the energy conversion medium. This problem is exacerbated for plasmas of more complex shape but one possible solution is the use of a double layer coil set consisting of an inner coil layer, thin enough to be virtually transparent to neutrons, which provides the magnetic boundary for plasma control, and an outer coil layer, situated behind the blanket and perhaps even outside the toroidal field coils, which is programmed to minimize the currents in the inner coil layer (*Ref. 4-30*).

4.1.3 Fault Problems

Protecting the vacuum vessel and poloidal field coil system against faults which result from plasma disruptions produces severe design constraints on these systems. The tokamak must be protected against the transfer of plasma energy by three mechanisms. The vacuum vessel walls and plasma limiters must be protected against the transfer of thermal energy from the plasma during disruptions. The poloidal field coils and power supplies must be protected against voltage spikes and/or the transfer of magnetic energy into the poloidal field coil systems. Both the coils and the vacuum vessel must be protected against the forces arising from the large currents induced during plasma disruptions.

No design effort would guarantee that the vacuum vessel, or limiter, could absorb the thermal energy of the plasma with the local concentration that appears to occur during a disruption, without sustaining some structural damage. At best the limiters can be designed to protect the vacuum vessel and absorb the plasma energy with modest local melting and sublimation.

The time constant for plasma disruptions plays an important role in determining the protection required against magnetic energy transfer. The phenomenon cannot occur more rapidly than the transit time for the discharge at the Alfvén velocity in the poloidal field; for reactor size devices, this time constant is typically 10–100 microseconds. If large plasma currents were extinguished at this rate, the voltages induced in the poloidal field coils would be extremely large, particularly if the vacuum vessel has a ceramic insulator, which would eliminate the electrical protection provided by the vacuum vessel for the poloidal field coil systems. It should be noted that although continuous vacuum vessel designs ameliorate this problem as well as the development problem associated with an insulating break that can tolerate a high neutron fluence, the current induced in a continuous vessel gives rise to a severe (~ 10 atmosphere) pressure pulse which may be a life-limiting feature of reactor designs (*Ref. 4-31*). Fortunately, experiments on PLT indicate the typical disruption time exceeds 1 millisecond, while disruptions with time constants less than a millisecond are relatively rare occurrences.

The approximate voltage seen by an open circuited poloidal field coil will be the typical single turn voltage multiplied by the difference between the number of positive and negative turns. Coil protection can be gained via a continuous vacuum vessel design or by increasing the vacuum vessel

conductivity at the expense of the poloidal field penetration time. Two additional methods can be used to protect the poloidal field coil system. The first method is to provide a fault protector across the individual coil pancakes which will conduct if sufficiently high voltages occur. In addition, the maximum turn to ground voltages can be minimized by interleaving coil bundles with positive and negative voltages.

The vacuum vessel and its internal structure must have sufficient strength to withstand the forces resulting from the interaction of the induced current during disruption with the applied external fields. Of particular importance is the interaction of the saddle currents, resulting for example from a neutral beam port, with the toroidal field.

4.2 IMPURITY CONTROL

The necessity for maintaining plasma purity has long been recognized as a fundamental concern for tokamak operation. The presence of impurities makes startup and heating more difficult and degrades the power density in a reactor grade plasma. Generally speaking, plasma cleanliness has improved markedly in recent tokamaks as increased operating experience has provided information, much of it empirical in character, about wall preparation and about the choice of materials that comprise the plasma-matter interface. But the problems to be faced in burning plasmas, with the associated increased particle and energy fluxes, longer burn times, and ash removal concerns, will keep impurity control in the forefront of tokamak research.

The requirements on plasma purity, and thus on the effectiveness of an impurity control system, vary substantially for the different phases of the discharge and for different plasma conditions. For example, the permissible impurity concentration during the early ($T_e < 50$ eV) stage of the discharge is determined by the ohmic heating power density (and thus the maximum one-turn voltage) since line radiation from low-Z impurities generally dominates the energetics in this temperature range. High-Z impurities are a greater concern at higher temperatures ($T_e \gtrsim 500$ eV) because the low-Z impurities are fully stripped; there are thoroughly documented instances in which radiation from high-Z impurities provides such a large heat sink that further increases in electron temperature via auxiliary heating become impractical. For a burning plasma, the major concern is that impurity accumulation, whether generated via surface effects or via the production of helium, will result in a premature termination of the burn.

A variety of approaches have been suggested to alleviate the impurity control problem. Some of these approaches have not yet been subjected to meaningful experimental tests and those that have been studied in some detail in the laboratory have been explored under widely divergent operating conditions. The succeeding discussion will thus of necessity focus on the dimensions of the problem rather than elaborate on a particular proposed resolution.

4.2.1 Plasma-Wall Interactions

Interactions between the plasma and the surrounding structure are of considerable complexity and remain incompletely understood. For a review of the current status of the subject, see *Ref. 4-32*. The dominant surface effects appear to have their origin in the photons and energetic particles escaping the plasma and in the space charge initiated unipolar arcing. In a reactor, the neutrons carry off the majority of the fusion energy; although they do contribute to wall heating and are responsible for bulk irradiation damage, they are implicated to only a minor degree in wall erosion. The small neutron sputtering yields may still have significance for safety and maintenance by contributing to the spread of radioactivity into coolant fluids.

The importance of plasma-wall interactions was recognized early (*Ref. 4-33*) and was among the first fusion technology areas to be pursued vigorously. The effort has been international in scope and has become increasingly sophisticated. In particular, the increasing availability of and emphasis on measurements on operating tokamaks has proven very beneficial and has provided considerable guidance for focusing laboratory experiments.

One area of prime concern is the degree of cleanliness of the surfaces exposed to the plasma. In a high vacuum, this is normally set by outgassing, as least for quasi-steady-state operation. Practice therefore dictates system baking at elevated temperature to purge the dissolved/adsorbed gases to a level where the residual outgassing rates are acceptable. Because release rates have an approximately exponential dependence upon temperature, reactors, which must operate routinely at elevated temperatures, will require greater attention to wall pretreatment than do present devices.

The degree of care required to achieve adequate wall cleanliness depends on the life cycle of impurity atoms. Simple recycling through the plasma redeposits the impurity on the wall surface, to be re-released in a subsequent discharge. Actual removal from the chamber can occur only by divertor action or by vacuum pumping. Outgassing rates are very dependent on the choice of material as well as the fabrication and preparation treatments, and relatively little data is available at elevated temperature. Rates of about $10^{15} - 10^{16}$ molecules/m²-sec have been measured for metals and low-Z materials operating at their maximum outgassing temperatures of 400–500°C (*Ref. 4-34*). At these rates an impurity concentration of a few percent could accumulate in the plasma during a 1000-second burn. Such a rate of accumulation may be tolerable assuming relatively efficient purging from the chamber; if the purging is inadequate, reduced outgassing rates are required. These rates can be reduced by orders of magnitude, but at the expense of very high temperature pretreatment (*Ref. 4-35*). Pretreatment prior to construction can be employed, but it is likely that some short term over-temperature capability will be required to be designed into the plant as well. The temperature differential required will depend on detailed specifications of materials and operating modes.

Vaporization of wall or limiter materials, such as may be caused by excessive local power deposition, can pose problems for both plasma purity and the structural integrity of the device. Control measures are of course reduced energy loading or improved cooling design. Alloying elements or metallic impurities may be more volatile than the host metal but the specifics depend sensitively on materials selection and pretreatment.

As suggested above, the wall surface can act as a reservoir for adsorbed gaseous impurity molecules. Adsorption energies can be so large that some impurities will not be removed by any pretreatment procedures. However, plasma photon and particle spectra are sufficiently energetic to lead to stimulated desorption of these impurities. Reduced plasma chamber surface-to-volume ratios in power reactors will diminish the importance of this problem and, depending on the effective purge rate from the chamber, it may be eliminated altogether except for initial startup or following extended down times. In any case, "discharge cleaning" is a process now well in hand (*Ref. 4-16*) and alternate approaches (*Ref. 4-36*) could prove to be even more effective.

Momentum transfer from impacting energetic particles (hydrogenic species, helium, or impurity atoms) may also result in sputtering of surface atoms of the wall or limiter material itself. This is a process long recognized (*Ref. 4-37*) and has received a good deal of emphasis recently for systems of interest for fusion, covering impact energies of 0.1–10 keV. Sputtering yields are generally in the range 10^{-3} to 10^{-2} atoms/ion and can generally be reproduced in different laboratories to within about a factor of two. Values for different target materials are usually within about an order of magnitude. The much larger radiative loss ratios for the ions of these materials in the plasma gives a clear preference for low-Z materials for interfacing with the plasmas unless particle

energies can be kept below a few tens of eV, i.e., below sputtering thresholds for the higher Z materials. Energy dependences at these low energies are less clearly defined, and the dependence on the incident angle to the surface needs clarification. Continued measurements of the sputtered species are also needed for input to detailed transport calculations. Recent results with computer calculations of sputtering yields (*Ref. 4-38*) for light ion projectiles have shown improved correlation between theory and experiment.

Surface temperature generally has little or no influence on sputtering yield. An important exception is carbon, for which an enhancement of about an order of magnitude in carbon erosion rate is observed due to the formation of methane; in this case, sputtering peaks in the range 400–600°C (*Ref. 4-39*). Boron and silicon carbides have also displayed increased erosion at elevated temperature (*Ref. 4-40*).

Plasma particle impact energies with the first wall are sufficient for implantation to a depth ranging from angstroms to micrometers. Super-saturation due to chemically active atomic hydrogen may also occur. A steady state recycling of plasma ions will be developed, including some directly reflected at the surface. The storage capacity for tritium must be considered in inventory and safety considerations and recycling of these atoms into the plasma edge must be accounted for in model calculations.

These implanted atoms may also be responsible for additional modes of wall erosion. Even in the presence of atomic hydrogen alone, small blisters have been observed to form by hydrogen gas recombination within the crystal lattice, apparently leading to sloughing off of wall metallic particles (*Ref. 4-41*). Accumulation of the implanted energetic particles trapped relatively deep below the surface may also lead to surface exfoliation and blistering, to a depth comparable to the implantation depth. Repetition of this process can lead to large effective erosion rates, in some cases generating equivalent sputtering ratios in excess of unity. The possible occurrence of blistering depends on many considerations, prime among them of course being the details of the high energy particle spectrum impinging the wall (*Ref. 4-42*).

Finally, unipolar arcing was the major concern at the very initiation of consideration of fusion technology (*Ref. 4-43*), and interest in it has now come full circle. Renewed investigations have begun and the known characteristics were reviewed at a recent workshop (*Ref. 4-44*), but no definitive conclusions have been drawn. It is conceivable that the quantity of wall and limiter material removed by arcing in operating tokamaks accounts for all the metal impurities observed in the plasmas. Whether this is indeed the source is unknown because of the lack of information on the fraction that in fact penetrates into the plasma, a question that is relevant for essentially all considerations of impurity problems in tokamaks.

To summarize, the experimental investigation of the various plasma-wall interactions is an active field with useful progress being made. Important work still remains to adequately understand and quantify the significant processes, and to expand experimentation into areas exploring possible synergistic influences. Detailed theoretical and empirical modeling of plasma edge effects and the subsequent impact on global plasma dynamics is needed to cast the results to date into a form with predictive capability.

4.2.2 Impurity Radiation

The significance of impurity radiation in tokamak plasmas has been well established. Typically, moderate-Z impurities such as carbon and oxygen can contribute significantly to the plasma resistivity, and can give rise to substantial amounts of radiated power from the edge region of the

discharge where such impurity ions are incompletely stripped. High- Z impurities such as molybdenum, tungsten, and gold are typically not present in sufficient quantities to cause significant contributions to the plasma resistivity, but the radiation from these ions, especially from the hot central part of the discharge, can be very substantial since they are incompletely stripped even at temperatures of a few keV.

A number of calculations have been performed to quantify the impurity radiation losses for specified impurity concentrations using detailed information about individual transitions for all the charge states involved (*Ref. 4-45*). These calculations determine coronal equilibrium values for the emission from steady-state solutions of a set of coupled rate equations which include the effects of electron collisional ionization and excitation, radiative and dielectronic recombination, $\Delta n = 0$ and $\Delta n \neq 0$ line transitions, and bremsstrahlung. Because of the uncertainties of the atomic physics and the large number of excitations that must be taken into account, this approach has been limited in practice to low- and moderate- Z elements and, until recently, only order of magnitude estimates have been available for such materials as molybdenum and tungsten, substances which have subsequently turned out to be important factors in the energy balance in the core of many tokamak discharges.

A recent development has been the use of a statistical approach, the average-ion model, to compute radiative losses (*Ref. 4-46*). The results for moderate- Z ions, such as oxygen and carbon, and for iron agree very well with other calculations, which adds to the confidence in the results for the higher- Z ions. For high- Z ions, e.g. tungsten, the radiation is found to be substantially higher than that predicted by previous order-of-magnitude estimates. This appears to be due primarily to the effect of dielectronic recombination, which lowers the equilibrium degree of ionization substantially below that given by radiative recombination alone.

Figure 4-2 shows the total radiated power as a function of electron temperature for a variety of impurity ions (*Ref. 4-46*). The presence of radiating impurities can have a substantial effect on the conditions for ignition, or for power amplification, in a reacting plasma. Figure 4-3 shows the $n_e \tau_E$ requirements to achieve ignition, or various Q -values, in a 10 keV D-T plasma as a function of impurity concentration, for tungsten and iron impurities (*Ref. 4-47*). Here Q is defined as

Fig. 4-2 Radiated power as a function of electron temperature T_e for a variety of impurities. The impurity density is n_Z and the electron density n_e .

the ratio of fusion power to injection power, taking fusion power as 17.6 MeV/reaction, and assuming (in the driven cases) that the plasma temperature is maintained by 200 keV deuterium beams in addition to the alpha-particle heating.

The effect on energy balance of low- Z impurities, at least in moderate concentrations, may well be beneficial because the radiated power is strongly peaked at low temperatures. Thus the central plasma parameters are not significantly degraded while the radiation at the plasma edge provides an effective mechanism for transferring the plasma energy benignly to the wall; this may act as a control mechanism for the production of high- Z impurities. However, it has been observed that the accessible operating regime is diminished

with increasing Z_{eff} (see Section 2.1.2, in particular Fig. 2-5). Thus low- Z impurities may exacerbate the problem of producing the densities needed for ignition-level $n\tau$. Further, even a modest amount of low- Z impurities may make startup more difficult (see Section 4.1.1) since a large energy density is required to "burn through" the low ionization states. In this case, a fully dynamic (non-coronal) treatment is needed to accurately model the physical situation. For burning plasmas, the key issue is not radiation but rather the resulting dilution of the reacting fuel. The reaction rate (or equivalently, the fusion power) for fixed n_e and T_i is reduced by a factor

$$\left(\frac{Z_i - Z_{\text{eff}}}{Z_i - 1} \right)^2 \quad (4-2)$$

for a single impurity of charge Z_i in comparison to the impurity-free case. For the case of carbon and $Z_{\text{eff}} = 2$, this factor is 0.64, which translates into a 50% penalty in the cost of electricity.

4.2.3 Non-Divertor Methods of Impurity Control

A number of tokamaks have proven capable of producing clean plasmas without resorting to the use of a divertor. The keys to this success appear to be some combination of choice of wall and limiter materials, surface treatment of those materials, and care in plasma positioning, especially during startup. Although the physics of the plasma boundary is not sufficiently well understood to permit confident extrapolation to the reactor regime, it seems a certainty that impurity control will become more difficult at the higher temperatures, higher wall loadings, and longer burn times

Fig. 4-3 Requirements on $n_e \tau_E$ to achieve ignition, or various Q -values, in a plasma driven by a 200 keV deuterium beam, as a function of the concentration of iron and tungsten impurities.

characteristic of reactor operation. Such operation must cope not only with the generation of impurities at the plasma edge but also the inevitable manufacture of helium from the fusion reaction itself.

Section 4.2.3.1 attempts to quantify the extent of the impurity control problem and the degree to which the problem can be alleviated by choice of first wall material. Subsequent sections address a number of possibilities to modify the physics of the plasma edge to either enhance the radial diffusion of impurities or to encourage radiative cooling, thus limiting impurity generation.

4.2.3.1 Impurity Accumulation — The most technologically attractive structural materials for fusion reactors are composed primarily of medium- and high-Z elements. Unless the plasma edge temperatures can be kept very low or the charged particles emerging from the plasma are diverted to impact another surface, the sputtering rates from such materials combined with the radiative energy loss rates exhibited in Fig. 4-2 are sufficient to preclude a sustained burn.

Since the great heat loads associated with plasma disruptions in large tokamaks suggest the use of a sacrificial surface to protect the structural integrity of the first wall (Section 2.1.3), it is natural to consider constructing this surface, whether it be a coating or a structure that stands off from the wall, of a substance that has more attractive impurity production properties. Such an analysis reveals that low-Z substances such as beryllium, carbon, and perhaps silicon carbide, are good candidate materials. For such materials, impurity accumulation is a gradual process, permitting plasma burns of order one minute or longer, the exact value depending on the plasma edge temperatures and, in the case of carbon where chemical sputtering is an important concern, the surface temperature as well (*Ref. 4-48*).

The burn may also be quenched by the fuel depletion associated with the retention of helium. In what follows, a 0-D model is employed both to compute the pulse length under the assumption of perfect recycling and to deduce the recycling rate required to produce a steady state burn (*Ref. 4-49*).

The assumptions made in the calculation are $n_D = n_T$, fixed n_e , no fast alpha losses, $Z_{\text{eff}} = 1$, parabolic $n(r)$ and $T(r)$, Alcator scaling, and an unspecified burn control mechanism which acts to maintain a constant β [thus preserving the relation $\partial(nT)/\partial t = 0$]. Power balance reads

$$P_\alpha = P_T + P_R + P_{BC} \quad (4-2)$$

where the subscripts refer to the α -power, the transport losses, the radiation losses (assumed to be purely bremsstrahlung), and the losses due to the burn control mechanism.

As the helium builds up, P_α decreases and when n_α reaches $n_{\alpha, \text{max}}$ the burn will be extinguished by virtue of P_{BC} being reduced to 0. The burn pulse length may then be calculated by integrating the α production equation,

$$\frac{\partial n_\alpha}{\partial t} = n_D n_T \langle \sigma v \rangle_{D-T} - \frac{n_\alpha (1 - R)}{\tau_{p, \alpha}} \quad (4-3)$$

forward in time until $n_{\alpha, \text{max}}$ is obtained. Here, R is the fraction of helium leaving the plasma that recycles and $\tau_{p, \alpha}$ is the α particle confinement time (actually, this is the effective recycling time, which may be considerably less). Figure 4-4 illustrates results for a typical ignition-size plasma ($a = 1.2$ m) for $R = 1$, i.e., perfect recycling. It is seen that ash accumulation limits the burn to tens of seconds assuming perfect recycling or, equivalently, perfect α particle confinement.

Steady state solutions are possible provided α particles leave the system sufficiently rapidly. The required magnitude of the effective recycling parameter $(1-R) \tau_{E,DT}/\tau_{P,\alpha}$ for a steady state burn falls in the narrow range (0.025, 0.04) for T between 5 and 10 keV and $4\% < \beta < 8\%$.

These results are very model dependent, and a meaningful 1-D simulation is difficult owing to the large uncertainties associated with particle transport. One general conclusion that can be drawn, however, is that it is advantageous to increase the ratio $\tau_{E,DT}/\tau_{P,\alpha}$ as well as the sticking coefficient $1-R$. There is some indication that the former may occur naturally due to the presence of a ripple loss cone so the requirements on the sticking coefficient may be relatively modest.

Assuming the alpha particle confinement time (or recycling time) is no more than three times the ion energy confinement time, it follows that if a fraction $1-R \gtrsim 10\%$ of the alpha particles leaving the plasma does not recycle, a steady state D-T concentration sufficient for ignition will be retained. One possibility to achieve the needed effective sticking probability without using a divertor is to employ as a first wall or as a coating on the first wall a material which acts as a getter for helium. Although this possibility has not yet been thoroughly explored, many substances can absorb helium in the needed quantities for periods on the order of minutes while operating in the temperature range of interest for fusion reactions; in other temperature ranges, these same materials will purge the helium giving rise to a potentially attractive operational cycle (Ref. 4-50).

4.2.3.2 Impurity Flow Reversal and Gas Blankets — The theoretically predicted (Ref. 4-51) alteration of impurity transport by poloidally asymmetric injection of hydrogen gas has been seen in titanium gettered discharges in ISX-A (Ref. 4-52), albeit for times too short to gain insight into the steady state situation. The existence of this effect, which is predicted by neoclassical theory, coupled with the agreement between measured and neoclassically predicted transport times (Ref. 4-53), argues that theory can be used to extend the concept to the edge of a reactor plasma.

More gas driven flow reversal tests are planned on ISX-B and Doublet III in FY 1980. In addition, tests are planned on ISX-B to attempt to verify the theoretically predicted effect of neutral beam produced momentum on impurity transport (Ref. 4-54). The momentum transferred to the plasma by means of traveling rf waves is another attractive possibility, but no tests of this are presently planned.

Fig. 4-4 Length of burn of a typical ignition-size plasma as constrained by helium accumulation assuming perfect recycling. The solid curves denote Alcator scaling for the energy confinement time while the dotted curves represent a factor of two improvement over Alcator scaling.

One promising means of achieving the required poloidal asymmetry, thus driving the flow reversal, is the use of gas recycling from a toroidal limiter at the top or bottom of the chamber since the gas throughput needed is typically comparable to the recycling rate. However, the use of this scheme in a reactor context does raise the issue of the ability of limiters to withstand the intense heat fluxes anticipated in that environment. Establishing the proper polarity for whatever fueling technique is employed is another means of enhancing the flow reversal.

To date, the study of the gas blanket approach has been limited to edge modification by hydrogen gas injection. By proper programming of injection, high-Z impurities have been reduced in ORMAK (*Ref. 4-55*) and PLT (*Ref. 4-56*). The reactor analog of this is some method of edge cooling by gas injection or by maintaining a plasma edge that is intentionally impurity contaminated. The advantage of this is that the energy flux from the plasma could be spread almost uniformly over the walls as radiation while wall erosion would be lowered due to the low temperature edge.

In principle, this cold, contaminated edge might be achieved by impurity flow reversal based on hydrogen injection or limiter recycling, or it might occur naturally through temperature gradient screening (*Ref. 4-57*). Under some circumstances, the sharp temperature profiles at the edge that will result from edge cooling may act as a barrier to impurity penetration of the plasma core (*Refs. 4-57, 4-58*). While the introduction of a cool boundary layer is favorable to plasma energy balance, it must be kept in mind that the resultant increased peaking of the current profile makes plasma disruptions and plasma shaping increased concerns.

4.2.3.3 Other Edge Modification Techniques — One difficulty in maintaining a layer of "cold" plasma to thermally insulate the hot central core from the wall lies in the high thermal conductivity required to produce a high density, low temperature region that is broad enough to permit dissipation of the large power emerging from the core. If the plasma conductivity could be increased at will, then the temperature profile could be appropriately adjusted in order to handle the requisite heat load. One possible way to make a low-temperature plasma turbulent is by driving an electric current along the direction of the magnetic field. Under such conditions, a "current-convective" instability can be excited (*Ref. 4-59*). This instability is due to the coupling between a steady-state current (parallel to the main magnetic field) and the transverse variation in resistivity which may arise, e.g. from transverse temperature gradients. Preliminary theoretical investigations have suggested that the use of an alternating electric field of a frequency chosen to obtain a skin depth of the order of the desired thickness of the layer may in fact be a viable impurity control alternative (*Ref. 4-60*).

In a similar vein, it has been suggested that the propagation of rf waves in the plasma edge may enhance the diffusion coefficient by up to an order of magnitude over the empirical value (*Ref. 4-61*). Such transport coefficient modifications would serve to increase the impurity recycling rate, and the subsequent enhancement in the radiative losses would ameliorate the technological burdens involved in stimulating the turbulence.

Another proposed edge cooling scheme involves modifying the poloidal magnetic field configuration such that the volume of the low temperature boundary is substantially larger than in the conventional approach (*Ref. 4-62*). An example of such a modified plasma configuration in which the top outermost flux surfaces are expanded is depicted in Fig. 4-5. In such an approach, most of the heat flows across the magnetic surface near the midplane, where the temperature gradients are steepest, and into the top expanded volume, where it is dissipated by radiative cooling. For a sufficiently low edge temperature or a sufficiently large boundary, a significant temperature variation may occur along the outermost flux surfaces. Specifically, this results when

the heat flow as a result of the temperature gradient along the field line can be dispersed by radiation. Such an effect, which should be operative in large tokamaks, would further enhance the operating characteristics since the radiative losses increase substantially as the temperature is lowered at fixed plasma pressure.

4.2.4 Divertors

By judicious choice of field-shaping coils, some magnetic field lines that pass through the plasma region may be diverted, i.e., extended to regions far from the discharge. Since particles following such diverted field lines do not impact matter until they are far removed from the main body of the plasma, impurity generation will be reduced; in addition, such a system can shield the plasma from wall generated impurities because the latter, upon ionization, may be diverted before entering the plasma core. These two functions

are termed unloading (or sweeping) and screening (or shielding), respectively. One significant advantage over other approaches to impurity control is that the sweeping action reduces recycling, a particularly important point in coping with the accumulation of helium, the by-product of the D-T reaction. Thus, on physics grounds, divertors appear to be the logical route to the achievement of long pulses, particularly in a burning plasma. Here, as is conventional, the term "divertor" refers both to the volume occupied by the diverted flux and to the associated hardware.

Various types of magnetic field divertors have been devised for toroidal devices. All are based upon the idea of generating a null at some point (or along a line) in a component of the magnetic field. This generates a separatrix and field lines on the outside of this separatrix are carried away from the plasma surface as they pass the neutral point (null point). Magnetic flux surfaces on the inside remain closed within the plasma volume. The separatrix, sometimes called a magnetic limiter, then becomes the boundary of the toroidally confined plasma.

The types of divertors that have been incorporated into existing or planned toroidal devices are displayed in Fig. 4-6. The toroidal divertor was the first employed. It has been used on stellarators (*Ref. 4-64*) but is inappropriate for tokamaks because it generates an unacceptably large deviation from axisymmetry, typically of order 50% on axis. A variation on the toroidal divertor, the so-called bundle divertor, was proposed by a Culham group (*Ref. 4-65*). This approach, utilizing two opposing current loops to divert a bundle of magnetic flux, has been tested on DITE (*Ref. 4-66*) and will also be implemented on both ISX-B and TEXTOR. A null in the poloidal field can be achieved by inclusion of poloidal coils which carry current in the same direction as the plasma current. Such poloidal divertors, which have the virtue of preserving axisymmetry, can assume a variety of configurations, two of which are shown in Fig. 4-6. DIVA has achieved encouraging results with a single-null poloidal divertor (*Ref. 4-67*) and major tests of the double-null approach are planned for PDX and ASDEX. A single-null poloidal divertor with the separatrix located on the large major radius side of the plasma is an important feature of JT-60, a large tokamak now under construction.

Fig. 4-5 Schematic of an expanded radiative edge cooling region.

Experiments with divertors have been successful in improving plasma purity. The toroidal divertor in the Model C stellarator reduced the oxygen concentration by two orders of magnitude. The energy balance in DITE is radiation dominated without divertor action while the radiated power is only a small fraction (10–15%) of the ohmic power when the divertor is activated. DIVA has reported that the divertor has resulted in a factor of 3 reduction in radiation. These experiments have documented the unloading, screening, and magnetic limiter properties without observation of deleterious effects on plasma confinement. Tests at more reactor relevant temperatures and energy densities will be forthcoming in the next generation of experiments.

Of equal importance to the underlying physics are the engineering and technological implications on device design of the choice of impurity control method. Common to all divertors are the problems of handling the particle and heat loads. The unload function means that greater pumping capacity will be required and, in the case of a D–T device, an increase in the tritium inventory is necessitated. The heat load problem mandates a large surface area for the collector. These two features alone guarantee that a divertor will significantly alter the overall design of the device. In fact, as seen below, the design constraints in a reactor context are so severe that no acceptable divertor design has been generated to date.

The bundle divertor has the advantage of removing the particle and heat collection functions external to the toroidal field coils. They may also be replaced without major disassembly. However, achieving the needed area for heat collection in a reactor grade plasma will require a complex field-shaping coil system to expand the diverted flux tube. Furthermore, complex coil designs generating higher order multipole fields will be needed to achieve acceptable ripple in the center of high temperature plasmas, and this problem increases in severity in the reactor case because of shielding requirements. The forces on the divertor coils were at one time regarded as so serious a problem as to preclude high field application, but recent advances in design (*Ref. 4-68*) have reduced these concerns.

Poloidal divertors preserve axisymmetry and naturally provide a large collector area, but they have the drawback of requiring high current poloidal coils near the plasma. This introduces two problems in a reactor application. One is to provide space to adequately shield the internal poloidal coils from neutrons and the second is to develop a coil system that allows remote maintenance of the tokamak. While the first problem is clearly solvable, albeit at the cost of increasing the size of the device, the second problem remains as a fundamental question. One possible solution is to situate the divertor coils outside the toroidal field coils but the force problems are severe in this instance.

Fig. 4-6 Types of divertors that have been incorporated into toroidal devices (from *Ref. 4-63*).

4.3 EXTENDED OPERATION

In the conventional mode of operation, the maintenance of the toroidal current against dissipation by Coulomb collisions necessitates a continuously increasing flux through the centerpost region. Thus, unless an alternative means of driving the plasma current is employed, the tokamak is a pulsed device. In addition to this intrinsic limitation, the practical restriction to date to pulse lengths of less than a second has precluded the exploration of a number of reactor-relevant physical processes. Since high duty cycle operation is a necessity for economic competitiveness of the tokamak as a reactor candidate, long pulse operation must eventually be demonstrated.

The analysis of what constitutes a long pulse is the subject of Section 4.3.1. The succeeding sections deal with two issues pertinent to long pulse operation. The first, fueling, is needed in present generation tokamaks to achieve high density and is essential to replace spent fuel in long pulse reactor operation. The second topic discussed is the possibility of driving the plasma current other than via the ohmic heating coil.

4.3.1 Long Pulse Operation

As a guide to answering the question of what constitutes an adequate discharge duration to produce a true equilibrium plasma condition, a list of the important physical processes in hydrogen operation and their characteristic time scales is given in Table 4-1. The dependence of these time scales on the plasma and machine parameters is illustrated together with order of magnitude estimates for these quantities for "typical" PLT discharges with several hundred kilowatts of auxiliary heating power. Note, for reference, that such discharges in PLT have a pulse length (τ_{pulse}) of approximately one second, with typical neutral beam pulses of 100 msec.

The first three entries in Table 4-1 refer to transport processes. Since the discharge duration (about 1 sec) exceeds these characteristic times by an order of magnitude or more, it is clear that existing ohmically-heated tokamaks are steady-state with respect to transport processes. This is further evidenced by the fact that the plasma density and temperature profiles are observed to be time-independent during the flattop portion of the discharge. In addition, existing tokamaks can in principle be steady state with respect to the heating processes taking place in the plasma, the relevant time scale for beam heating being the fast ion slowdown time (i.e., 0.01–0.1 sec). Unfortunately, the experimental situation on this point is not always definitive because of technological factors that have limited the duration of the heating pulse.

4.3.1.1 Impurity Time Scales — The situation with regard to impurities, which is partly determined by the coronal equilibrium, impurity diffusion, and plasma-wall equilibration time scales

TABLE 4-1
SCALING AND CHARACTERISTIC TIMES FOR A TYPICAL
PLT DISCHARGE FOR IMPORTANT PHYSICAL PROCESSES
IN NON-BURNING PLASMA

Process	Scaling	Time in PLT (sec)
Collisions (τ_{ei})	$T^{3/2}/n$	$10^{-5} - 10^{-4}$
Electron-ion equilibration	$(m_i/m_e) \tau_{ei}$	$10^{-2} - 10^{-1}$
Particle and energy confinement	na^2	$10^{-2} - 10^{-1}$
Fast ion slowdown	$(m_i/m_e) \tau_{ei}$	$10^{-2} - 10^{-1}$
Coronal equilibrium	$1/n$ (?)	10^{-2}
Impurity diffusion	$(T^{1/2}/n) a^2$ (?)	$10^{-2} - 10^{-1}$
Plasma wall equilibration	(?)	$\gtrsim \tau_{\text{pulse}}$ (?)
Resistive diffusion	$a^2 T^{3/2}$	$\gtrsim \tau_{\text{pulse}}$

in Table 4-1, is considerably more complex. Provided the density is not too low and impurity transport is not too rapid, plasmas certainly achieve a steady state with respect to impurity atomic physics, i.e., the impurities are in a state of coronal equilibrium. Impurity generation and transport are not well understood; but, again, the observed asymptotic density and temperature profiles suggest that a steady state has been reached, at least for ohmically-heated discharges, in which impurities play an important role in the overall particle and/or energy balance. Caution must be exercised in drawing such a conclusion, however, for two reasons:

- First, Table 4-1 indicates that the plasma is probably not in equilibrium with the surrounding medium, implying that the impurity generation processes may still be slowly changing, even in existing, low power density devices;
- Second, a large power load on the first wall of long duration, such as may be achieved by long pulse auxiliary heating or significant fusion energy production, may alter the plasma-wall interaction sufficiently to change the character of the impurity problem. (Wall temperatures of 300–600°C with loadings of 20–40 W/cm² are expected for reactor operation.)

The plasma-wall equilibration time is unknown at present, and there is a complex, interlocking set of processes which determine it. Some of these are:

- **Gas Trapped and Re-Emitted by the Wall.** Fuel atoms diffuse into the metal, reach a steady state, and may be re-emitted later in the pulse.
- **Saturation of Impurity Control Mechanisms.** The gettering system in the flow reversal scheme, or the divertor collection plates in bundle or poloidal divertors, may reach saturation or thermal limits in a matter of seconds.
- **Chemical Stabilization of the Surface Layer.** Desorption of hydrogen monolayers, or the equilibration of the impurity constituents on the surface (C, O, Ti) and the migration of limiter-produced impurities through recycling require many particle confinement times. This has a strong effect on physiochemical sputtering.

To obtain a crude estimate of the plasma-wall equilibration characteristic time scale τ_{pw} , assume that a steady state is obtained when the discharge has been sustained long enough to desorb a monolayer of absorbed gas from the walls. Then

$$\tau_{pw} = \tau_p \frac{N_w}{N_p} \approx 4a \quad (4-4)$$

where τ_p is the particle confinement time, N_w is the number of particles in a monolayer, and N_p is the number of particles in the plasma. The numerical estimate was obtained by assuming a $5 \times 10^{19} \text{ m}^{-2}$ coverage of the walls and by taking $\tau_p = 4\tau_E$ where τ_E is the global "Alcator" energy confinement time. This estimate amounts to 2 sec for PLT and Doublet III but, in reality, τ_{pw} for this process may well be considerably smaller; depending on the method of wall preparation employed, the wall coverage has probably been over-estimated and the relevant τ_p for recycling is probably considerably shorter than the global particle confinement time. If such a model is valid, present day plasmas are essentially in equilibrium with the walls. The remaining concern, however, is the extrapolation of our experience to date to the much higher wall loadings and wall temperatures characteristic of reactor operation.

4.3.1.2 MHD Time Scales — The resistive diffusion time, the final entry in Table 4-1, is the process with the longest time scale among those listed. The time scale of interest is actually that in which the flux configuration changes sufficiently to modify the equilibrium distribution of magnetic fields within the plasma. We adopt the skin time as the basic time scale. Assuming for simplicity a uniform density and temperature plasma, it follows that

$$\tau_{\text{skin}} = \frac{15}{Z_{\text{eff}}} \frac{2\kappa^2}{1 + \kappa^2} a^2 T^{3/2} \quad (4-5)$$

where κ is the vertical elongation, a is the minor radius in meters, and T is the temperature in keV.

The key concern with this process of current relaxation is the question of MHD stability. Although dynamical paths to high beta equilibria have been demonstrated via MHD studies that simulate intense auxiliary heating, there is the concern that these "flux conserving equilibria" will relax to MHD unstable states. To examine this question experimentally and, indeed, to study the entire issue of the control of high beta plasmas requires that the pulse length exceed some fraction of the plasma skin time.

The current profile relaxation also strongly affects the probability of a disruptive instability. As gas puffing, recycling, beam or pellet fueling proceed, the contraction of the current channel can lead to disruption, a phenomenon that has been observed at the highest currents in present tokamaks. Since the highest possible current will be chosen to optimize fast-ion and alpha confinement, there may be very little margin of safety in this area. Slow-growing instabilities, the time constants scaling as fractional powers of resistivity (e.g. resistive ballooning), may also play a role as pulse lengths are extended.

Other time constants may also be relevant, depending on the chosen mode of operation. Gas puffing as a means of fueling tokamaks has been observed to modify density profiles on a time scale which is considerably faster than the confinement time. While this is not fully understood, it is anticipated that the characteristic time scale for the physical processes that govern gas puffing will not exceed the particle confinement time τ_p . Pellet injection is somewhat more complicated because more variables enter the problem, but the same conclusion holds in cases of interest; in this case, of course, "steady-state" is really a misnomer, and, in fact, refers to a time-averaging over the pellet injection period.

To summarize the situation for hydrogen operation, a pulse length of 1 second is found to be adequate for quasi-steady-state operation of an ohmically heated plasma while on the order of 10 seconds may be required for studying plasma-wall interactions/impurity effects and current profile relaxation in an auxiliary heated plasma. For reactor relevance, current profile relaxation should be studied in a high beta context; hence, it is most profitably carried out at modest toroidal field in a vertically elongated plasma.

4.3.1.3 Time Scales for Burning Plasmas — New time scales enter for the case of a burning plasma. The charged particle fusion products, 3.5 MeV alpha particles in the D-T case, transfer their energy to the plasma over a period of 0.1–1 second, the exact value being most sensitive to the plasma temperature. Because this time scale is short and because, by definition, alpha heating dominates the energetics of an ignited plasma, it may be that the previously identified "long pulse" questions, i.e., plasma-wall equilibration, current profile relaxation, and the possibility of long time scale impurity phenomena, may be quite different in a burning plasma from those in a hydrogen plasma. In fact, the differences in plasma pressure and current profiles may give rise to quite different transport processes for burning plasmas. Additional phenomena unique to burning plasmas are associated with the process of fuel burning and the concomitant problem of ash removal.

Typical reactor fuel depletion rates are on the order of 1%/sec so a pulse length of tens of seconds is needed to study these effects. Control of fuel density profiles is another non-trivial aspect of D-T operation with long pulses. The diffusion properties of thermal alpha particles will be a major determinant of the useful burntime; however, this is amenable to experimental investigation via injection of helium into a hydrogen plasma. It is worth noting that the combination of large size and high temperature associated with reactor operation translates into such a long skin time ($\sim 10^3$ sec) that it is by no means certain that reactors must operate for periods in excess of a skin time in order to be economically attractive. It would be valuable, however, to know whether current relaxation can be accommodated in a reactor-grade plasma in order to assist in the prognosis for possible true steady-state tokamak operation.

Another phenomenon unique to burning plasmas is the thermal excursion that inevitably follows ignition. The mechanisms which will provide the energy losses to govern the burn equilibrium point are not understood. Depending on the underlying physical process, whether it be increased transport, enhanced MHD activity, or some feedback control scheme for an inherently unstable operating point, the associated characteristic time may vary over orders of magnitude and again may alter conclusions drawn in the case of hydrogen operation.

4.3.2 Fueling

Fueling is an essential ingredient of tokamak reactor operation. Without fresh fuel being periodically supplied to the plasma core to replace the fuel consumed by fusion reactions, the burn time would be limited to tens of seconds, a constraint which is probably incompatible with economic realities. Depending on the method of impurity control employed, additional burdens may be placed on the fueling system; for example, both gettering and divertors reduce the recycled fraction. Furthermore, some means of fueling is required in the startup phase since breaking down the gas and stripping out the impurities at the densities needed for ignition would put an intolerable burden on the poloidal field system. Various techniques for refueling, i.e. neutral beam injection, pellet injection, edge fueling, and plasma guns will each be discussed briefly below.

4.3.2.1 Fueling by Neutral Beams — Neutral beams are a proven means of fueling plasmas with modest opacity, defined as the product of the density and the minor radius. In fact, high power, low energy injection raises the problem that the plasma density will build up to a level that will impede beam penetration, a problem which is aggravated by the cold gas streaming from the beam ducts. On the other hand, the beams themselves will be a poor source of particles for reactor grade plasmas because the neutral beam energy required for adequate penetration is so high that a relatively small number of particles is injected. To see this quantitatively, note that the beam energy demanded to simultaneously satisfy particle and energy balance is given by

$$E_b = \frac{\tau_p}{\tau_E} \frac{3\bar{T}}{1 + Qf_\alpha/5} \quad (4-6)$$

where t_α is the fraction of alpha power deposited in the plasma and Q is the ratio of the fusion power to the beam power. For reasonable reactor startup scenarios, this leads to $E_b \lesssim 10 \bar{T}$. Such energies are wholly inadequate to penetrate a plasma with $n_\alpha > 10^{20} \text{ m}^{-2}$, which, according to Alcator scaling, is a prerequisite for a reactor-grade plasma. This conclusion may be invalidated if ripple-assisted injection is employed. However, electric fields resulting from the substantial injection currents needed to provide the bulk of the fueling are a source of concern in this approach. Unless such problems are avoided, it can be concluded in general that neutral beams will not provide an adequate means for refueling large tokamaks.

4.3.2.2 Pellet Injection — One thoroughly studied method for fueling reactor tokamaks is the injection of high velocity pellets of D-T ice. Once in the plasma, the impinging flux of energetic particles rapidly ablates the pellet and the key issue is to develop an injector capable of accelerating pellets of sufficient size to the velocities required for adequate penetration. Fortunately, the ablation rate is retarded through the formation of a protective layer of previously ablated material surrounding the pellet; this reduces the required velocities to difficult but at least technologically conceivable values. As will be discussed below, the physical models developed to describe the ablation process are in relatively good agreement with experiment so the major unanswered questions in this field are the technological ones.

Key experimental evidence regarding pellet injection was gathered on ISX-A (*Ref. 4-69*). Pellets of diameter 0.6 mm were injected at speeds of $3-4 \times 10^2$ m/sec into an ohmically heated discharge. The observed ablation rate followed a similar pattern to what was seen on ORMAK (*Ref. 4-70*) and is in general agreement with theoretical predictions of *Ref. 4-71* which employed a model built upon earlier work (*Refs. 4-72, 4-73*). The key element of these models is the description of the gas shielding effect, i.e. the interaction of the background plasma with the molecular cloud surrounding the pellet.

By varying the plasma density while holding the pellet size fixed, information was gathered regarding the effect of a large localized cold particle source on plasma stability and confinement. No adverse effects were found even for cases in which a single pellet gave rise to a 30% increase in average density. This observation implies that relatively large pellets can be utilized in reactor-grade plasmas. Since larger pellets take longer to ablate, the velocity requirement and hence the technological problems are eased. For a plasma with an average density of 10^{20} m⁻³ and an aspect ratio of 4, the requirement that the pellet give rise to a density perturbation satisfying the inequality $\Delta n/n < 0.3$ translates into the limit

$$\frac{r_p}{a} < 6.9 \times 10^{-3} \quad (4-7)$$

on the pellet radius r_p . Full penetration of such a pellet into a reactor-grade plasma ($a \approx 1$ m) demands a pellet velocity of about 10^4 m/sec. However, if recent speculation that less than complete penetration may be preferable (*Ref. 4-74*) is verified, pellet velocities of $2-3 \times 10^3$ m/sec may prove adequate.

A number of techniques have been proposed to accelerate pellets to velocities in excess of 10^3 m/sec: light gas guns, electrostatic acceleration, electromagnetic acceleration, and a variety of mechanical acceleration approaches, the latter generally employing rotating machinery. The mechanical approaches have been the most thoroughly explored in practice but they encounter strength of material limitations at velocities of 5×10^3 m/sec or less.

Another concern for pellet fueling is the difficulty of retaining the pellet intact and cryogenic during the injection process. This is of particular significance for D-T pellets because β -decay related effects greatly alter the electrical and thermal conductivities and may degrade the mechanical properties as well.

4.3.2.3 Gas Puffing — Certainly the most desirable fueling technique from a technological perspective is simply to admit gas at the vacuum-wall boundary. This technique has proven successful in many tokamaks in which rapid density increases have been produced and the resulting discharges have been maintained for many particle recycling times. The most striking aspect of these experiments is the increase of the entire plasma density profile on a time scale that is short

compared to the confinement time. This feature, while highly desirable, has been difficult to explain theoretically; in fact, the speculative nature of the suggestions offered to date to explain gas puffing data make extrapolation to the reactor regime quite conjectural. On a more positive note, the observed success of gas puffing has been so universal that an optimistic attitude toward applicability in new regimes is not uncalled for.

A scheme that is often suggested for fueling reactor tokamaks utilizes a gas blanket, a term introduced by Alfvén and Smars (*Ref. 4-75*) to describe a cool, dense layer of plasma that thermally isolates the plasma from the vessel walls. Such a region is generally envisioned as satisfying $n_i \ell > 10^{19} \text{ m}^{-2}$, where n_i is the ion density and ℓ is the thickness of the layer. A fluid description of this layer and the surrounding layer of neutrals has been presented by Lehnert (*Ref. 4-76*). Such a gas blanket should be useful for fueling large, dense plasmas, particularly if an inverted density gradient can be sustained. However, a gas blanket has not yet been observed, apparently because of the large appetite of the plasma for gas at the edge. The gas load problem is more serious for the gas blanket approach but it is a potentially viable mode of operation, one that may take the place of conventional gas puffing in reactor operation.

There are two advantages of gas puffing that supplement its experimental successes: first, no technological advances are needed for its implementation and second, the apparatus required involves negligible capital cost and energy investment. One drawback is that raising the plasma density in the core by means of edge fueling may be incompatible with divertor operation.

4.3.2.4 Other Fueling Techniques — The intriguing possibility exists that plasma guns of the Marshall or deflagration types may provide a mechanism for refueling tokamaks. Such guns produce a high particle flux at a high efficiency and there is evidence that careful operation can produce fuel with low impurity contamination. However, theoretical uncertainty exists concerning the ability of the plasma so produced to penetrate across the magnetic field into a tokamak plasma. Penetration across magnetic field lines has been observed (*Ref. 4-77*) but the mechanism responsible, believed to be related to the formation of a transverse electric field which induces an $E \times B$ drift motion, is poorly understood, particularly when injecting into an existing plasma.

The use of cluster beams is another possible fueling mechanism. However, this technique suffers from the same penetration difficulty as neutral beams so the outlook for fueling a large reactor plasma is highly unfavorable, unless ripple injection or some other scheme is employed to reduce the required beam energy.

4.3.3 Steady-State Operation: Current Drive Techniques

Present tokamaks depend on inductive processes for the maintenance of the plasma current. Because the flux-swing capability of ohmic heating coils is limited, such devices must operate in a pulsed mode. Continuous operation of tokamaks may become possible if alternative steady-state mechanisms for providing the plasma current drive could be demonstrated.

Continuous operation of tokamaks would provide a number of important engineering simplifications that might enhance their acceptability by utility companies. Naturally, steady state operation eliminates the need for thermal storage to prevent from pulsing the power output to the electrical grid. Continuous operation would also produce steady power output without cyclic thermal excursions of any portion of the reactor; thermal cycling of the first wall is a particularly serious problem that must be avoided if the structural integrity of the wall is to be maintained over long periods. Non-inductive steady-state operation greatly reduces the engineering required for the ohmic heating coil and its related switching gear; however, ohmic drive may be required in a reactor

during the startup phase, even if current drive is successful. Further, pulsed magnetic fields will be vastly reduced; this eases the design constraints on the use of superconductors, which are sensitive to time-dependent fields. Finally, the large losses that occur in the plasma during startup could be nearly eliminated; this both reduces the capital cost of the facility and increases the useful plant power output.

In general, techniques for maintaining a plasma current in tokamaks rely on the preferential acceleration of electrons relative to ions. Collisions tend to reduce differences in the mean relative motion of electrons and ions. Hence, power must be supplied to the plasma so that the resistance resulting from electron-ion collisions can be overcome and the plasma current can be maintained. The current drive techniques explored in subsequent sections are unidirectional neutral beam injection and the injection of traveling rf waves.

Before proceeding to this discussion, it is worth noting that most of the advantages to be gained from true steady state operation are present to a degree when sufficiently long pulse lengths are employed. The most straightforward way to increase the allowed pulse length is to increase the bore of the ohmic heating coil to gain additional volt-second capability. Because the one-turn voltage for a reactor grade plasma is much less than one volt, a considerable increase in pulse length can be obtained by a modest increase in aspect ratio. It is to be kept in mind, however, that present theory suggests that the limiting value of beta is inversely proportional to the aspect ratio so significant increases in aspect ratio are to be avoided if possible.

Another attempt to deal with the apparent restriction of tokamak devices to a pulsed mode is to design a field-shaping coil set which is capable of simultaneously creating one plasma while burning another. In one such approach, an extension of an earlier suggestion by Yoshikawa (*Ref. 4-77*), the plasma discharge in a burn chamber is moved to a burial chamber for quenching and is replaced by another plasma, all on a time scale which is short compared to the thermal time constant of the surrounding media (*Ref. 4-78*). In another approach, the forces between a newly formed plasma and an existing plasma lead to a merging of the two in a burn chamber (*Refs. 4-79, 4-80*). By programming the field-shaping coils correctly, the new plasma winds up in the center of this merged plasma so the process potentially deals with the refueling and impurity control problems as well. To date, neither of these approaches has been examined experimentally.

Yet another possibility for steady-state operation is the utilization of the bootstrap current, a current which has been predicted on the basis of neoclassical theory, but which has not been observed experimentally. The bootstrap current owes its existence to the net flux of charged particles; such a constraint is incompatible with edge fueling. The great promise of this approach is that the magnitude of the current for reactor grade plasmas is sufficient to eliminate any need for volt-seconds from the ohmic heating coil during the burn phase. Furthermore, the current profile associated with the bootstrap current is outwardly peaked, a very desirable profile from the standpoint of plasma shaping.

4.3.3.1 Neutral Beam-Driven Currents — When an ion beam is injected parallel to a uniform magnetic field in a pure hydrogen plasma (e.g., by neutral injection), the ion current is exactly compensated by an electron shielding current. Essentially, the frictional force between the beam ions and the background electrons will tend to cause electron streaming, and this electron streaming will be opposed by the frictional force between the electrons and the background ions. When the frictional forces on the electrons are in balance, the electron current exactly cancels the beam current.

If the background plasma contains impurities, the frictional force opposing the electron streaming will be increased, with the result that the beam current is not completely compensated

by the electron current (*Ref. 4-81*). Similarly, in a toroidal geometry, the trapped electrons add to the frictional force on the streaming electrons, again leading to a lack of complete compensation of the beam current. A simple formula for the net current per unit area may be given that includes both of these effects (*Ref. 4-82*):

$$j_{\text{net}} = j_{\text{beam}} \left[1 - \frac{1}{Z_{\text{eff}}} + \frac{1.5}{Z_{\text{eff}}} \left(1 + \frac{0.7}{Z_{\text{eff}}} \right) \left(\frac{r}{R} \right)^{1/2} \right] \quad (4-8)$$

$$j_{\text{beam}} = 2\pi R e \int v_{\parallel} f_{\text{beam}} d^3v$$

In present-day tokamaks, with high injection power and low-to-moderate plasma density, this net beam-driven current should attain experimentally observable values.

Neutral-beam injection appears to be an undesirable method of maintaining a steady-state current for ignited or near-ignited operation. The principal reason is that minimizing the injected power while injecting large numbers of fast ions to carry the current demands relatively small beam energy, so that a low opacity (na) must be used to permit neutral-beam penetration. But low opacity is generally associated with low $n\tau_E$, which implies a small Q and hence an unattractive energy balance. However, the neutral-beam-driven current, together with the equilibrium-field system, could possibly be used to start up the plasma current without the assistance of a central ohmic-heating transformer (*Ref. 4-83*). This current would decay on either the "skin time" scale, or the L/R time, both of which can be several hundred seconds for a plasma of ignition size and temperature. By this technique, an extremely compact ignited tokamak reactor might be made feasible, although operation would of necessity be pulsed.

4.3.3.2 RF Driven Currents — The time averaged momentum ($\vec{p}$) and energy ($\dot{\mathcal{E}}$) density associated with traveling plasma waves are related by

$$\vec{p} = (\vec{k}/\omega) \dot{\mathcal{E}} \quad (4-9)$$

At steady state, deposition of toroidal rf momentum on the electrons and subsequent dissipation of the electron momentum at a rate ν results in an average charge carrier velocity v given by

$$\dot{p} = nm\nu v \quad (4-10)$$

where n is the number density of current carrying electrons. Accordingly, the ratio of current to required rf power is

$$j/\dot{\mathcal{E}} = \frac{e}{m\nu v_{ph}} \quad (4-11)$$

where v_{ph} is the toroidal rf phase velocity. Experimental studies of rf current generation (*Refs. 4-84, 4-85*) have been performed in collisional plasmas wherein the traveling plasma waves are absorbed via non-resonant collisional damping; agreement with theory is obtained (*Ref. 4-85*) for ν constant and proportional to ν_e defined in Section 2.2, Eq. (2-7).

Recent approaches to rf current drive have sought to maximize the current to power ratio. If the rf momentum is deposited on the electrons at $v \sim v_{ph}$, and is dissipated via electron-ion collisions, then

$$j/\dot{\mathcal{E}} = \frac{e}{m\nu_e v_{th,e}} \cdot \begin{cases} C_1 (v_{ph}/v_{th,e})^2 & v_{ph} > v_{th,e} \\ C_2 (v_{th,e}/v_{ph}) & v_{ph} < v_{th,e} \end{cases} \quad (4-12)$$

Experimental evidence for collisionless absorption of slow Alfvén waves at $v_{ph} \sim v_{th,e}$ and associated rf current generation has been obtained in the C-Stellarator (Ref. 4-86). The high phase velocity approach to maximizing $j/\dot{\epsilon}$, utilizing the fact that the collision frequency is proportional to v^{-3} for tail electrons, has been examined for lower hybrid waves and fast Alfvén waves (Refs. 4-87, 4-88). The numerical constant C_1 must be determined by detailed Fokker-Planck code analysis (Refs. 4-89, 4-90), a linearized version of which yields a value equal to 0.23 (Ref. 4-90). A practical realization of the low phase velocity approach with plasma waves utilizes Alfvén magnetosonic cavity modes (Ref. 4-91) absorbed by transit-time damping and Landau damping. Two-dimensional Fokker-Planck computations give $C_2 \approx 1.0$ (Ref. 4-92).

To examine the feasibility of maintaining equilibrium in a reactor entirely by rf induced currents, the required total rf power, P_{RF} , must be compared to the fusion power, P_F . If the requisite current is determined by the condition $\beta_p = R/a$, then

$$\epsilon \equiv \frac{P_{RF}}{P_F} = \begin{cases} \frac{2.5 \times 10^{10}}{\bar{n}^{1/2} \bar{T}^{3/2} a^{1/2} R^{1/2} (v_{ph}/c)^2} & v_{ph} > v_{th,e} \\ \frac{2.8 \times 10^{12} v_{ph}/v_{th,e}}{\bar{n}^{1/2} \bar{T}^{5/2} a^{1/2} \bar{T}^{1/2}} & v_{ph} < v_{th,e} \end{cases} \quad (4-13)$$

where parabolic density and temperature profiles have been assumed and the units are MKS except that $\bar{T}$ is in keV. The fusion power has been taken to be proportional to $n^2 T^2$ and, as in Eq. (4-12), $j/\dot{\epsilon}$ is evaluated at the average density and temperature. Both approaches favor large, dense, hot plasmas. For the UWMAK-I design, $\epsilon = 4\%$ for the first case with $v_{ph}/c = 0.5$, and $\epsilon = 2\%$ for the second case with $v_{ph}/v_{th,e} = 0.2$. Optimization of ϵ based on depositing the rf momentum in the low density peripheral region of the plasma has been proposed (Ref. 4-93), but stability of the resulting equilibria has not been examined.

In the low phase velocity approach, limitations of the above theory as a result of low $v_{ph}/v_{th,e}$ have not been fully investigated. In particular, the interaction with toroidally trapped particles may decrease the efficiency of this approach (Ref. 4-94).

In addition to transferring rf momentum directly to electrons, it has been suggested (Ref. 4-81) that the neoclassical bootstrap current may be enhanced by injection of rf waves which are absorbed by the toroidally trapped particles.

In summary, the rf current drive concept presents several promising approaches to attain steady-state toroidal current. However, much theoretical and experimental effort is still required before the utility of the concept for reactors can be thoroughly evaluated.

4.4 DIAGNOSTICS

The importance of plasma diagnostics to progress in tokamak research cannot be overestimated. Diagnostic systems are the eyes and ears of the experimental physicist; they facilitate quantitative comparisons of experiment and theory and permit systematic monitoring and adjustment of discharge parameters. Fortunately, tokamak diagnostics have evolved considerably since the time of the earliest experiments. The trend has been away from diagnostics which give gross volume-averaged information, towards systems capable of high radial resolution. The next step has been to measure entire radial profiles in a single shot, and finally some of the newest systems give radial profile information continuously as a function of time. For non-circular cross-section tokamaks, it may eventually prove necessary to follow the time evolution of some parameters as functions of two spatial variables.

In contrast to the dramatic progress in providing multi-dimensional information on some important plasma properties, relatively little advance has been made towards developing diagnostics for certain other quantities — most notably $Z_{\text{eff}}(r)$ and $q(r)$. While techniques have been demonstrated for measuring both of these parameters, none has as yet been used as a "standard" diagnostic for tokamak research. Some new plasma properties have come under diagnostic scrutiny recently, however. As prime examples, bolometric measurements of total radiated power have become standard tools in the study of electron energy confinement, and new high-resolution X-ray crystal spectrometers have become available to measure impurity temperatures and charge-state distributions.

In this overview, no attempt will be made to discuss in detail all of the diagnostic systems used on present-day tokamaks or planned for future devices. In any event a recent comprehensive review of the subject by the TFR Group (*Ref. 4-95*) is available. The primary focus of this treatment concerns what quantities can and cannot be measured in current tokamaks, touching briefly on the implications that this has for our understanding of tokamak physics. A secondary theme revolves around diagnostic development through discussion of the technical successes and the limitations of the major systems currently in use, with perhaps some extra emphasis on recent progress since the TFR review paper. Where appropriate, mention is made of the problems associated with plasma diagnostics in the hostile environment of future reactor-grade tokamak experiments.

4.4.1 Magnetic Measurements

In order to assess the total ohmic input power to the plasma column (IV , or equivalently the Poynting flux $E_{\phi} B_{\theta}$), it is necessary to measure the total plasma current, and the toroidal electric field at the plasma surface. The plasma current is measured using a Rogowski belt around the plasma minor circumference, which gives a voltage proportional to $\int \dot{B}_{\theta} r d\theta$. Using an active operational amplifier integrator, one then obtains $I_p(t)$. Any current flowing toroidally in a vacuum vessel without an insulating break must be estimated and subtracted. An evenly wound Rogowski belt evaluates the line integral of $\vec{B}$, and so gives the true total enclosed current, being insensitive to local eddy currents and small misalignments. The situation with the toroidal electric field measurement, however, is somewhat more difficult. In devices with insulating breaks in the vacuum vessel, one measures the voltage across this break, while in tokamaks with resistive vessels, a separate loop is wound close to the discharge chamber, and the voltage around this loop gives the "loop voltage", V . In neither case, however, is the measuring device located at the plasma surface, so changes in the plasma current and position give rise to changes in the magnetic flux linked between the plasma and the measuring loop. These flux changes can be estimated, however, and a "surface voltage" can be deduced with reasonable accuracy.

Two major difficulties still remain in determining the ohmic input power to the plasma, however. On the one hand, the increasing use of SCR power supplies, rather than capacitor banks, to power the ohmic heating transformers in long-pulse devices gives rise to a large amount of switching noise in the loop voltage. The massive analog electronic and/or digital smoothing required to find the true signal introduces either phase lag or spurious precursor effects into analyzing transient phenomena. This makes determining the loop voltage during relatively short-pulse auxiliary heating experiments, for example, extremely difficult. A second more fundamental problem arises from the fact that one can only measure the total Poynting flux into the plasma. Some of this power flow may be going into the buildup of poloidal field energy, or alternatively poloidal field energy may be being converted into resistive heating within the plasma core. Thus, if there is reason to believe that the poloidal flux is changing significantly, the resistive component of the measured IV becomes difficult to determine. Typically, experimenters subtract $L_i dI_p/dt$ from the loop voltage, in an attempt to estimate solely the resistive component. The internal inductance, L_i , is generally determined by assuming that the plasma resistivity is neoclassical, with Z_{eff} independent of r . This procedure allows some correction for changes in the plasma current, with the result that the input resistive power can be determined with perhaps $\pm 15\%$ accuracy during the quasi-steady-state phase of a discharge. During startup, shutdown, and transient changes, however, the uncertainty in $dL_i I_p/dt$, combined with the SCR noise, can make the uncertainties as large as a factor of two.

A second frequent use of magnetic loops is to measure the position of the plasma column. This has become increasingly important in modern tokamak devices without copper shells, which use active feedback control of the externally imposed vertical field. Various approaches are used in the design of the magnetic probes for plasma position. The classical approach is to use four individual $\hat{B}_\theta$ loops located top, bottom, inside and outside. Another possibility is to use two continuous Rogowski belts, one with a winding density proportional to $\cos \theta$, and the other proportional to $\sin \theta$. A third technique is to locate "saddle coils" around the plasma column, such that a properly centered discharge has no poloidal flux linking any of the saddles. With any of these loops, a further interesting question arises as to whether one should measure the fields both from the plasma and from the poloidal field system, or whether the signals from the various machine field coils should be nulled out using summing amplifiers. In one approach, the philosophy is to measure the total field, and thus the "true" flux surfaces. In the second approach, the thought is that the near fields from individual poloidal field coils distort the total field locally, and so one does better to measure only the field from the plasma. Which technique works better depends largely on the particular machine design. Both have been used successfully, however, as have the three different types of probe geometry. In general, absolute measurements of plasma position can be made to ± 1 cm or less, during the quiescent steady-state phase of a discharge, and relative measurements of ± 2 mm are typical. During startup and shutdown, however, eddy currents in the vacuum vessel and field coils cause substantial pickup, and plasma position is generally known much less accurately.

More general magnetic techniques for measuring the shape of non-circular plasmas have been developed at General Atomic Company as part of the Doublet experiments there. The most accurate method involves measuring the poloidal flux and magnetic field at a number of locations around the discharge. The flux values are used as input boundary conditions for a free-boundary MHD equilibrium code which solves for the corresponding plasma equilibrium in the system of poloidal shaping coils. In this manner, information can be gained about the plasma current profile by demanding that the poloidal field computed by the code agree with that measured in the experiment. Typically, the plasma shape can be determined to ± 2 cm in Doublet III.

The same sort of local $\dot{B}_\theta$ probes which are classically used for plasma position measurements can also be used to measure fluctuations in the poloidal magnetic field, known as Mirnov oscillations. These fluctuations typically have poloidal mode numbers, m , in the range of $m=1$ to 5, and toroidal mode number, n , equal to one. The frequency of these modes is usually in the high audio range (5–20 kHz), and they are observed to rotate in the electron diamagnetic direction, at a speed comparable to the electron diamagnetic drift velocity. Correlations of the Mirnov oscillations, which are measured externally, with fluctuations in the soft X-ray flux emitted from the plasma core, have permitted increasingly detailed comparisons with resistive MHD theory. While the key effects have not yielded to detailed theoretical understanding, models based on fluid analysis nonetheless do explain many aspects of the measurements and this is one of the most rapidly growing fields of study in tokamak physics (*Ref. 4-96*).

Recently, the TFR group has also reported (*Ref. 4-97*) observing very high-frequency oscillations (up to 300 kHz) on $\dot{B}_\theta$ loops. Since this frequency range is characteristic of drift waves, the measurement of a fluctuating magnetic field at these frequencies strongly suggests the presence of electromagnetic (or finite β) drift waves in the plasma. Similar results have been obtained in the Macrotor tokamak, by actually inserting $\dot{B}_\theta$ loops into the relatively low-temperature and density plasma column. Recent theoretical studies (*Refs. 4-96, 4-98*) have suggested that the magnetic braiding resulting from high frequency magnetic fluctuations could be responsible for the observed anomalous electron thermal transport.

4.4.2 Interferometry

The principal application of microwave techniques to tokamak plasma diagnostics has been for measurement of the line integral electron density. The phase shift of a microwave beam polarized parallel to $\vec{B}$ (ordinary mode) passing through a plasma column is given by

$$\Delta\phi = \int \left(1 - \frac{\omega_p^2}{\omega^2}\right)^{1/2} \frac{2\pi}{\lambda_0} dl \quad (4-14)$$

where $\omega_p^2 = 4\pi n_e e^2/m_e$ and ω is the microwave signal frequency. Thus for $\omega \gg \omega_p$, the phase shift experienced by the wave compared to that of a wave traveling an equal distance in vacuum is approximately proportional to the line integral of n_e . In order to preserve the linearity of this line integral, and to avoid refractive effects, 2 mm microwave systems are typically used for line average densities of $5 \times 10^{13}/\text{cm}^3$ and below, while 1 mm systems (which are difficult to obtain) can be used up to $\bar{n}_e = 2 \times 10^{14}/\text{cm}^3$. Beyond this density range HCN, alcohol, and CO_2 lasers are used.

The principal advances in microwave interferometry in the last few years have been in the use of "fringe-counting" electronics and in the development of multi-channel systems. Until recently, video-type zebra-stripe displays were used to generate interferometric fringes on an oscilloscope CRT, with the result that a certain amount of care and experience was needed to obtain accurate fringe counts, particularly when the plasma density changed rapidly, such as at startup, or during a disruption. Fast electronic circuits now are used to keep track of the phase shift, and while they occasionally also fail in the presence of fast changes, or strong plasma noise, they considerably improve the accuracy of measurement. In addition, the direct readout phase-shift can be straightforwardly digitized and stored. The multi-channel systems, coupled with crude Abel inversion techniques, also provide a measurement of the electron density profile as a function of time.

The second important application of microwave techniques has been to provide measurements of density fluctuations within the plasma column (*Ref. 4-99*). In these experiments, typically a 2 mm microwave beam is directed into the plasma, and a very small fraction of the incident power is scattered into a receiving horn set at a location and angle where it does not directly detect the incident beam. The scattered wave is shifted in frequency and wave vector from the initial wave, and its intensity as a function of $\Delta\vec{k}$ and $\Delta\omega$ reflects the strength of the density perturbation, $\tilde{n}/n$, at frequency $\omega = \Delta\omega$, and wave vector $\vec{k} = \Delta\vec{k}$. These experiments typically detect broadband fluctuations in the region of 0.1–1 MHz with wave-vectors directed perpendicular to $\vec{B}$, and a spectrum of wavelengths centered around $\sim \rho_i/2$. Generally the total fluctuation level, $\tilde{n}/n$, is of the order of 1%. The fluctuations are believed to be due to non-linearly saturated drift modes, possibly driven by trapped-electron effects, but their contribution to transport can only be very roughly estimated on the basis of qualitative arguments. For a direct evaluation of the effect of electrostatic waves on transport one needs the phase angle between $\tilde{n}$ and $\tilde{\phi}$, which cannot be measured by these techniques. Furthermore, the radial localization of the fluctuations is in considerable doubt, since the measurements to date have used rather large scattering volumes within the plasma.

Recently on PLT, with high power neutral beam injection, a new class of density fluctuations has been observed (*Ref. 4-100*). At ion temperatures in excess of 4 keV, the fluctuation amplitude at certain wave-vectors was seen to increase by nearly an order of magnitude, and the frequency spectrum was broadened and broken up into distinct bands separated by ~ 50 kHz. The nature of these new fluctuations, and even their total $\tilde{n}/n$ have not as yet been determined, so speculation as to their importance must await further measurements.

4.4.3 Laser Techniques

Since the joint British-Soviet measurements on T-3 in 1969, Thomson scattering has become the mainstay of modern tokamak plasma diagnostics. In this technique a Q-switched ruby laser is used to shine a ~ 15 nsec pulse of 6943 Å light into the plasma column. A tiny fraction of the incident photons scatter incoherently off electrons in the plasma, and are collected in a multi-channel polychromator located at 90° to the laser beam. The total flux of scattered light is a measure of the plasma density at the location where the laser beam and the polychromator sight-line cross, and the doppler spreading of the scattered photons reflects the local temperature.

Problems associated with these measurements are background plasma light, stray reflected laser light, poor statistics due to the inefficiency of photo-multipliers at these wavelengths, and shot-to-shot variation of plasma conditions, when a radial profile is taken on a point by point basis over many discharges. Adequate solutions to each of these problems have been found, involving fast-gated detectors, beam dumps at a distance from the vacuum chamber, razor-blade type viewing dumps, new types of photocathodes and, recently, multi-spatial-channel systems capable of measuring entire radial profiles in a single discharge. Figure 4-7 shows a diagram of the PLT multi-channel Thomson scattering system (*Ref. 4-101*), based on a micro-channel-plate image intensifier, coupled with a cooled silicon vidicon tube. The spectrometer is capable of providing 56 channels of spatial resolution, and divides the scattered spectrum from each location into 20 wavelength elements. The accuracy of each spatial channel of a system of this sort is somewhat inferior to what can be obtained in a single-channel system, but an uncertainty of $\pm 15\%$ in each individual channel is typical, and by averaging over a number of shots, quite good accuracy can be obtained.

Fig. 4-7 A top and side view showing the components of the PLT multi-channel Thomson scattering light collection optics, spectrometer, and TV camera.

Since the Thomson scattering measurement of $T_e(r)$ and $n_e(r)$ is the most reliable and the most localized measurement of these profiles, there is considerable impetus to develop appropriate high-power multi-pulse or cw lasers in order to follow the time evolution of the profiles. To date, little progress has been made in this direction. An interesting problem with future multi-channel laser systems is that non-circular cross section plasmas and tokamaks that utilize major radius compression are not well diagnosed by measuring T_e and n_e along a vertical chord; the precise major radius location of the central flux surface may not be well known, and will change in time. Indeed a measurement of the location of the magnetic axis may be helpful for directly deducing β . Thus in future tokamaks, the laser beam must cross the discharge in the horizontal midplane, and the polychromator must view the beam from above or below. This geometry presents a host of (hopefully soluble) problems with regard to access and beam dump design.

Laser techniques have been proposed for a number of other diagnostic purposes: laser-induced fluorescence for neutral atom density measurements, Faraday rotation of the polarization of a far infrared laser beam for the measurement of B_θ , and collective Thomson scattering in the far infrared for Z_{eff} and ion temperature measurement. To date, however, none of these advanced techniques has as yet been conclusively demonstrated. On the other hand, HCN (*Ref. 4-102*), alcohol (*Ref. 4-103*) and CO_2 (*Ref. 4-104*) lasers have been used as the basis for sophisticated multi-channel interferometers capable of operation at extremely high electron densities. Also cw CO_2 lasers have been used in near-forward coherent scattering experiments to measure plasma density fluctuations, in much the same way that large-angle microwave scattering has been used. Finally, in a very difficult experiment on DITE (*Ref. 4-105*), the possibility of using 90° Thomson scattering of ruby laser light for a diagnostic of $B_\theta(r)$ has been demonstrated. In this experiment, one takes advantage of the fact that the scattered light polarized precisely perpendicular to $\vec{B}$ is broken up into distinct harmonics of ω_{ce} . Using a tuned Fabry-Perot interferometer, the scattered peaks are superimposed, and the direction of $\vec{B}$ at the scattering point, and therefore the magnitude of B_θ , can be determined.

4.4.4 Neutral Particle Diagnostics

The basic diagnostic for measuring ion temperature on tokamaks has been the neutral particle analyzer, or charge-exchange detector. This detector measures the energy spectrum of neutral hydrogen atoms emitted from the plasma column. The H^0 atoms are generated from resonant charge-exchange, which permits a hot magnetically confined ion to pick up an electron from a neutral atom within the plasma column. The resulting hot neutral then has a chance of escaping the discharge without becoming ionized, and carries out with it information about the ion energy distribution in the plasma core. By detecting the fast neutrals, and determining their energy distribution, one can reconstruct the plasma ion temperature.

The neutral atom density at the periphery of the discharge is much greater than the density on axis, because the hydrogen atoms, which cascade inwards through successive generations of charge exchange, are simultaneously strongly attenuated by electron impact ionization. As a result, most of the charge-exchange neutrals are created near the outside of the discharge. In order to discriminate against these low energy neutrals, one typically measures the neutral atom flux in the range of 3–10 times $T_i(0)$. A measurement of this sort generally reflects the ion energy distribution in the central region of the discharge, and gives a good determination of the central ion temperature.

A problem with this technique is the reliance on the Maxwellian nature of the ion energy distribution function out to very high energies. In the case of neutral beam injection heating, it is certainly true that the distribution of the injected ion species is strongly non-Maxwellian,

and in various rf-heating schemes one also induces a strongly non-Maxwellian distribution. Fortunately, however, when one injects one neutral hydrogen isotope into a plasma consisting of a second isotope (e.g. energetic H^0 injection into a D^+ plasma), the energy distribution of the second isotope remains Maxwellian as it is heated (*Ref. 4-100*). In addition, with ICRF minority species heating, one finds that the majority species remains highly Maxwellian (*Ref. 4-106*). As a result, a charge-exchange detector with mass discrimination is able to provide good ion temperature measurements even when one of the species present has a strongly distorted energy distribution.

Figure 4-8 illustrates one approach to constructing a mass-discriminating charge-exchange system (*Ref. 4-107*). The fast neutrals are stripped to ions in a gas cell, typically filled with hydrogen, helium or nitrogen. A magnetic field is then used to provide momentum analysis of the ions, and a rather crude electrostatic energy analysis provides mass separation. The analyzed ions are usually detected with channeltrons operating in the pulse-counting mode, and count rates are stored in a computer.

Charge-exchange systems of this sort are also used to study the thermalization of injected beam ions. In this case, the circulating beam ions undergo charge exchange with neutrals in the plasma, and the outflux of beam neutrals as a function of time, angle, and energy provides information about the slowing-down rate and the pitch-angle scattering of the beam ions within the plasma. In turn, this information can be used as a measure of various plasma parameters, for instance Z_{eff} , which can be determined from the fast ion angular scattering rate.

Since charge-exchange detectors require large apertures into the tokamak vacuum chamber, tritium contamination will be a concern in reactor applications. As tritium builds up in the analyzer

Fig. 4-8 Layout of the ISX mass analyzing charge-exchange detector showing the principal components.

vacuum systems, or in the electron multipliers used as ion detectors, β particles emitted from the tritium will cause spurious counts. Without careful vacuum engineering, this appears to be a potential problem even in TFTR (*Ref. 4-108*).

4.4.5 Particle Beam Techniques

The two neutral-particle diagnostic techniques described in the last section, the charge-exchange determination of ion temperature and the measurement of fast ion thermalization rates, can both be substantially improved through the use of relatively small (5–15 A, 15–80 keV) collimated neutral hydrogen beams. In the case of the ion temperature measurement, the neutral hydrogen beam provides a target for charge-exchange reactions with the thermal ions. By using a time-modulated "doping" beam, and then detecting the modulated neutral flux at the analyzer, one can make a local measurement of ion temperature in the region of the plasma where the beam path and the detector sight-line cross. At very high density, or in very large plasmas, the intrinsic neutral atom density of the core region of the discharge is very small. As a result, the neutral flux emanating from the plasma core is so low that, without a doping beam, no measurement is possible at all. For very large, high density plasmas, there is the further complication that neutrals generated in the plasma core are very likely to be reionized on their way out to the detector. This problem appears to be surmountable, however, through the use of relatively high-current doping beams and neutral particle analyzers which view a large solid angle, as well as a large spectral region.

The use of a collimated neutral beam in the beam-ion thermalization measurements is quite different. In this case the purpose is to create a radially localized population of fast ions, so that measurements of their orbits and pitch-angle scattering rate give local measurements of plasma parameters. To this end, the diagnostic neutral beam is aimed into the plasma tangentially, outside of the center of the discharge. As a result, a "hot ring" of beam ions is created near the minor radius of the tangency point of the injected beam. Since fast ion orbits are displaced from the magnetic surfaces by a quantity $\Delta = qv/\Omega_{ci}$, a measurement of the orbit shift of the "hot ring" gives a local measurement of q . This technique has been used successfully in the ATC tokamak to determine q near the magnetic axis (*Ref. 4-109*). In addition, measurements of the pitch-angle scattering rate of the ions localized in the "hot ring" have been used in ATC to determine $Z_{eff}(r)$ (*Ref. 4-110*). The penetration of hydrogen neutral beams through tokamak plasmas has also been used as a diagnostic for plasma composition, and therefore Z_{eff} . Recent calculations and measurements, however, indicate that the impurities likely to be found in these plasmas have beam stopping powers very similar to the stopping power of the hydrogen ions they replace, thus making the measurements extremely difficult to interpret.

High energy ($\lesssim 100$ keV), low current (a few microamps) beams of heavy ions are used in another diagnostic technique. In the ST tokamak (*Ref. 4-111*), a 100 keV Th^+ beam was injected into the plasma column, where it was further ionized to Th^{++} . By detecting the outgoing Th^{++} beam current as a function of position and energy, it was possible to determine the electric potential across the plasma cross section, and also to measure local fluctuations in the plasma density. This technique presents problems in large high-field, high-density plasmas, due to the difficulty of obtaining adequate beam penetration, and because it requires the beam ion Larmor orbit to be comparable to the plasma minor radius. Through the use of a neutral injected beam, it will nonetheless always be possible to make measurements near the plasma surface.

A final particle-beam technique, which has had a successful preliminary trial on Pulsator, is the measurement of the poloidal magnetic field from Zeeman splitting of neutral lithium. In this technique, a neutral lithium beam is injected into the plasma, and the σ -components of a Zeeman

split resonance line are measured. By determining the polarization of this radiation, the direction of $\mathbf{B}$, and therefore the value of B_θ , can be measured. Unfortunately, this method cannot be used to measure $q(0)$ since B_θ vanishes on axis. The advantages of the approach are its conceptual simplicity and its potential for application in large, high density plasmas.

4.4.6 Spectroscopic Measurements

Visible and ultra-violet spectroscopy are the most direct methods for measuring plasma composition. Impurity atoms within the plasma are collisionally excited, and emit characteristic line radiation, whose intensity can be related through atomic physics calculations and measurements to the absolute impurity density. In addition, the excitation rates for certain transitions bear an unambiguous relationship to the ionization rates in some cases, so the source of impurities of a given ionization state can be determined from absolute measurements of radiation intensity. When this technique is applied to hydrogen atoms as well, a direct measurement of the total electron source, and therefore the particle containment time, is possible.

A second important application of spectroscopic techniques is the measurement of ion temperature from the Doppler-broadening of line radiation. Except for extremely high temperatures ($\gtrsim 5$ keV) this measurement is impossible in the vacuum ultra-violet, due to the limited resolution available. This limits the measurement to visible and near ultra-violet lines ($\lambda \gtrsim 600 \text{ \AA}$) which has, until recently, meant that the ion temperature could only be determined spectroscopically in cold regions of the plasma, since the dominant radiation from highly stripped impurities is in the vacuum ultra-violet. Recently, however, forbidden line radiation at $2665 \text{ \AA}$ from Fe XX has been identified (*Ref. 4-112*), and used for ion temperature measurements in the hot core region of the plasma.

The principal limitation of spectroscopic techniques is the fact that fully stripped impurities do not generate line radiation, and so cannot be detected through spectroscopy. For light impurities, which often make the largest contribution to Z_{eff} , this means that the impurity density cannot be measured over much of the plasma column. In addition, since recombination rates are typically very slow, light impurities could be diffusing into the plasma center, becoming fully stripped, and then traveling out again, and contributing, undetected, to Z_{eff} even in the cold region of the plasma. A possible resolution of this problem has come from an unexpected quarter. As a consequence of the worry about the stopping power of impurities for the fast neutrals used in neutral beam injection heating, much theoretical and experimental effort has been directed towards determining the cross-sections for charge-exchange between various tokamak impurities and energetic neutral hydrogen atoms (see Section 3.2.2). It appears that the reaction rates are high enough that the modest diagnostic beams used for doping the plasma with neutrals in the charge-exchange measurement of T_i should be expected to generate a substantial local population of hydrogen-like impurities through charge-exchange with fully stripped ions. Measurements of beam charge-exchange with low ionization states of oxygen on ORMAK (*Ref. 4-113*) lend credence to this expectation. No direct experiments however, have been made along this line.

The main technical difficulty with spectroscopic measurements on tokamaks is a reflection of their primary advantages: good spectral and spatial resolution. Because of the multiplicity of different lines each impurity radiates, it is impractical to attempt to compile a complete survey of the visible and ultra-violet radiation spectrum from the plasma with the usual technology of monochromators and photomultipliers. Photographic plates are invaluable for providing spectral information, but they do not provide adequate time resolution. As a result, typically a few lines are measured from each of the major impurities, and the overall radiation level is estimated from

models and experience. The introduction of micro-channel-plate, vidicon, and CCD technology into spectroscopy, however, will make much more information available from each discharge. The first example of the use of one of these potentially revolutionary technologies is a multi-spatial-channel, micro-channel-plate based, ultra-violet spectrometer used on Alcator A (*Ref. 4-114*). This instrument provides radial profiles of ultra-violet emission at a given wavelength as a function of time on a single-shot basis, rather than requiring a tedious shot-by-shot scan. It is reasonable to suppose as well that polychromators with a very large number of wavelength channels will soon come into use. These developments are strongly driven by the need to make measurements with a minimal number of machine shots, and by the fact that other diagnostic systems are progressing rapidly toward being able to make complete measurements in a single discharge.

In common with other diagnostic systems which use electron multipliers based on secondary electron emission, ultra-violet spectrometers are subject to noise from hard X-rays and γ -rays. In a reactor environment, (n,γ) reactions in neutron shield materials can be expected to cause a significant spurious background signal, which must be blocked out after the neutrons are attenuated. Already, a small but detectable "neutron noise" has been observed on micro-channel-plate electron multipliers in beam heated PLT discharges.

4.4.7 X-Ray Measurements

At higher electron temperatures, more plasma radiation is found in the soft X-ray region, and less in the visible and ultra-violet. Inner-shell line radiation, free-free bremsstrahlung, and free-bound recombination radiation are all found predominantly in the photon energy range near, and somewhat above, the electron temperature. Three different experimental techniques are generally used to study this radiation: pulse-height analysis for modest resolution spectral information and crystal spectrometry for high resolution over a narrow spectral range, both on a slow time scale, and analog spectrally integrating detectors, for very fast time resolution.

The pulse height analysis systems are typically based on a liquid-nitrogen cooled lithium-drifted silicon detector, which generates electrical pulses whose amplitudes are proportional to the energy of incoming X-ray photons. These pulses are amplified and sorted electronically and then stored until a full energy spectrum is accumulated. The principal difficulty with these systems is that, due to the exponentially decreasing flux of bremsstrahlung and recombination radiation as a function of energy, most of the X-rays are found near the low energy end of the spectrum. As a consequence, when the aperture is opened wide enough to allow good statistics at the high energy end of the spectrum, low energy pulses tend to pile up on one another too rapidly to be discriminated by the electronics, and the resulting measurement is significantly biased. With a multi-channel system (*Ref. 4-115*) it is possible to circumvent this problem by using large apertures for high energy photons, with filters to block out the low energy flux, and narrow apertures without filters primarily for the low energy photons. With these pulse height analysis systems it is possible to measure a broad region of the continuum spectrum. The magnitude of the flux reflects the overall impurity level, and the energy dependence determines the electron temperature. In addition, inner shell line radiation from incompletely stripped heavy impurities is superimposed on the X-ray continuum, and can be used to measure absolute impurity concentrations.

For more detailed information on the line radiation, an X-ray crystal spectrometer composed of a large semiconductor crystal used for Bragg scattering, coupled with a position sensitive X-ray detector (such as a multi-wire proportional counter) is used. With these systems it is possible to resolve details of a cluster of closely packed lines, from which one can deduce the charge state distribution of a given impurity. Furthermore, with very high resolution instruments, it is possible to

make Doppler-broadening measurements of individual lines, and so deduce ion temperature (*Ref. 4-116*). By choosing a line of an impurity with a high-ionization potential such as helium-like iron, one can thereby make ion temperature measurements in the very hottest central region of the plasma.

Since the spectrally integrated X-ray emission is very sensitive to electron temperature and to impurity content, measurements of rapid fluctuations in the X-ray emission give information about fluctuations in the plasma. In order to study MHD behavior in the internal region of tokamak plasmas, analog X-ray detectors, typically silicon surface-barrier detectors, are assembled in large arrays to view the plasma along as many chords as can be arranged. Insofar as possible, one attempts to imitate the techniques used in medical X-ray tomography, where very complete chordal measurements permit reconstruction of detailed internal structures. With the use of detectors located in vacuum, and with various thin foils used as X-ray filters, it is possible to obtain at least qualitative spectral information, and fluctuations can be measured on time scales as fast as one microsecond. These measurements are currently being used to make detailed comparisons with MHD theory (*Ref. 4-117*), and to provide insight into the mechanism, or mechanisms, of plasma disruptions (*Ref. 4-118*).

It is worth noting that X-ray detectors are the most neutron-sensitive of all tokamak diagnostics. Since X-rays are extremely difficult to deflect or focus, these detectors probably must in some way face into the neutron-producing plasma, and in the case of a reacting D-T system, this will result in very rapid destruction of any semiconductor device. Thus the technology of tokamak X-ray measurements, which has relied heavily on electron-hole production in semiconductor materials, will need to be drastically re-oriented for D-T devices.

4.4.8 Electron Cyclotron Emission

Perhaps the most exciting recent advance in tokamak diagnostics has been the advent of measurements of electromagnetic radiation from the plasma at multiples of the electron cyclotron frequency. Simple theoretical considerations (*Ref. 4-119*) suggest that even at modest densities tokamak plasmas should emit extraordinary mode radiation at $2\omega_{ce}$ essentially at the blackbody level. There are also indications (*Ref. 4-120*) that the emission at ω_{ce} is blackbody as well. Thus spectrally resolved measurements of the natural plasma emission in the millimeter microwave region and in the very far infrared constitute time and space resolved measurements of electron temperature. Higher harmonics may also be used to study non-Maxwellian features of the electron distribution function.

In order to measure $T_e(r,t)$, one must view the plasma along a major radius and cover a spectral range of width $\Delta\omega \sim \omega$. For the ω_{ce} diagnostic, one has the inconvenience that the view must be from the smaller major radius side. Heterodyne microwave detection techniques can be used for ω_{ce} in moderate field devices, while quasi-optical far-infrared technology is used for $2\omega_{ce}$ in most cases, and for ω_{ce} measurements in high field experiments. Here one uses either a grating polychromator (*Ref. 4-121*), or a scanning Michelson interferometer (*Ref. 4-122*). The infrared sensors are cryogenically cooled Putley detectors (*Ref. 4-123*) or similar devices. Any of these systems must be calibrated in situ against the Thomson scattering measurements. While they do not have the spatial resolution of the laser scattering due to finite spectral resolution and viewing angle, nor the absolute calibration, the ω_{ce} and $2\omega_{ce}$ diagnostics are particularly useful for electron temperature measurements in auxiliary heating experiments, where good time resolution, and good relative measurements of fairly modest changes in temperature are often required.

An important difficulty with the electron cyclotron emission diagnostic for the next generation of experiments is the fact that it relies on the unambiguous relationships between $|B|$ and R in a low- β tokamak. As β increases, $|B|$ no longer falls as $1/R$, and in extreme cases a local magnetic well can be created, making $|B|$ no longer monotonic with major radius along a line of sight. In this case, to provide complete T_e profiles from the edge to the magnetic axis, the plasma must be viewed from the inside in major radius even for the measurement at $2\omega_{ce}$, which could be problematic. In addition, a rather complex analysis will be needed to deduce the spatial dependence of T_e from the emission spectrum. On the other hand, such an analysis will be needed to some degree in any case, in order to understand the equilibrium in high β experiments. Because of their exceptional capability of providing complete profile information as a function of time on a single shot, it seems likely that electron cyclotron emission diagnostics will soon set the standard by which other, more cumbersome measurement techniques will be judged.

4.4.9 Bolometry

Visible and ultra-violet spectrometers, and even very soft X-ray detectors all suffer from the fact that they can only measure a limited region of the electromagnetic spectrum during an experimental run, either due to the intrinsically limited spectral range of the particular instrument, or because of its narrow spectral resolution, and the finite number of machine pulses available. With bolometric techniques one attempts to overcome this problem by measuring the total radiation power emitted by the plasma. The usual experimental technique involves exposing to the plasma a thin metal foil, with a thermister attached to the back. The total radiation power deposited is found by measuring the increase in temperature of the foil as a function of time. In the spectral range of interest, most metals are quite black (*Ref. 4-124*), so one need not be concerned either about reflecting incident radiation from the detector, or about radiation reflected from the far wall. It is important, however, to shield the detector from direct contact with the plasma in the shadow of the limiter.

By using a scannable bolometer with a collimated view, one can determine the chordal integrated radiation pattern, and through standard Abel inversion techniques find the volume radiation emission rate. This gives directly the volume sink of electron energy, which is critical information for power balance calculations used in the understanding of electron thermal transport. Perhaps the most dramatic example of the use of bolometers for the understanding of electron power flows comes from the hollow electron temperature profile experiments on PLT, using tungsten limiters (*Ref. 4-125*), where it was possible to demonstrate a local detailed balance between ohmic input power and radiation loss, an identification that played an important part in the PLT experimental plan.

The principal difficulty with present bolometric techniques is that it is not now possible to separate radiation impinging on the detector from the neutral particle flux due to charge-exchange. A thin carbon foil which would attenuate most of the charge-exchange flux would also block out a significant portion of the plasma radiation. This problem is aggravated by the fact that the neutral particle emission from the plasma violates the assumption of transparency needed for Abel inversion, so a small error in the measurement may be amplified in the inversion process. However in many cases the power carried out by charge-exchange is expected to be much smaller than the radiated power, so this difficulty is circumvented.

As with other diagnostics, bolometry is also moving towards multi-channel systems, rather than single-channel scanning procedures. This will permit studies of individual discharges, investigation of asymmetries, and immediate feedback to the machine operators on overall impurity conditions.

4.4.10 Neutron Diagnostics

As tokamaks approach closer to reactor conditions, and the neutron production rate climbs, a whole range of new neutron diagnostic techniques becomes available. This is particularly true for the case of auxiliary heating by neutral beam injection of D^0 beams into D^+ plasmas, where the neutron emission rate may rise by five orders of magnitude over its pre-injection value. In this case high time resolution measurements, spatially resolved measurements, and spectral measurements all become possible. In a D-T reacting plasma the possibilities are expanded even further, but the technology of detecting and shielding against 14 MeV neutrons is considerably more difficult than what is required at 2.5 MeV.

The standard diagnostic tool for measuring gross neutron flux from a tokamak in order to determine ion temperature is the moderated BF_3 "long counter". This is a proportional counter filled with BF_3 gas, which can absorb a thermal neutron through the $^{10}B(n,\alpha)^7Li$ reaction. The reaction products have a total energy of 2.8 MeV, and so generate ionization in the gas, and the resulting voltage pulse causes a count to be registered. The neutrons coming into the detector are thermalized by passing through a paraffin or polyethylene shield around the counter, and the whole apparatus is shielded against X-rays with a lead casing. Hard X-ray pickup is further reduced through pulse-height discrimination in the counting electronics. In order to be able to measure neutron fluxes over a wide range in energies, an array of detectors of different sensitivities is sometimes employed (Ref. 4-126).

To obtain neutron energy spectra, recoil proton spectroscopy in liquid scintillators or in gas proportional counters can be used. Alternatively 3He detectors employing the $^3He(n,p)^3H$ reaction can be used. In any case, the best spectra are obtained by locating the spectrometer inside a massive neutron collimator (Ref. 4-127), which permits the detector to observe unmoderated neutrons from the plasma column, but shields out neutrons scattered from surrounding apparatus. In order to obtain spatial resolution of neutron production, nuclear emulsions can be located close to the vacuum vessel (Ref. 4-128), and the longest recoil proton tracks show the direction of incidence of the neutron flux.

One critical difficulty with neutron diagnostics is the problem of absolute calibration. The usual technique is to place a PoBe source of known strength in the vacuum vessel, in order to estimate the response of the detector system. The need to move the neutron source around the torus and then average the detection efficiency can be eliminated by locating the BF_3 counters along the machine axis of symmetry. Unfortunately, no available neutron source has a simple emission line at 2.5 MeV, so the scattering effects and resulting detection efficiency can only be determined within a factor of ~ 2 . In the ranges of interest at present, however, neutron production scales approximately as T_i^4 , so a factor of two uncertainty in neutron rate amounts only to a 20% uncertainty in T_i .

At relatively low ion temperatures and low plasma densities, photo-neutron production in high-Z limiter materials can compete with thermonuclear neutron production (Ref. 4-129) and produce significant errors. Gamma rays produced by runaway electrons stopping in the limiter blades can detach neutrons from the heavy nuclei. This process can be discriminated against by locating two BF_3 counters symmetrically with regard to the plasma, but with one much closer to the limiters. A higher count rate in the detector closer to the limiters is then a sign of strong photo-neutron production. A more insidious effect is direct electro-disintegration of deuterium in the plasma by runaways, which will be a symmetrical process. This can only be detected through neutron spectroscopy, but in discharges that are not dominated by runaway effects, it appears negligible.

4.4.11 Plasma—Wall Interactions

The traditional application of surface physics techniques to tokamak studies has been to insert probes into the vacuum vessel which are designed to simulate a part of the vessel wall. The probe is then exposed to tokamak discharges, or to discharge cleaning, and changes in its surface condition are monitored, through a number of different sensitive surface physics techniques. The removal of light impurities by discharge cleaning can be monitored as can, for instance, the deposition of limiter material on the walls by high-power discharges. This provides a basis for evaluating different approaches to impurity removal, and for understanding the interaction of the plasma with material surfaces.

An alternative approach to studying tokamak plasmas with surface physics techniques is to insert a probe into the scrape-off plasma in the shadow of the limiter (*Ref. 4-130*). With this technique one can monitor the flow of impurities both along and across field lines in this region. By using a shuttered aperture, or a moving target behind a continuously open aperture, one can make time-dependent measurements of these fluxes, which can then be correlated with plasma behavior and spectroscopic measurements in order to study the processes by which impurity atoms are introduced into the plasma.

4.5 REACTOR DESIGN

Reactor studies play a central role in the development of any new power source. They serve to establish the technical utility of the scientific program, to provide a continuous monitor of the economic competitiveness of the reactor concept, and to guide the development and technology programs. For nearly a decade, groups throughout the world have been engaged in studies of reactor systems based on the tokamak concept. These studies have spanned a broad range of devices and applications. Near-term devices of the ignition test reactor class have been examined (*Refs. 4-131 through 4-135*) as well as long-range devices of the power reactor class (*Refs. 4-136 through 4-146*). The tokamak has also been proposed as the fusion core of hybrid fusion-fission energy systems (*Ref. 4-147*) in the context of both fissile fuel factories and fission power generators. Moreover, within any of the classes of devices or applications described above, a variety of fuel cycles may be contemplated (*Ref. 4-148*). Thus, tokamak reactor design is an extremely broad subject and for the purposes of this document it appears appropriate to limit the scope of the discussion to D—T fueled systems of the power reactor class.

In D—T fueled reactors, the fusion reaction of primary interest is

Provided the plasma current is sufficiently large (≥ 2.5 MA), the bulk of the alpha particle energy will be deposited within the plasma, leading to self-sustaining or ignited operation if plasma confinement and the ion temperature satisfy certain criteria. Since the plasma is virtually transparent to the high energy neutrons, a blanket is placed around the plasma to convert the neutron energy into high grade heat for power production. The radioisotope tritium (half-life ~ 12.6 years) is a limited resource and must be manufactured or bred by the reactor itself. This would be accomplished in a lithium based blanket via the neutron-induced reactions

and

The D-T fuel cycle has the disadvantages of requiring the handling of tritium and producing substantial amounts of neutrons that result in radiation damage and induced radioactivity within the reactor structure. These considerations have considerable impact on the design of D-T reactors, especially in the areas of materials, radiation shielding, maintenance and safety. It is to be noted that such problems, which may well govern the timetable for commercialization, are generic to the D-T fuel cycle rather than specific to tokamaks.

Figure 4-9 shows schematically the major systems of a D-T tokamak reactor. The magnetic systems consist of both a steady-state toroidal field system and a pulsed poloidal field system. Power balance considerations require that the toroidal field coils be superconducting. Poloidal field design on the other hand is somewhat more flexible; while the ohmic heating or induction coils are likely to be superconducting, the coils used to shape and position the plasma may well be a mix of superconducting, normal, and cryoresistive coils. Auxiliary heating will be required to reach the ion temperature needed for ignition; currently, neutral beam injection and several varieties of radio-frequency heating are under consideration for this purpose. In order to sustain long plasma burn times, fresh fuel will have to be introduced into the fusing plasma. Both pellet injection and gas puffing are being studied as means of plasma fueling. The impurity contamination of the plasma must be maintained at a sufficiently low level to permit ignition and to sustain an economically acceptable burn time. Both divertor and non-divertor approaches are being pursued as means for plasma purity control in tokamaks. The bulk of the fusion energy would be deposited within the blanket where it would be recovered as heat through conventional techniques. Both steam cycles and direct cycle helium power conversion systems are being considered. Handling and control of substantial quantities of tritium is required not only for the tritium bred in the blanket, but also for the unburned tritium contained in the plasma exhaust stream. The blanket structural

Fig. 4-9 The major systems of a D-T tokamak reactor.

system must be capable of giving economically acceptable service life in the presence of a harsh radiation environment and under pulsed conditions. A number of structural alloys are being considered for this application. The pulsed power delivery system drives the poloidal field system in a tokamak. In current experiments, motor-generator flywheel sets are employed for this task. Homopolar generators and superconducting energy storage systems are also being considered for this application. The pulsed power delivery system drives the poloidal field system in a tokamak. In current experiments, motor-generator flywheel sets are employed for this task. Homopolar generators and superconducting energy storage systems are also being considered for reactor applications. A reactor embodiment of the systems identified in Fig. 4-9 is illustrated in Fig. 4-10, which shows a cross-sectional view of NUWMAK, a recent reactor design (*Ref. 4-146*).

Reactor design studies have been carried out to varying levels of detail and with varying degrees of optimism with regard to both the extrapolation of our current understanding of the central physics issues and the future course of technological development. For most such investigations, plant economics is the overriding concern with the focus being on identifying and making advances in cost-sensitive areas. Some general considerations in reactor economics are discussed in Section 4.5.1. Section 4.5.2 deals with some of the unique features of tokamak power plants and the subsequent implications for reactor design. The data base needed for reactor design is discussed and compared to the information flow anticipated from the existing fusion program in Section 4.5.3. Finally, Section 4.5.4 enumerates a list of critical issues to be explored in future studies.

4.5.1 Economics Considerations in Reactor Design

In this section, the role of plant costs in reactor designs is reviewed (*Refs. 4-144, 4-149*). Power generation cost is generally divided into three components: capital cost, fuel cost, and operation and maintenance cost. For example, the light water fission reactor power generation cost is presently distributed with about 70% capital cost, 20% fuel cost, and 10% operation and maintenance cost. Fission power generation is said to be highly capital-cost intensive, i.e., the power cost is to a large extent determined by capital cost investment and to a lesser extent by fuel cost. Fusion reactors will be even more capital-cost intensive. About 90% of the power generation cost for fusion reactors will be capital expenditures and about 10% will be operation and maintenance; the fuel cost of fusion power should be less than 1% of the total. The capital cost (in mills per kilowatt hour) is derived from the installed plant cost (dollars per kilowatt electric) and the plant capacity factor (the time averaged ratio of the average generated load and the total rated capacity of the plant). The capacity factor is a measure of the plant performance, which reflects system reliability and maintainability.

In assessing the evaluating plant costs it is important to consider both the total cost (\$) and the unit cost (\$/kW_e). For a given choice of blanket structure, blanket coolant, and power conversion system, the balance of plant cost will be relatively insensitive to the details of the tokamak energy source. In this case, the reactor systems cost of the tokamak power plant offers the greatest potential for cost reduction. Some general trends in tokamak reactor system costs can be deduced by considering the following set of simplified cost scalings:

$$\begin{array}{l} \text{Reactor Systems} \\ \text{Total Cost (\$)} \end{array} \sim \begin{array}{l} \text{Volume of} \\ \text{Reactor} \\ \text{Structure} \end{array} \sim a^2 \Delta \quad (4-18)$$

Fig. 4-10 Cross-sectional view of NUWMAK.

where a is the plasma radius and Δ is the total thickness of reactor structure including the blanket, shield and magnet systems. Thus,

$$\begin{aligned} \frac{\text{Reactor Systems Unit Cost (\$/kW}_e)}{\text{Reactor Power}} &\sim \frac{\text{Volume of Structure}}{\text{Reactor Power}} \sim \frac{a^2 \Delta}{Pa^3 \eta} \\ &\sim \frac{\Delta}{Pa \eta} \end{aligned} \quad (4-19)$$

where P is the fusion power density and η is the overall plant efficiency. Since power density times radius is proportional to wall loading, L_w , reactor systems unit cost can be written as

$$\frac{\text{Reactor Systems Unit Cost (\$/kW}_e)}{\text{Reactor Power}} \sim \frac{\Delta}{L_w \eta} \quad (4-20)$$

Note that materials and thermal-hydraulic considerations will set an upper bound on the value of L_w . Furthermore, Δ is determined to a first approximation by the penetration of 14 MeV neutrons in matter.

It follows from Eq. (4-18) that low reactor systems total cost is favored by small plasma radius. From Eq. (4-20) we see that low reactor systems unit cost is favored by high wall loadings. For a fixed value of the wall loading, low reactor systems total cost is favored by high power density and, therefore, small plasma radius. Naturally, high overall plant efficiency also reduces unit cost, assuming that no cost penalty is incurred by the technology advancement required to achieve the increase in η .

In the following, reactor design characteristics are discussed relative to the observations cited above. Because the overall plant efficiency of fusion power plants is intimately tied to the plasma operating characteristics, this issue is discussed separately from the issues of size, power density and wall loading.

4.5.1.1 Size, Power Density and Wall Loading — On the basis of the preceding discussion, it is concluded that more competitive reactor systems costs, that is, both total and unit costs, are favored by small plasma radius, high power density, and high wall loading. Indeed, the trends in tokamak reactor design studies have been in these directions. These trends are indicated in Table 4-2, which compares the parameters of interest for the three University of Wisconsin designs (*Refs. 4-141, 4-142, 4-146*), UWMAK-I (published in 1973), UWMAK-III (published in 1976), and NUWMAK (published in 1978). It is seen that the plasma volume has been reduced by a factor

TABLE 4-2
THE TREND IN REACTOR DESIGN STUDIES HAS BEEN TOWARDS REDUCED SIZE,
INCREASED POWER DENSITY AND INCREASED WALL LOADING

	UWMAK-I (1973)	UWMAK-III (1976)	NUWMAK (1978)
Plasma radius (m)	5	2.7	1.1
Plasma volume (m ³)	6400	2370	150
Power density (MW/m ³)	0.8	2.3	12
Wall loading (MW/m ²)	1.25	2.5	5.0

TABLE 4-3
CURRENT ESTIMATES INDICATE A MARKED IMPROVEMENT
IN TOKAMAK REACTOR SYSTEMS COST

	UWMAK-I (1500 MW _e)	ORNL/DEMO (750 MW _e)	LMFBR (1200 MW _e)	PWR (1000 MW _e)
Reactor Systems*	\$1000/kW _e	450	100	75
Balance of Plant	800	800	700	525
Total	\$1800	1250	800	600

*The reactor systems do not include the heat transport system. This system is included in the balance of plant.

of 40 in going from the UWMAK-I design to the NUWMAK design. Furthermore, the fusion power density in the NUWMAK design is a factor of 15 higher than that in the UWMAK-I design and the neutron wall loading is a factor of 4 greater in the NUWMAK design than in the UWMAK-I design. When current tokamak reactor designs such as the NUWMAK are compared in size to the next generation of large tokamaks (TFTR, JET, and JT-60), it appears that we may be reaching a limit in size reduction for tokamak power reactors. In any case, further reductions in plasma size may not be accompanied by a corresponding reduction in unit cost because a minimum thickness of roughly 1 meter of matter must surround the plasma to absorb the energy of the fusion neutrons.

Table 4-3 compares current cost estimates (in 1976 dollars) for the UWMAK-I design, the ORNL/DEMO design (*Ref. 4-144*), a liquid metal fast breeder reactor (LMFBR) (*Ref. 4-150*), and a pressurized water reactor (PWR) (*Ref. 4-151*). These power reactor plants were chosen for comparison since they all use stainless steel as the structural material of the reactor systems and a steam cycle as the power conversion. The cost estimates in Table 4-3 represent *direct* costs and do not include interest during construction or escalation.

The following points are noted with regard to the estimates given indicated in this table:

1. The ORNL/DEMO design showed a factor of 2 improvement in the reactor systems cost relative to the UWMAK-I design. This reflects the smaller size (430 m³), higher power density (5.3 MW/m³) and higher wall loading (3 MW/m²) of the ORNL/DEMO design.
2. The balance of plant costs are essentially the same in the UWMAK-I design and the ORNL/DEMO design. This reflects the similarity of the two designs in the balance of plant.
3. Fusion power plants can be expected to have considerably larger reactor systems cost than fission reactor plants. This reflects the more diverse technological requirements of the fusion reactor systems and the inherently lower power density.
4. The balance of plant costs should not be dramatically different between an LMFBR and a fusion reactor.

As a final point it is useful to consider the amount of structure in the reactor systems as an indicator of reactor systems cost. Table 4-4 compares the amount of stainless steel structure per kW_e for the UWMAK-I design, the ORNL/DEMO design, an LMFBR, and a PWR. This table

also shows that the structure ratio is very close to the cost ratio for the PWR and the LMFBR and also for the ORNL/DEMO and the UWMAK-I design. However, the cost ratio between the fusion and fission systems is about a factor of 2 greater than the structure ratio for these systems. This is probably due to the fact that the materials requirements of fusion systems are more demanding than those of fission systems and the structural requirements represent only part of

the total materials requirements for the fusion reactor systems. Nevertheless, the cost ratios and structure ratios are in reasonable agreement and the cost estimates indicated in Table 4-3 are not inconsistent with the amount of structure.

TABLE 4-4
THE AMOUNT OF STRUCTURE IN THE REACTOR SYSTEM
IS A USEFUL INDICATOR OF REACTOR SYSTEMS COST

	kg Structure/ kW _e	Structure Ratio	Cost Ratio
UWMAK-I	21	8.8	13.3
ORNL/DEMO	7	3	6
LMFBR	3.1	1.3	1.3
PWR	2.4	1.0	1.0

4.5.2 Unique Features of Tokamak Power Plants

In contrast to most other power sources, fusion power reactors are inherently power amplifiers since they require recirculating power to establish the fuel conditions necessary for appreciable power release. This feature may be quantified in terms of the parameter Q , designated the power amplification factor and defined as the ratio of the total nuclear power release, P_t , to the power input, P_i :

$$Q \equiv \frac{P_t}{P_i} \quad (4-21)$$

The essential features of the power balance in a fusion reactor are illustrated by the schematic diagram of Fig. 4-11. The power terms in this figure have the following meanings:

- P_i $\equiv$ power input required to initiate and sustain the plasma burn
- P_t $\equiv$ total nuclear power release
- P_g $\equiv$ gross electrical power output
- P_c $\equiv$ circulating power, that is, the power required by the confinement, heating, and sustaining systems
- P_n $\equiv$ net electrical power output
- η $\equiv$ overall plant efficiency.

The plasma confinement, heating, and sustaining systems are characterized by a power efficiency, η_i , and may incorporate pulsed as well as steady-state power components. The power conversion systems are characterized by a power efficiency, η_c , and may incorporate direct energy conversion as well as thermal energy conversion.

In order to achieve a useful overall plant efficiency, a fusion reactor must yield an attractive combination of Q and η_i (Ref. 4-152). This point is illustrated in Fig. 4-12, which gives the values of Q and η_i required to achieve an overall plant efficiency, η , of 30% when a power conversion efficiency, η_c , of 40% is assumed. For relatively low Q devices, $Q < 10$, the plasma systems power efficiency, η_i , must be high, $\sim 80\%$, to achieve a useful plant efficiency. For relatively high Q devices, $Q > 100$, the plasma system power efficiency can be low, $\sim 10\%$, and still allow a useful plant efficiency.

Fig. 4-11 Fusion power reactors can be viewed as power amplifying devices.

Fig. 4-12 In order to achieve a useful power balance a fusion reactor must yield an attractive combination of Q and η_i .

Another unique feature of tokamak power producers is their fundamentally cyclic nature. The length of the burn time will be limited by one of the following:

- Plasma quenching owing to one or more of the following factors: enhanced radiation associated with impurity buildup, fuel depletion resulting from the inability to replace fuel, and fuel dilution resulting from the inability to remove the helium ash.
- Loss of plasma control, perhaps through the evolution of the plasma current profile.
- The available magnetic flux through the center of the torus, a quantity which determines the duration of the plasma current.

TABLE 4-5
SUMMARY OF POWER CYCLE CHARACTERISTICS FOR SOME
CURRENT D-T TOKAMAK FUSION POWER REACTORS

	GA Doublet Prototype	ORNL/DEMO	MIT/PPL HFCTR	University of Wisconsin NUWMAK
Cycle time, sec	177	1260	588	218
Burn time, sec	160	1200	500	200
Down time, sec	17	60	88	18
Duty cycle	0.90	0.95	0.85	0.92

In the absence of fueling and ash removal capabilities, it appears that the burn time would be limited to tens of seconds; the constraint imposed by impurity accumulation may be of the same order. The plasma control issue is difficult to quantify and will likely require actual operating experience on reactor-grade plasmas since both the plasma profiles and the plasma energetics may be quite different in burning plasmas from present laboratory plasmas. The magnetic flux consideration, albeit a fundamental limitation, does not preclude burn times of many minutes because the one-turn voltage in a reactor grade plasma is very small; in addition, possibilities exist for eliminating the volt-second limitation on pulse length (see Section 4.3).

The time between burns, the so-called down time, is constrained by pumping requirements and energy storage and delivery considerations to be at least 10–20 seconds. Since a duty cycle in excess of 90% is deemed desirable, a burn time of several minutes is a necessity. Hence, it is essential to come to grips with the impurity control, fueling, and ash removal questions identified above in order to demonstrate the economic competitiveness of the tokamak as a reactor candidate. The power cycle characteristics of some current tokamak power reactor designs are summarized in Table 4-5.

Not the least of the issues associated with cyclic operation is the problem of cyclic fatigue. Although the design implications depend on the details of the burn cycle so that generalizations are difficult, there is the concern that cyclic fatigue may be the life limiting feature for some systems and hence adversely affect device availability.

4.5.3 Required Data Base

Reactor studies have provided considerable guidance to the tokamak program, both in assessing the relevance and completeness of the existing data base and by providing a framework to judge between programmatic alternatives. The data base assessment is particularly important for the design of near term, post-TFTR devices. Figure 4-13 presents a brief synopsis of the information flow anticipated from important elements of the fusion program. It is seen that the bulk of the information needed to make key reactor design decisions will be available by the mid-1980s.

From the point of view of plasma physics, the relevant elements of the existing program are the various currently-operating and planned tokamak major scaling experiments, including PLT and PDX at Princeton Plasma Physics Laboratory, Doublet III at General Atomic Company, Alcator C at Massachusetts Institute of Technology, and ISX at Oak Ridge National Laboratory. On a longer time scale, the key physics information is also anticipated from TFTR at Princeton Plasma Physics Laboratory and from the European (JET) and Japanese (JT-60) fusion programs. For comparison,

Fig. 4-13 Information flow anticipated from current and planned program elements.

TABLE 4-6
SUMMARY OF REACTOR PHYSICS REQUIREMENTS FOR SOME CURRENT D-T TOKAMAK
FUSION POWER REACTORS

	GA Doublet Prototype	ORNL/DEMO	MIT/PPL HFCTR	University Of Wisconsin NUWMAK
Major radius, m	6.8	6.2	6.0	5.1
Plasma radius, m	2.1	1.55	1.2	1.13
Aspect ratio	3.2	4.0	5.0	4.5
Height-to-width ratio	2.8	1.6	1.5	1.65
Shape	Doublet	Dee	Dee	Dee
Average toroidal beta, %	10	10	4	6
Peak ion temperature, keV	16	13	12.4	13
Peak ion density, 10^{20} m^{-3}	1.8	1.7	5.2	2
B_T on axis, T	3.7	3.4	7.4	6.1
Safety factor, q	2.5	3	3	2.6
Plasma current, MA	22	3.9	6.7	7
Z_{eff}	2	1.1	1.2	1-1.5
Burn time, sec	160	1260	588	200
Down time, sec	17	60	88	18

the reactor physics requirements of some current reactor designs are summarized in Table 4-6. It is interesting to note that for most parameters in this table the range among the various designs is within a factor of two.

It should be mentioned, however, that the presently planned program will not supply all of the physics information needed to guarantee that current reactor designs will achieve the specified plasma performance. In particular, none of the presently planned experiments will address the questions associated with ignition and post-ignition dynamics, for which the plasma energy balance will be dominated by alpha-particle heating, or the questions associated with long-pulse operation, such as impurity and ash accumulation. Indeed, several stages of devices intermediate between TFTR and commercial reactors will be needed to address such questions as well as the vital questions governing the economic competitiveness of the concept (*Ref. 4-153*).

The reactor technology requirements typical of several current design studies are exhibited in Table 4-7. Engineering and technology input to the program will come from a variety of sources. Solutions to a number of tokamak engineering problems associated with D-T operation will be tested on TFTR and the European device JET. The Large Coil Program at ORNL will provide engineering data and fabrication experience for the large superconducting toroidal field coils. Questions relevant to tritium processing and control will be addressed by the Tritium Systems Test Assembly at Los Alamos Scientific Laboratory. Critical and pertinent materials tests will be performed at Hanford Engineering Development Laboratories, Lawrence Livermore Laboratory, and Oak Ridge National Laboratory.

TABLE 4-7
SUMMARY OF TOKAMAK REACTOR TECHNOLOGY REQUIREMENTS
BASED ON CURRENT D-T REACTOR STUDIES

System	Requirements
Magnetic	
Toroidal field coils	Superconducting 8–12 T
Induction coils	SC/Normal 5–10 T/sec
PF pulsed power supply	1–2 x 10 ³ MVA 2–4 GJ
Plasma heating	50–100 MW
Plasma fueling	10 ²² –10 ²³ particles/sec with central fueling
Plasma purity control	If divertor: 10 ²² –10 ²³ particles/sec 1–2 kW/cm ² in collectors
Blanket structural	6–10 year lifetime at elevated temperature and in radiation field
Tritium handling	~10 kg T/day or ~10 ⁸ curies/day <10 curies/day leakage
Energy recovery and conversion	$\eta \geq 30\%$

4.5.4 Critical Issues for Future Tokamak Reactor Design Studies

Although considerable progress is being made in tokamak reactor design studies, it is important to identify a number of critical issues which must be pursued in future studies. A suggested list of issues is given below.

1. **Dealing with the pulsed nature of the tokamak** — How can the burn time be extended? Are there possibilities for making the tokamak operating mode steady state?
2. **Evolving a credible divertor particle-collection power-handling scheme** — Although a number of concepts have been proposed for the divertor, there is no concept which adequately deals with the problems of particle collection and power handling in a reactor context.
3. **Increasing the feasibility of tokamak maintenance** — What additional engineering options can be advanced which might ease maintenance in the complex geometry of tokamak designs? Can the geometrical constraints be relaxed?
4. **Developing a workable design for the first-wall blanket system** — A number of concepts have been advanced for the first-wall blanket system, but no approach has yet been carried through to the point where it represents an acceptable engineering design.

5. **Improving the coupling between the reactor design studies and the scientific and technological programs** — The ideas advanced in reactor systems studies must find a vehicle for verification within the scientific and technological programs. At the same time the results from the scientific and technological programs must be incorporated into the reactor design studies.
6. **Evolving a plan for implementation** — What intermediate steps are required between the present experiments and a demonstration of fusion power? How do we progress from our current technology base to the technology base required for reactors?

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