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TIBER II/ETR Final Design Report

Volume 1

1.0 Introduction 2.0 Plasma Engineering

Technical Editor: J. D. Lee

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Argonne National Laboratory (ANL)
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September 1987

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EXECUTIVE SUMMARY

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EXECUTIVE SUMMARY

1. INTRODUCTION

The Tokamak Ignition/Burn Experimental Reactor (TIBER) [1] shown in Fig. 1 was pursued as the U.S. option for an International Engineering Test Reactor (ETR). This concept evolved from earlier work on the Tokamak Fusion Core Experiment (TFCX) to develop a small, ignited tokamak [2]. While the copper-coil versions of TFCX became the short-pulsed, 1.23-m radius, Compact Ignition Tokamak (CIT) [3], the superconducting TIBER [4] with long pulse or steady state and a 2.6-m radius became a candidate for international collaboration. In this final report the design has been updated to TIBER II, to accommodate the conservative Kaye-Goldston confinement scaling, double-poloïdal divertors for impurity control, steady-state current drive, and nuclear testing.

While TIBER II satisfies the perceived mission for an ETR, the design has several additional objectives. The first is that it be reactor relevant. That is, the physics and technology embodied, including blanket test modules, must lead to an attractive practical concept. Second, steady-state operation should be achieved. Not only is steady state a likely condition for a power reactor, but it also may be necessary for engineering testing of nuclear components. Plasma cycling causes time-dependent changes in environmental testing conditions, such as thermal conditions, tritium-concentration profiles, eddy-current effects, and failure or fracture modes. In fact, pulsing effects can be more dominant than the steady-state effects for which testing is desired. At a minimum, pulsing complicates tests and complicates the extrapolation to reactor conditions. For these reasons, TIBER II is designed with sufficient auxiliary current-drive power for steady-state operation with the intention of stimulating further current-drive experiments in the U.S. and world fusion programs. As a contingency, however, a limited inductive drive capability was provided without increasing size thereby allowing a limited number of pulses.

A third design objective, low cost, is enhanced by steady-state operation. The usual large ohmic heating (OH) coil required in the center of the tokamak for conventional inductive operation leads to a larger overall device size. This is because the size of the center post area is determined

① Shield stabilizer	⑩ Poloidal field coils
② PF support structure	⑪ Upper support shield
③ Ohmic heating coils	⑫ Outer vacuum wall
④ Inboard shield	⑬ Toroidal field coils
⑤ OH support structure	⑭ Lower divertor
⑥ Outboard shield	⑮ Lower support shield
⑦ Blanket test module	⑯ Divertor vacuum ducts
⑧ Cryogenic support struts	⑰ Shield support pedestal
⑨ Cryogenic support plate	

Figure 1. TIBER II is the U.S. option for the international ETR.

by the OH coil radius, the toroidal field-coil current density, and the inner neutron shielding. It is necessary to minimize all of the dimensions (especially the OH coil) to keep the reactor size and cost low. Pulsed OH current drive would mandate lower magnet current densities, significant eddy-current heating loads, and, because of fatigue, the need for larger structures.

2. PLASMA ENGINEERING

The TIBER II design incorporates physics requirements and neutron wall-loading with fluence parameters that make it compatible with a nuclear testing mission. Reactor-relevant physics can be tested by using high-beta, non-inductive current drive, and steady-state operation. Although the design accommodates several current-drive options, the base case uses a combination of lower-hybrid and neutral-beams for the main current drive with electron-cyclotron radio-frequency power for localized control of the current profile.

2.1. Equilibrium and Stability

The plasma equilibrium for TIBER II was obtained from the NEQ equilibrium code [5]. Free-boundary and fixed-boundary calculations were iterated to minimize the number of poloidal field (PF) coils required to produce a desired plasma (specified by elongation κ and triangularity δ) corresponding to a crescent plasma shape of the form described by Yamazaki [6]. Profile factors for the plasma pressure, toroidal flux function, and poloidal beta were adjusted until an equilibrium was obtained with the desired average beta and safety factor on axis, $q(0)$. A coil configuration was determined that produces an equilibrium at the desired plasma current subject to maximum current-density limits on the superconducting coils. The resulting plasma shape, toroidal current, and volume-averaged beta value were then used in the power balance at the 95% flux surface. The density and temperature profiles assumed in the power-balance code are adjusted to give a pressure profile that is consistent with the equilibrium solution.

The operating beta is taken to be the lesser of the Troyon beta limit [7], given by $\beta_{crit} = 3(\%) I(\text{MA}) / a(\text{m}) B(\text{T})$ and 0.75 times the ballooning limit of Yamazaki [6]. The latter takes into account the plasma shape and aspect ratio.

Fast-alpha pressure and contributions from alphas and impurities at the bulk ion temperature are included in the beta calculation. The model used for the fast-alpha pressure [8] assumes that the alpha distribution function satisfies a steady-state kinetic equation with a monoenergetic source at the alpha birth energy. The fast alpha confinement time is assumed to be much longer than the slowing down time. Typically the fast alphas contribute 10 to 15% of the total beta.

2.2. Transport

Radial density and temperature profiles were assumed to be:

$$x(r) = \hat{x} [1 - (r/a)^2]^{\alpha_x},$$

where x represents either density n or temperature T , and $\hat{x} = x(r = 0)$. The same profile shapes are assumed for both electrons and ions, and the exponents α_n and α_T are adjusted so that the resulting pressure profile matches the pressure in the equilibrium solution. Particular density and temperature profiles used are not unique; other combinations are certainly consistent with the equilibrium pressure. Because the equilibrium flux surfaces are shifted outward in TIBER II, separate exponents α_n^- , α_n^+ and α_T^- , α_T^+ for the density and temperature profiles are used on the inboard and outboard sides of the magnetic axis, respectively. Somewhat peaked density profiles $\alpha_n^- = 1.1$ and $\alpha_n^+ = 0.9$ have been chosen in accordance with the Tokamak Fusion Test Reactor (TFTR) experiments and the assumption of pellet fueling of TIBER II. The corresponding temperature profiles consistent with the equilibrium pressure are $\alpha_T^- = 0.6$ and $\alpha_T^+ = 0.55$.

Gross energy confinement is taken to be described by the Kaye-Goldston scaling [9] times an enhancement factor H_f representing H-mode operation (CIT). The value of $H_f = 1.1$ required as an operating point control is somewhat less than the possible enhancement $H_f = 1.5$ to 2.0 for the type of double-null poloidal divertors used in TIBER II. It was assumed that all of the alpha power and the auxiliary power (current drive) degrade confinement. Electron and ion energy confinement are taken to be equal (ion diffusivity is much larger than neoclassical).

Finally, steady-state burn requires some mechanism by which the thermalized alpha particles can be transported out of the plasma. Operating points are calculated from power balance equations assuming a fixed concentration of thermal alphas (typically 5% of the fuel ion density). In steady state, this implies a thermal alpha particle confinement time that is long compared to the bulk energy confinement time. Fuel ions and thermal alphas are assumed to be removed at the same rate. The resulting convection power loss, even assuming fuel ions were lost with $5/2$ times the ion energy, would be a small fraction of the total energy loss calculated, assuming equal ions and electron diffusion rates. The level of turbulence required to produce the required alpha particle removal rate is thought to be consistent with the assumption that the ion thermal diffusivity is much larger than neoclassical.

2.3. Steady-State Current Drive

Several different methods of non-inductive current drive have been considered for TIBER. Of these, a combination of lower hybrid (LH) fast waves for the outer region and neutral beam (NB) injection for the core, has the best experimental and theoretical support and was chosen for TIBER. Additionally, electron cyclotron resonance heating (ECH) has proven to be useful for current-profile control to avoid disruptions. The baseline operation of TIBER includes 10 MW of ECH in addition to the roughly 45 MW of NB and 20 MW of LH power.

The efficiency of current drive is generally given by the quantity

$$\gamma_{nb} (A/W\text{-m}^2) \propto \langle n_{e20} \rangle I(A) R_o(m) / P(W)$$

where $\langle n_{e20} \rangle$ is the volume-averaged electron density in 10^{20}m^{-3} and P is the absorbed power required to drive the current I . A simple scaling law for NB current drive efficiency is [10,11]

$$\gamma_{nb} = \langle T_e \rangle J(\langle x \rangle, \langle y \rangle) F_e / \ln A_e$$

where $x^2 \approx E_b / 10 A_b T_e$, $y \approx 0.8 Z_{eff} / A_b$, F_e accounts for the reverse current carried by the electrons, and J is given approximately by

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

The above relation for γ_{nb} is used as a scaling relation for full computations carried out with a Fokker-Planck code [12]. The efficiency for LH current drive was computed using the relativistic formulation of Fisch [13]. This efficiency is degraded by electron trapping effects. This effect has been calculated with the formulation of Cohen [14]. A summary of the plasma parameters for TIBER is given in Table 1.

3. CONFIGURATION

Minimum neutron shielding, compact structures and divertor, and high magnet-current densities all contribute to compact tokamak size. By minimizing the TIBER II size, costs can be reduced. Most costs of the tokamak components, such as magnets, first-wall structures, and vacuum vessels vary with surface or the square of the major radius. Other components, such as plasma-heating systems, are related to volume and vary as the cube of the major radius. Fortunately, because facilities, diagnostics, and other auxiliary systems do not change as much, the net change is approximately linear with the radius.

3.1. Compact Divertor

Experiments and analyses on ASDEX, PDX, and D-III point to a better understanding of the conditions that a divertor must produce to allow good confinement of the hot plasma. In particular, it was found that the divertor configuration and the divertor plasma must prevent gas from penetrating into the confined plasma region. Experience from NET [15] and CIT [16] lead to the divertor configuration that is proposed for TIBER II (seen in Fig. 2). The newer designs allow for a more open configuration and rely on a dense edge plasma to intercept the gas produced at the neutralizer plate.

The TIBER edge plasma and compact divertor operate in a high-recycling mode. This mode generates the high density needed to allow the edge plasma to shield the core plasma from the gas impurities generated at the divertor plates. It also allows the charged-particle power to be carried out by a high

Table 1. TIBER parameters.^a

Major radius, R_0 -m	3.00
Minor radius, a-m	0.83 ^b
Elongation, κ	2.4 ^b
Triangularity, δ	0.4 ^b
Magnetic field, B_0 -T	5.55
Plasma current, I_p -MA	10
Total fusion power, MW	314
Q	5.0
Average (peak) wall loading, (MW/m ²)	1.3(2.0)
Neutral beam power, MW	44
Lower hybrid power, MW	19
Neutral beam driven current, MA	3.4
Lower hybrid driven current, MA	3.9
Bootstrap current, MA	2.7
LH current-drive efficiency, ^a $10^{20} \text{ m}^{-2} \text{ A/W}$	0.47
NB current-drive efficiency, ^a $10^{20} \text{ m}^{-2} \text{ A/W}$	0.25
Overall current drive efficiency, $10^{20} \text{ m}^{-2} \text{ A/W}$ (NB + LH + bootstrap currents)	0.52
Average (peak) electron temperature, keV	19 (25)
Average (peak) ion temperature, keV	20 (26)
Fast-ion mean energy, keV	264
Average (peak) electron density, 10^{20} m^{-3}	1.1(2.1)
Fast ion density, 10^{20} m^{-3}	0.031
Thermal alpha concentration	0.05
Thermal impurity (O^{8+}) concentration	0.024
Effective ion charge (Z_{eff})	2.1
Gross energy confinement time (τ_E), s	0.68

^aThe volume-average density over the outer region of LH current drive is used in evaluating the current-drive efficiency for LH; a volume average over the entire plasma is used for the NB efficiency.

^bAt the 95% flux surface. The corresponding quantities evaluated on the separatrix are $a = 0.855$, $\kappa = 2.74$ and $\delta = 0.77$.

Figure 2. Divertor configurations for TIBER II.

flux of low-energy particles to reduce the sputter erosion of the plates. Identical divertors are located at the top and bottom of the machine to provide the surface area of the plates needed to keep the average heat load below about 2 MW/m^2 .

In steady-state operation, gas must be removed from TIBER at the rate at which fuel is added. The fueling rate for the TIBER baseline is 244 A (equivalent) of DT. For a fusion power of 314 MW, the production rate for helium is 18 A (particle equivalent). The steady-state gas load is therefore about 20 Torr•L/s of D/T and about 3.2 Torr•L/s of helium.

3.2. High-Current-Density, Radiation-Tolerant Magnets

High current density in the toroidal field and central plasma-shaping coils is used to limit the inner radial build. The TF current density must be compatible with the higher neutron flux heating resulting from the use of a thin neutron shield [17]. In the recently developed cable-in-conduit conductors (CICC) where supercritical helium is forced through the conductor, the heat removal capability is large and predictable. In addition, a high degree of transient conductor stability is available at about twice the current density of a typical cryostable system, even with high radiation damage to the copper. Further radiation tolerance of the magnet is achieved by using polyimide insulators that can withstand more than 10^{11} rads of radiation dose.

Figure 3 is representative of the effects of high-energy neutron fluence on Nb_3Sn conductors, both pure and modified [18]. The samples tested were irradiated with 14.1-MeV neutrons at room temperature and then tested at 4.2 K. The lower horizontal axis in Fig. 3 accounts for the "softer" shielded neutron-energy spectrum calculated for TIBER II. Initial increases in current density are explained by the damage giving rise to increased normal-state resistivity of the superconductors, resulting in an increase in the upper critical field. In the modified materials, the upper critical field has already been optimized by additions of a third element to increase the normal-state resistivity. Therefore, any further increase due to damage might be expected to be small. At high fluence a decrease in performance is caused by a drop in critical temperature with increased damage. For unmodified Nb_3Sn , it is clear that the superconductor will retain its performance beyond the

Figure 3. Critical current vs neutron fluence. Data were taken by using a 14.1-MeV neutron source. The lower horizontal axis accounts for the "softer" neutron energy spectrum expected in TIBER II.

10^{19} n/cm² fluence in TIBER II. Modified Nb₃Sn is useful for attaining the high field performance in the central solenoid of the PF system, where the high-field windings are shielded by the TF system and the outer, lower field turns as well as the inboard shielding.

Stability in a CICC is obtained primarily from the enthalpy available in the internal helium over the temperature margin provided by operating the conductor at less than its critical current. As a result, there is no large penalty on stability as the conductivity of the copper is reduced by accumulated damage. However, to protect the TIBER II magnets in the event of a quench, knowledge of copper resistivity vs damage is essential [19].

In a CICC system, the strain in the Nb₃Sn filaments determines critical current at operating conditions. Differential strains in the sheath and the cable result from formation at around 1000 K, cool down to the operating temperature around 4 K, and electromagnetic forces due to charging to full

field and current. An experimental program is underway to examine the conductor/sheath interaction through a series of electromechanical tests [20].

In the first series of tests, the critical currents of model CICC's were measured while the axial load applied to the sheath was increased by small increments until the measured critical current was observed to pass through a maximum. The ratio I_{c1}/I_{cm} of initial or "no-load" critical current to maximum critical current was observed to vary with the fraction of the internal space available for helium in the otherwise identical specimens. These data are plotted in Fig. 4 along with other values for similar conductors inferred from data in the literature [21,22].

3.3. Poloidal Field Coils

Although the baseline TIBER II design is a steady-state machine with non-inductive current-drive, we also investigated, as a contingency, pulsed inductive operation employing whatever limited volt-second capability was available from the poloidal field system. The original PF configuration had separate "pusher" and "puller" coils in the central stack to provide the desired plasma shape. In the final configuration (Fig. 5), all the available space in the core of the machine was filled with a conventional transformer stack. This still allowed high shaping to be attained but increased the total volt-seconds capability to 57.5 V.s. At a plasma current of 8 MA, this provided for a flat-top, ignited burn of ~ 140 s, and for a lower hybrid (LH) assisted burn of ~ 500 s employing LH current-maintenance during transformer reset.

The pusher coil requires a total current of 18 MA, which can be achieved with a coil of 1.1-m inner diam $\times$ 1.9-m outer diam $\times$ 1.13-m height, having an average winding-pack current density of 40 A/mm^2 . To achieve the required average winding-pack current density, conductor grading was necessary as in Table 2. The first grade operates at 14 T. Subsequent grading steps were chosen at 11 T and 8.5 T for convenience (an integral number of layers per grade) and for a margin of improvement in overall current density beyond the previous higher field grade. All conductors have the same generic shape and cable pattern with adjustments made to external dimensions; conduit wall

Figure 4. Ratio of initial (unloaded) critical current of CICC samples to maximum critical current as a function of void or helium fraction inside the conduit.

thickness; wire diameter; and fractions of conductor, copper, and helium to satisfy the requirements of the particular grade. All grades are connected in series and operate at the same current.

3.4. Toroidal Field Coil

The TF system produces 5.55 T at the plasma major radius of 3 m. The space requirements for the plasma, blanket, shield, etc., dictate that the TF coil have an opening of about 6 m on the vertical axis and 4 m on the horizontal axis. These and other electromagnetic requirements can be met with a 16-coil set, an average current density of 40 A/mm², and a maximum field of 12 T, as in Table 3.

Design of both the plasma shaping and TF coil systems will be driven by protection requirements due to high current density and high stored energy.

Figure 5. Configuration with OH coils filling the core.

Table 2. Parameters for the pusher-coil winding pack.

Parameter	Grade 1	Grade 2	Grade 3
Field B (T)	14	11	8.5
Current density J (A/mm ²)	66.4	69.0	76.3
Conductor fraction	0.73	0.64	0.67
Helium fraction	0.27	0.35	0.33
Copper fraction	0.53	0.66	0.77
Operating to critical			
Current, I_{op}/I_c	0.58	0.53	0.55
Bath temperature (K)	4.2	4.2	4.2
Current sharing temperature (K)	6.4	7.5	8.0
Stability parameter	0.09	0.16	0.14

The TF system, because of its nearness to the plasma, reduced shielding, and compactness also accommodates the effects of nuclear heating, neutron fluence to the superconductor of the order of 10^{19} n/cm², damage to the copper stabilizer on the order of 0.006 dpa, and radiation exposure to insulating materials in the winding pack on the order of 10^{10} rad.

3.5. Design for Maintenance

All internal components of TIBER II can be removed or repaired by use of remote tooling employing simple straight-line motions assisted by specialized jigs and fixtures. Identified short-lifetime parts, such as the divertors and test modules, are installed in locations where they can be removed with minimum disturbance to the rest of the system. (Refer to Fig. 6 in the following discussion.)

The assembly sequence is started by mounting the cryogenic magnet base-plate along with the warm support pedestals which penetrate it. The outer poloidal field coils can next be assembled on the PF coil jacking system, which permits their withdrawal into a retracted position in an annular basement. This "parking" location for the outer PF coils is necessary during assembly or removal and replacement of a TF coil.

Table 3. Parameters for the TF coil winding pack.

Total current (MA)	5.2
Winding-pack cross section (m^2)	0.130
Average winding-pack current density ($A\ mm^{-2}$)	40
Number of turns	144
Conductor current (kA)	36
Effective area of a turn (mm^2)	904
Conductor + insulation (mm)	33.6 x 26.9
Overall winding-pack material fractions	
Fraction of steel	0.37
Fraction of insulation	0.6
Fraction of conductor	0.34
Fraction of helium	0.23
Conductor cable-space cross section (mm^2)	515
Fraction of conductor in the cable space	0.60
Fraction of copper in the conductor	0.60
Number of cable strands	375
Strand diameter (mm)	1.02
Coil dump voltage (kV)	10
Estimated maximum hot-spot temperature (K)	130
Estimated maximum quench pressure (MPa)	40

Sixteen TF coil subassemblies are prepared one at a time in a staging area in the reactor room. Each consists of one TF coil, the upper and lower support shields, one inboard shield segment, one outboard shield segment, and the overturning wedge members which, when keyed, resist the large overturning forces on the outboard legs of the TF coils. These 250 tonne subassemblies are lowered onto the baseplate and are jacked toward the center to make the necessary close contact at the noses of the TF coils. Then shear pins are installed radially outward from the cylindrical space not yet occupied by the OH coils. The inboard legs of the TF coils are thus pinned to transmit shear thereby resisting the large overturning loads. Shear keys are next installed

Figure 6. TIBER II assembly sequence.

between the wedges preassembled to the outboard legs of the TF coils. All sixteen TF coil subassemblies thus act in concert to resist the outboard leg overturning loads.

The OH coil sub-assemblies can next be lowered in place by the overhead crane. There are four solenoids in each of four sub-assemblies.

The outer PF coil assemblage of four coils can then be raised from its parked location and joined to the TF coil outboard case structure where it hangs after the jacks are retracted.

All of the above activities can occur with or without the main outer vacuum shell in place. During initial assembly it may be desirable to build the outer vacuum wall after the TF coil installation.

After the main vacuum shell is in place the divertors can be installed through the service doors (16) in the vacuum shell. These move at a 29.5 degree angle from the horizontal and are guided into place by permanent linear bearings. Nominal clearance of one-half centimeter is provided between adjacent divertors and adjacent blanket test modules.

Once the upper and lower divertor are in place between any pair of TF coils, the blanket test module can be installed between them by simple horizontal straight line motion.

The blanket test modules will be the most frequently replaced components, but their removal least disturbs the system because major components do not need to be moved to provide access. Warming of the magnets and opening the vacuum vessel is all that is necessary to remove these test modules. In some cases re-entrant spaces in the test region or air-locks will allow replacement of short term test articles without warming or admitting the atmosphere. Should repair of a divertor be necessary they can be accessed by first removing the one blanket test module above or below that divertor segment.

3.6. Structural Details

Large forces on the coils, coupled with the requirement for simple assembly/disassembly procedures, place constraints on the structural design. The centering force on the TF coils generates an effective pressure of 97 MPa on the central solenoid. Vertical, separating loads in the TF coils are mutually shared by the conductor sheath and the coil cases. The out-of-plane loads on the TF coils are supported by the coil cases acting as a cylinder loaded in shear. These cases are connected together with a series of removable

match-machined pins, thereby providing the combined TF coil legs with the ability to resist shear loading. A preliminary finite-element stress-analysis calculation indicates that this concept will be satisfactory for TIBER II. Additional forces on the TF coil are carried by the coil case, an intercoil structure at the top and bottom of the TF coils, and the structure connecting the large poloidal field (PF) coils. Out-of-plane forces on the outside of the TF coils are transferred to the large PF coil structure, which includes the coils, their cases, and a structural cylinder that separates the coils.

The entire cold structure (consisting of the coils, base plate, center post, and large PF-coil structure) weighs approximately 1000 tonnes. It is supported by 16 plate columns, 1 m high, that provide thermal isolation between the vacuum vessel and the cold magnet structure. These plate columns are thinned at the ends so that they have an elastic hinging action to allow for thermal contraction of the coil structure while still providing position control of the magnet set.

All of the active components of the machine are housed in a common vessel that is 13 m in diam and 14 m in height. This vessel provides the vacuum environment for the plasma and acts as a dewar for the superconducting magnets. To reduce the cryogenic heat load, the liquid-helium-cooled components are surrounded by liquid-nitrogen-cooled baffles. The magnets are all supported by a base plate and center post, both of which are cooled with liquid helium to simplify the cold-to-warm structure interfaces. Shielding and other internal components are supported by a separate warm structure attached to the same foundation as the cold structure.

3.7. Neutron Shielding and Tritium Breeding

Because the radial dimension of the center-post region of the tokamak determines the overall size of the machine, the inboard neutron shield is minimized. About 0.5 m of high-performance shield with tungsten, steel, and water limits the nuclear heating in the windings of the toroidal field coils to < 72 kw. This heat is removed by flowing helium. On the outboard section of the torus, the shielding thickness can be arbitrarily large, so that tritium breeding can be accomplished. Also, blanket test modules can be placed at the outboard equatorial plane. The remainder of the shield is a cold, low-technology breeding region, such as steel with a lithium salt in the

cooling water. With a beryllium neutron multiplier, it is possible to be tritium self-sufficient.

With a fusion power near 300 MW and a modest duty factor, it may be unnecessary to breed tritium in the ETR. Full performance of 30% availability in Table 4 would require about 5 kg/year of tritium. This amount is more than is available through Canada, but the supply might be supplemented by other sources. Lowering the duty factor to 10% would nearly match the Canadian source and still permit a meaningful test program. The decision to breed tritium in the ETR, the fluence, and the availability are uncertain at present. However, a small machine permits the flexibility of a choice of supply.

4. COMPARISON OF NATIONAL ETR STUDIES

It is understood that an international design effort will establish a common design. However, it is informative to that process to review and compare current efforts. At present there are five separate national or international conceptual-design studies in progress, as summarized in Table 5. Of these, TIBER II has the smallest radius but comparable or better performance. However, it is not ignited in the usual sense, but has a sustained steady-state burn because of current drive. The experimental basis for current drive will need to expand in the coming years to justify such a choice. Alternatively, the use of a backup ohmic heating drive in TIBER II was investigated for current ramp-up or burn.

Table 4. Projected TIBER II operation schedule.

	Duration (yr)	Emphasis	Availability (%)	Annual T_2 consumption (kg/yr)	Integrated 14-MeV fluence; peak/average (MW yr/m ²)
Phase I	1	H ₂ checkout	15	--	--
Phase II	2	DT operation	10	1.5	0.36/0.24
Phase III	10	Engineering testing	30	5.2	5.4/3.6

Table 5. The National ETR studies.

	INTOR	NET	FER	TIBER	OTR
<u>Machine Parameters</u>					
R(m)	5.0	5.18	4.42	3.0	6.3
a(m)	1.2	1.35	1.25	0.83	1.5
elongation	1.6	2.05	1.7	2.4	1.5
triangularity	0.25	0.48	0.2	0.4	0.3
I(MA)	8.	10.8	8.74	10.	8.
B(T)	5.5	5.	4.61	5.55	5.8
<u>Operating parameters</u>					
$\langle T \rangle$ (keV)	10.	10.	12.	19.2	8.
$\langle n \rangle$ ($10^{20}/\text{cm}^3$)	1.6	1.7	1.14	1.06	1.7
Murakami parameter	19.35	23.42	14.54	7.62	24.56
Greenwald density	1.77	1.88	1.78	4.62	1.13
<u>Confinement parameters</u>					
confinement time required for operation (sec)	1.4	1.9	1.7	0.68	1.7
aB^*	36.3	33.8	26.6	25.6	50.5
$t_{\text{neo-alcator}}$	7.7	11.8	4.6	1.54	19.
$t_{\text{ASDEX-H}}$	4.0	5.6	3.9	3.0	5.
$t_{\text{Kaye-Goldston}}$	0.90	1.3	0.93	0.56	1.03
$t_{\text{Kaye-Goldston-H}}$	1.77	2.5	1.75	1.02	2.1
<u>MHD limits</u>					
beta limit	4.90	5.6	5.3	6.	3.2
I/aB	1.21	1.6	1.52	2.17	.92
Equivalent Troyon coeff.	4.04	3.5	3.49	2.76	3.48
<u>Impurity control method</u>					
divertor	SN	DN	SN	DN	SN
<u>Operating characteristics</u>					
Fusion power (MW)	585	650	406	314	500
Burn time (sec)	150	350	800	c.w.	600
No. of pulses	200,000	80,000	18,000	N.A.	300,000
Heating	ICRF	TBD	ICRF	LH+NBI	ICRH
Aux. power (MW)	50	50	40	63	50
Q	ignited	ignited	>20-30	≥ 5	> 5

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1.0 INTRODUCTION TO THE TIBER II/ENGINEERING TEST REACTOR

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1.0 INTRODUCTION TO THE TIBER II/ENGINEERING TEST REACTOR

1.1 BACKGROUND

The Tokamak Ignition/Burn Experimental Reactor (TIBER) has been the U.S. option for the international ETR. The concept evolved from earlier work on TFCX to develop a small, ignited tokamak. While the copper-coil versions of TFCX became the short-pulsed, 1.23-m radius, Compact Ignition Tokamak (CIT), the superconducting TIBER with long pulse or steady-state and a 2.6-m radius became a candidate for international collaboration. Recently, the design has been updated to TIBER II which can accommodate the conservative Kaye-Goldston confinement scaling, double-poloidal divertors for impurity control, steady-state current drive, and nuclear testing (see Fig. 1-1).

LEGEND: (1) Lower Support Shield; (2) Upper Support Shield; (3) Inboard Shield; (4) Outboard Shield; (5) Toroidal Field Coils; (6) Ohmic Heating Coils; (7) OH Support Structure; (8) Poloidal Field Coils; (9) PF Support Structure; (10) Cryogenic Support Plate; (11) Cryogenic Support Struts; (12) Shield Support Pedestal; (13) Shield Stabilizer; (14) Blanket Test Module; (15) Lower Divertor; (16) Upper Divertor; (17) Divertor Vacuum Ducts; (18) Outer Vacuum Wall; (19) Vacuum Dome.

Figure 1-1. TIBER II Engineering Test Reactor.

1.2 GENERAL OBJECTIVES: STEADY STATE, REDUCED COST

TIBER is a candidate for an international Engineering Test Reactor (ETR), which is to be the next major step in the tokamak program between the present generation of large tokamaks (TFTR, JET, JT-60, T-15) and the generation of demonstration reactors (DEMO). Anticipating that a reliable DEMO would need steady-state operation, and that relevant blanket testing in the ETR would therefore need steady-state conditions, TIBER is designed for steady-state operation. To make the ETR affordable to the world fusion program, we seek a reduced cost design through the use of selected innovative engineering design features. By reduced cost, we mean something between NET and CIT, and comparable to TFCX, about 1 billion current U.S. dollars in direct capital cost. In achieving these general objectives, the design should incorporate advances in physics and engineering anticipated over the next six years.

1.2.1 TIBER MISSION (PROGRAMMATIC OBJECTIVES)

- Demonstrate reliable steady-state operation of an integrated fusion device.
- Complete the physics database for long-pulse burn.
- Demonstrate the performance of nuclear components at reactor-relevant conditions.

TIBER must be designed to test reliability, safety, and availability of integrated fusion reactor components, at reactor level particle and heat fluxes. Reliability is to be a key part of the design, accomplished through specific approaches, including (a) reduction of cyclic stresses through maximum use of steady state, (b) use of non-inductive current drive for current profile control and disruption stabilization, and (c) high-fluence operation at nominal plasma and coil parameters (β , current I_p , poloidal field-coil field B_{PF} , and stress σ_{PF}) which are sufficiently below the maximum design values, such that plasma disruptions are significantly reduced and PF coil life is significantly enhanced.

TIBER must be designed to operate in a full steady-state mode, such that the neutron flux is essentially constant for single test runs of substantial duration (≈ 1 week) which are not design-limited by any component, except for

long-term limits set by particle erosion and neutron damage. A sufficient number of on/off cycles must be produced, such that there is reasonable time for engineering and physics tests to meet the programmatic objectives. In addition to steady-state operation, the same machine design must be capable of initial (first phase) operation using inductive current-drive assisted with lower hybrid (lower-hybrid current drive used for combinations of: current ramp up, edge-current drive during inductively-driven flat top at high Q, and transformer reset). In both initial pulsed operation and steady-state operation, driven, high Q operation is the primary objective, not pure self-ignition. Pure ignition capability of TIBER is not necessarily excluded, but because of the objective of achieving a steady-state, reliable, controlled burn, high Q operation, rather than ignition, is chosen to determine the design requirements.

By high Q, we mean $P_{\text{fusion}}/P_{\text{CD}} \gtrsim 5$ where P_{CD} is the absorbed current drive power. At this minimum Q, the DT alphas still play an important role in heating, while at the same time there is sufficient auxiliary current drive power allowed to drive the current, control the burn, and control the current profile (sawteeth, disruption stabilization). Also, achievement of $Q \approx 5$ in TIBER at $P_{\text{fusion}} \approx 300$ MW, would be a significant implication for higher Q, steady-state reactors, since fusion power is expected to scale up faster than current drive power in larger future reactors.

The purpose of the nuclear testing mission of the ETR is to demonstrate the performance of nuclear components at reactor-relevant conditions. "Nuclear Components" refer to first wall/blankets, shielding (bulk and penetration shields), tritium processing systems, nuclear aspects of high-heat-flux components (lifetime and energy extraction), and associated assembly/maintenance equipment. "Performance" has three key aspects:

1. Demonstrate that the fusion fuel cycle can be made self-sufficient.
2. Demonstrate that useful fusion energy can be extracted and used safely.
3. Demonstrate required radiation protection for reactor equipment and personnel.

It is anticipated that the ETR would be followed by some type of fusion demonstration power reactor. Thus, the testing conditions in the ETR should extrapolate to anticipated demonstration reactor conditions. This requires that the ETR be designed to achieve certain values of key parameters important to nuclear testing (e.g., average neutron first-wall loadings of $\sim 1 \text{ MW/m}^2$ and long-pulse or steady-state operation with high availability over extended testing periods of weeks or months). The testing mission will also deal with the materials and safety issues of nuclear components.

"Demonstrate" is taken to mean to show that certain performance levels can be achieved. It does not necessarily mean that a complete, integrated reactor system must be operated; but rather one must show that its individual subsystems can perform at the necessary conditions. For example, it is not necessary to actually produce electrical power but it is necessary to show that fusion energy can be extracted from a first wall/blanket test module at suitable temperatures that could result in useful power production. As another example, it may not be necessary for the ETR to be self-sufficient in tritium, but tests in the ETR should show that the proper conditions (high tritium burnup in the plasma, high local tritium breeding in the blanket, minimum tritium losses in the processing system, etc.) can be achieved so that tritium self-sufficiency could be attained.

In general, TIBER operation and the nuclear test components must meet the following probable and potential safety requirements:

Probable Legal Requirements:

1. The public will be protected against the consequences of any credible accident such that the maximum off-site 2-hour dose is 250 mSv (25 rem). (Could take credit for active systems and confinement building.)
2. The maximum radiological exposure to plant workers must be below 50 mSv/yr (5 rem/yr).
3. Public individual exposure from routine atmospheric effluents must be below 250 $\mu\text{Sv/yr}$ (25 mrem/yr).

4. Public individual exposure from routine surface water effluents must be below 40 $\mu\text{Sv/yr}$ (4 mrem/yr).
5. There must be a plan for handling waste and for decommissioning the facility.

Potential Requirements:

1. Individual members of the public will be provided a level of protection against the consequences of fusion facility operation such that no individual bears significant additional risk (0.1%) to life or health.
2. Plant workers will be provided a level of protection while at work such that workers will be safer at work than away from work. (Worker risk is below 10% of background.)
3. The facility will be protected against off-normal events such that there is only a small chance (1%) that off-normal events require abandoning the facility and its research mission.

1.2.2 TIBER TECHNICAL OBJECTIVES

1.2.2.1 Steady-State Operation

1. Steady-state DT burn with equilibrium impurity and alpha ash removal (including full steady-state operation of a continuous period of at least 1 week in the later stages of operation)
2. $Q \geq 5$
3. Average neutron wall loading $\langle \Gamma \rangle \geq 1 \text{ MW/m}^2$
4. Minimum plasma current $i_p = 10 \text{ MA}$
5. Toroidal beta $\langle \beta_T \rangle = 2.8 I_p / a B_t$
6. Safety factor $q_{95} \geq 3.0$ calculated at the 95% flux surface.
7. Availability $\geq 30\%$ (during nuclear testing phase III)

1.2.2.2 Inductive Operation

The baseline requirement for TIBER is a steady-state, current-driven machine. However, as a contingency, a modest inductive volt-seconds capability will be provided. The purpose is to permit an early decision to build the ITER design, before the final selection of steady-state current drive systems. Thus, besides being capable of full non-inductive current-driven steady-state operation, TIBER must demonstrate the following inductive capabilities in the same machine (i.e., same magnet set and reactor configuration):

1. A volt-seconds capability sufficient to inductively ramp the plasma current up to at least 70% of the full steady-state value, i.e. to $I_p = 7\text{MA}$.
2. A volt-seconds capability which is sufficient to sustain a limited flat top burn (~ 100 's seconds) and permit accumulation of a limited neutron fluence ($\geq 1 \text{ MW}\cdot\text{yr}/\text{m}^2$), using the assistance of lower hybrid current drive for initial current ramp up, edge current drive during the burn, and transformer recharge.

Ignition performance should be investigated, but is not an a priori requirement for TIBER operation.

1.2.2.3 Reactor-Relevant Technologies

1. Superconducting toroidal and poloidal coils
2. Noninductive current drive
3. Plasma composition control (e.g., divertor)
4. Plasma power balance control
5. Plasma heating and fueling
6. Blanket heat removal and tritium production
7. Tritium fuel cycle
8. Remote maintenance
9. Vacuum
10. Instrumentation and Controls

1.2.2.4 Engineering Objectives

TIBER is to be designed for reduced cost through selective innovative engineering design features. The components should be designed with advanced technology which is capable of verification during the Development and Demonstration phase before construction. For example, high current density magnets are to be demonstrated as part of the multipurpose coil program, and free electron masers are part of on-going research programs at LLNL and TRW. High heat flux components such as divertor plates must be consistent with experiments to be performed on the PICES facility at UCLA, or the Sandia National Laboratory electron beam facility.

At the beginning of the TIBER study, the approach to reduced cost was taken to be a set of innovative design features (discussed in Drivers of the TIBER design, elsewhere in this report) with the aim of minimizing the major radius of the machine. Preliminary designs for TIBER indicated the minimum major radius was approximately $R_0 = 3$ m. Instead of designing TIBER to a specific minimum dollar cost (which would be very difficult), the major radius was set at $R_0 = 3$ m for the remainder of the study, as an equivalent means of keeping the cost down.

1.2.2.5 Shielding and Tritium Breeding Design Objectives

The primary shielding design objective for TIBER is for the shielding to limit SC coil heating and end-of-life damage to allow $\sim 10^8$ burn-seconds while having minimal impact on the machine's cost and size. The shield system is to be low risk and designed to be a lifetime component.

The primary tritium breeding objective is to test tritium breeding and extraction under reactor relevant conditions in individual blanket test modules.

Because of a potential tritium supply shortfall and tritium's possible high cost (~ 45 M\$/year for operation at 30% availability and a tritium unit cost of 10 M\$/kg), an important secondary objective of the TIBER study is to explore TIBER's tritium self-sufficiency potential from the standpoints of practicality and cost effectiveness. This secondary objective involves modifying the shield to breed as much tritium as possible without compromising shielding effectiveness or increasing the machine size. The modification explored is to add a soluble lithium compound to the shield/first wall cooling water and, where space allows, replace shield material with Be for neutron

multiplication. Whether or not baseload tritium breeding is actually performed on TIBER depends on the availability and unit price of the external tritium supply. Cost-benefit studies would then show the relative desirability of in-situ breeding at given fractions of full self-sufficiency. Our current rationale is to at least demonstrate that TIBER can breed tritium at full self-sufficiency.

The purpose of this secondary tritium breeding objective is to explore alternative ways to supply the requisite TIBER tritium needs to carry out the nuclear test mission (2.3 above). This approach of tritium breeding in the shields is not required to be reactor-relevant (i.e., at high temperatures) as is required in the blanket test modules of the nuclear test program.

1.2.2.6 Maintenance Objectives

1. All systems in the reactor vault designed for fully remote maintenance.
2. First wall and divertors designed for frequent replacement.
3. All reactor components, including any magnet designed to be replaced in less than 6 months.
4. Magnets designed for rapid warm-up and cool-down for each maintenance operation.

1.2.3 TIBER OPERATIONAL PHASES

Within a total machine lifetime of 12 years (2 years checkout + 10 years operation), TIBER is expected to increase pulse length, availability, and fluence in successive phases. These phases are taken to be:

- I. Checkout Phase - 1 year hydrogen/deuterium checkout followed by 1 year DT checkout, both at low machine availability (~ 5-10%).
- II. Physics Mission Phase - 3 year DT burning for physics missions at 10% availability.

III. Nuclear Testing Mission Phase - 7 year DT burning phase for nuclear testing mission at 30% availability. Capability for true steady-state (≥ 1 week) is to be demonstrated.

A minimum peak fluence goal at the end of Phase III is 3 MW yr/m^2 (which translates to a full power operating life of $\sim 4.7 \times 10^7$ burn-seconds at TIBER's peak wall-loading of 2 MW m^{-2}). The TIBER machine is to be designed for a minimum of 50,000 on/off plasma cycles and a total life of 7.9×10^7 DT burn sec and a peak fluence of 4.8 MW yr/m^2 . Any inductive test pulses in Phases I and II are to be included in this limit.

1.3 MACHINE PARAMETERS

Parameter	Description	Value
R_0 (m)	plasma major radius	3.0
a (m)	plasma minor radius ^a	0.834
κ	elongation ^a	2.4
δ	triangularity ^a	0.4
V_p (m ³)	plasma volume ^a	96
S_p (m ²)	plasma surface area ^a	170
B_0 (T)	toroidal field at R_0	5.55
$\Delta\phi$ (Vs)	volt-sec capability	57.5
P_{fus} (MW)	fusion power	314 (314) ^b
Q	fusion power/input power	5.0 (∞ , IM ~1.1) ^b
$\langle\Gamma\rangle; \Gamma$ (MW/m ²)	average (peak) wall loading	1.3; 2.0 (1.3; 2.0) ^b
P_{NB} (MW)	neutral beam power	44
P_{LH} (MW)	lower hybrid power	19
τ_{burn} (s)	burn pulse length	steady-state (141) ^b
ϕt (MWyr/m ²)	end-of-life fluence	5 (0.23-1.25) ^b
I_p (MA)	total plasma current	10.0 (8) ^b
I_{bs} (MA)	bootstrap current	2.7
I_{LH} (MA)	lower hybrid driven current	3.9
I_{NB} (MA)	neutral beam driven current	3.4
γ_{LH} (10 ²⁰ m ⁻² A/W)	lower hybrid efficiency	0.47
γ_{NB} (10 ²⁰ m ⁻² A/W)	neutral beam efficiency	0.25
γ_{CD} (10 ²⁰ m ⁻² A/W)	overall current drive efficiency (NBI + LH + bootstrap currents)	0.52
$\langle T_e \rangle; T_e$ (keV)	average (peak) electron temp.	19; 25 (10 av.) ^b
$\langle T_i \rangle; T_i$ (keV)	average (peak) ion temp.	20; 26 (10 av.) ^b
$\langle n_e \rangle; n_e$ (10 ²⁰ m ⁻³)	average (peak) electron density	1.1; 2.1 (1.78 av.) ^b
$n_{Greenwald}$ (10 ²⁰ m ⁻³)	Greenwald density limit	2.57
$n_{Murakami}$ (10 ²⁰ m ⁻³)	Murakami density limit	1.35
C_α	thermal alpha concentration	0.05
C_{oxygen}	thermal impurity concentration	0.024
Z_{eff}	plasma effective charge	2.1
τ_E (s)	gross energy confinement time	0.68

a. Values on the 95% flux surface.

b. Values in parentheses are for inductive ignited operation.

Power Balance Parameters

<u>Parameter</u>	<u>Description</u>	<u>Value</u>
P_{fus} (MW)	total fusion power	314
$P_{fus,beam}$ (MW)	beam-bulk fusion power	44
P_{α} (MW)	alpha power	63
$P_{\alpha,e}; P_{\alpha,i}$ (MW)	alpha power to electrons (ions)	42; 21
P_{LH} (MW)	lower hybrid power	19
P_{NB} (MW)	neutral beam power	44
$P_{NB,e}; P_{NB,i}$ (MW)	neutral beam power to electrons (ions)	12; 32
P_{brem} (MW)	bremsstrahlung power	7.3
P_{sync} (MW)	synchrotron power lost*	2.0
$P_{c,e}$ (MW)	electron conduction power lost	63
$P_{c,i}$ (MW)	ion conduction power lost	52
$P_{i,e}$ (MW)	ion-electron drag power	0.91
τ_E (s)	gross energy confinement time	0.68
$H = \tau_E / \tau_E$ (L-mode)	H-mode energy confinement enhancement	1.14

*Assuming 90% wall reflectivity

Baseline Current Drive Parameters

<u>Parameter</u>	<u>Description</u>	<u>Value</u>
P_{NB} (MW)	neutral beam power	44
P_{LH} (MW)	lower hybrid power	19
I_{NB} (MA)	neutral beam driven current	3.4
I_{LH} (MA)	lower hybrid driven current	3.9
I_{BS} (MA)	bootstrap current	2.7
$N_{ , LH}$	lower hybrid parallel refractive index	1.60
$N_{ , ac}$	accessible parallel refractive index (high field side)	1.55
$E_{ }/\langle T_e \rangle$	normalized tail parallel energy	8.0
r_{max} (m)	radius of maximum penetration (LH)	0.45
$\langle T_e \rangle_{LH}$ (keV)	average electron temperature (LH)	5.5
$\langle n_e \rangle_{LH}$ ($10^{20} m^{-3}$)	average electron density (LH)	0.75
$\omega_{pe}^2 / \omega_{ce}^2 _{r_{max}}$	dielectric constant at r_{max}	0.33
E_{crit} (keV)	hot deuterium critical energy	393
n_{hot} ($10^{20} m^{-3}$)	hot ion density	0.031
E_{hot} (keV)	hot ion average energy	264

1.4 SPECIAL FEATURES

The TIBER design incorporates physics requirements and neutron wall-loading and fluence parameters that will make it compatible with a nuclear testing mission. Reactor-relevant physics will be tested by using high-beta, current drive and steady-state operation. The base case uses a combination of lower-hybrid and radio-frequency power and high-energy neutral beams for current drive. Minimum neutron shielding, compact structures, high-magnet-current densities, and remotely maintainable vacuum seals all contribute to the compact size. Many of the special features of the current design effort are addressed in the subsequent sections of this report. These and other TIBER II features are briefly summarized in the following paragraphs.

1.4.1 PHYSICS BASIS

Kaye-Goldston scaling of electron confinement time is assumed with a moderate H-mode enhancement factor (given the double-null poloidal divertors in TIBER II) of 1.1. The energy confinement time of electrons and ions is taken to be equal. In addition, we conservatively assume that both the alpha power P_α and current drive power degrade confinement according to the Kaye-Goldston scaling law. The operating beta is taken as a conservative 2.8 I_p/aB_t limit, and the average plasma density is below the Murakami density limit.

1.4.2 CURRENT DRIVE

Several current-drive options have been considered for TIBER, including lower hybrid (LH), electron-cyclotron resonance (ECR), negative-ion neutral beams, and various combinations. It appears that TIBER would require about 60 MW of power to drive the nominal 10 MA of plasma current. The current-drive system is also intended to provide the necessary heating for startup, which is needed in any pulsed or steady-state tokamak. Our baseline is a combination of NB and LH and we note that recent data from TFTR has shown promise for neutral-beam drive if energies are raised to 500 keV and above. Bootstrap current might double projected current-drive efficiencies, which would decrease ETR equipment requirements or, conversely, lead to efficient reactors.

1.4.3 COMPACT DIVERTOR

The TIBER edge plasma and compact divertor operate in a high-recycling mode. This mode generates the high density needed to allow the edge plasma to shield the core plasma from the gas and impurities generated at the divertor plates. It also allows the charged-particle power to be carried out by a high flux of low-energy particles to reduce the sputter erosion of the plates. Identical divertors are located at the top and bottom of the machine to provide the surface area of the plates needed to keep the average heat load below about 2 MW/m^2 .

1.4.4 NEUTRON SHIELDING

Because the radial dimensions of the center-post region of the tokamak determines the overall size of the machine, the inboard neutron shield is minimized. About 0.5 m of high-performance shield with tungsten, steel, and water limits the nuclear heating in the toroidal field coils to $<72 \text{ kW}$ and allows for 2.5 fullpower years (4.8 MWY/m^2 peak). On the outboard section of the torus, the shielding space increases to 1.2 meters so tungsten is not required. Blanket test modules are placed at the outboard equatorial plane. The shield is cooled with low temperature water ($<100^\circ\text{C}$). By adding a lithium salt to the cooling water and replacing some of the outboard steel shielding with beryllium steel shielding tritium self-sufficiency is probable.

1.4.5 TRITIUM SUPPLY

Full performance of 30% availability would require about 5 kg/year of tritium. This amount is more than is available through Canada, but the supply might be supplemented by other sources. Lowering the duty factor to 10% would nearly match the Canadian source and still permit a meaningful test program. The decision to breed tritium in the ETR, the fluence, and the availability are uncertain at present. However, a small machine permits the flexibility of a choice of supply.

1.4.6 HIGH-CURRENT-DENSITY, RADIATION-TOLERANT MAGNETS

High current density in the toroidal field and central plasma-shaping coils is desirable to limit the dimensions of the center post. This requirement must also be compatible with the higher neutron fluence resulting from a thin shield. Fortunately, we can take advantage of the recent

development of Cable In Conduit Conductors (CICC). Because supercritical helium is forced through the conductor, heat removal is large and predictable. In addition, electromagnetic stability is more dependent on the heat capacity of helium than on the resistivity of copper, as with cryostability, permitting about twice the current density. Because the neutron-damage-induced resistivity rise in the copper is of second-order importance, annealing of copper is unnecessary for stability. Further radiation tolerance is achieved by using polyimide insulators that can withstand more than 10^{10} rads of neutron damage. Because the superconductor neutron damage limit for niobium-tin conductors is 2×10^{19} neutrons/cm², the ultimate life of the magnet is 10^8 s for the selected thin shield. Other more radiation-tolerant conductors, such as niobium-titanium, are available, but not at high-field. Newer, better conductors are being explored, but they will not be available in the time frame of the ETR.

1.4.7 COSTS

By minimizing the TIBER/ETR size, direct construction costs can be kept reasonable. Other, larger designs would have considerably higher costs. Most of the tokamak components, such as magnets, first-wall structures, and vacuum vessels vary with surface or the square of the major radius. Some components, such as plasma-heating systems, are related to volume and vary as the cube of the major radius. Fortunately, facilities, diagnostics, and other auxiliary systems do not change as much, such that the net change is approximately linear with the radius. However, most pulsed tokamaks with 100-V·s transformers have a major radius greater than 5 m, with an aggregate cost that increases linearly with the major radius, or about double that of TIBER II.

1.5 COST ESTIMATE

1.5.1 ESTIMATE BASIS

The cost estimate for TIBER II is based on the conceptual design described in this report, and it was developed concurrently with the design documentation. The technical basis for the estimate varies correspondingly with the level of design specification and detail provided for each system. TIBER II is assumed to be a government sponsored project on an undeveloped site, with plant systems and components procured through industrial subcontracts. First-of-a-kind design and manufacturing costs are estimated for components that are not commercially available, but the costs of the research and development programs to provide the design basis are not included. Also specifically excluded are the costs of the nuclear test modules and their supporting systems, although general plant services and operations space are provided.

The estimate includes direct costs, indirect costs, contingency and operating costs. The distinction between direct and indirect costs to the project is determined by the management structure, specifically the arrangements for plant engineering, procurement, construction and construction management. It is assumed that these services are provided by the managing organization or by subcontractors with a total project scope, so that the costs are recorded as indirect, rather than as direct costs to specific components and systems.

Direct cost consists of equipment and materials, installation labor and component engineering. Equipment cost is the purchase price for the component delivered to the site. The equivalent purchase price is estimated as the build-to-print cost if the component is not commercially available. The component engineering cost, or cost of generating the design for a build-to-print contract, is also estimated for first-of-a-kind equipment. This accounts for the non-recurring design costs, but not for the systems design for the specific application, which is included in indirect cost. The direct costs accounts follow the definitions and code of accounts described in the cost accounting system of Ref. [1].

The cost data base used depends on the specific account. Some of the equipment was designed in detail and original estimates were developed by potential manufacturers. Other costs are scaled from the historical base or

from estimates for similar systems developed by other design studies. All costs are escalated and reported in 1987 price levels.

1.5.2 DIRECT COSTS

The direct cost estimate is summarized in Table 1.5-1, and the data base used to develop the estimate is summarized by major account in the following.

1.5.2.1 Buildings Cost

The cost of the buildings is scaled with the external volume to the 0.8 power. Scaling factors were developed from the historical costs for similar buildings reported in the Energy Economic Data Base (EEDB, Ref. [2]) and from costs estimated for the TFCX project in 1984 (TFCX, Ref. [3]). There is a reasonable cost correlation between buildings of similar design requirements. More detailed building layout, internal design, service specification and structural design are required to provide the basis for a cost estimate based on materials, equipment and labor. This more accurate estimate is justified for the reactor and reactor maintenance buildings, due to their large cost and unique requirements.

1.5.2.2 Reactor Cost

The reactor accounts include the first wall and divertor, the blanket and shield, and the structural support. The first wall and divertor costs were scaled by area from TFCX.

The shield cost was developed based on a detailed design by a potential manufacturer. The materials costs are current prices for the steel, tungsten and beryllium quantities required. A manufacturing cost of 65 M\$ was estimated for the complex design, and the component engineering cost was 25% of the materials and manufacturing cost. This cost estimate is considerably greater than that used in other studies. In particular, the manufacturing cost for simple steel shields is generally estimated at about 20 \$/kg (e.g. TFCX). However, the present cost is consistent with those estimated for high performance (thermal and tritium breeding) blankets and with those for high performance heat exchangers. Further iteration of the design is clearly justified to simplify and better understand the manufacturing requirements.

The cost of the tokamak support structures was estimated at a uniform material and fabrication cost of 28\$/kg, with installation and engineering

Table 1.5-1. TIBER II cost summary.

Account	Description	Cost (millions of dollars)	
21	STRUCTURES	165	
211	Site Improvements	12.0	
212	Reactor Building	37.4	
214	Maintenance Building	50.0	
215	Tritium Building	26.1	
216	Electrical Building	8.0	
217	Other Buildings	31.3	
22	REACTOR EQUIPMENT	1218	
221	Reactor System	255.7	
221.1	First Wall		18.0
221.3	Shield		160.6
221.4	Structure		27.6
221.5	Divertor		49.5
222	Magnets	228.1	
222.1	Toroidal Field Coils		106.1
221.2	Outer Poloidal Coils		58.2
221.3	Inner Poloidal Coils		31.4
221.4	Solenoid		29.4
223	Power Injection	270.0	
223.1	ECRH System		30.0
223.2	Lower Hybrid System		60.0
223.3	Neutral Beam System		180.0
224	Vacuum System	25.7	
225	Power Conditioning	100.4	
225.1	TF Coil Power Supply		13.8
225.2	PF Coil Power Supply		42.9
225.3	Energy Storage		43.7
226	Heat Transport	64.9	
226.1	Reactor Cooling		14.8
226.2	Cryogenic System		34.0
226.3	Auxiliary Cooling		16.1
227	Fuel Handling	103.2	
227.1	Fuel Injection		37.1
227.2	Fuel Processing		7.3
227.3	Fuel Storage		6.7
227.4	Atmospheric Recovery		27.1
227.5	Water Recovery		25.0
228	Instrumentation	100.0	
229	Maintenance Equipment	70.0	
24	ELECTRIC PLANT	46	
25	MISCELLANEOUS PLANT	26	
26	HEAT REJECTION SYSTEM	15	
TOTAL DIRECT COST		1469	
INDIRECT COST		510	
CONTINGENCY		451	
CONSTRUCTED COST		2430	

costs of 10% each, based on the historical data. A more accurate estimate can be made by obtaining current fabricator estimates for the specific structural elements.

1.5.2.3 Magnet Costs

The magnet cost estimate is based on average scaling factors developed for the TFCX project: 3 \$/kA.m.T for the niobium-tin conductor, 120 \$/kg of winding pack for fabrication, and 75 \$/kg for the case. Installation is estimated at 5% and component engineering at 10%. The fabrication cost of the TIBER II magnets should be estimated independently by a manufacturer, since there is considerable variation in the data base, the design is unique, and the cost leverage is large.

1.5.2.4 Power Injection Systems

The cost estimates for the ECRH, lower hybrid and neutral beam systems are based on conceptual designs and associated equipment and materials lists. A cost factor of 3 \$/W is used for the free electron maser based on a reference design by a potential supplier, and the same factor is used for the lower hybrid system based on the conceptual design for TFCX. The neutral beam baseline was designed and estimated in detail for TIBER II by a potential manufacturer.

1.5.2.5 Fuel Handling Systems

The fuel injection system includes a high velocity injector during operation, as well as a mechanical injector and gas puffer during start-up. Most of the cost is the \$30 million estimated for a laser driven plasma jet based on a conceptual design. The fuel processing and fuel storage and receiving accounts are scaled and escalated from the historical data for the Tritium Systems Test Assembly (TSTA). The atmospheric processing account includes air detritiation units and monitors for the reactor, reactor maintenance and tritium buildings. The costs of the units are scaled from historical data. The water recovery system to extract tritium from the breeding shield coolant is estimated by Ontario Hydro based on historical data.

1.5.2.6 Other Plant Systems

The vacuum system estimate is based on unit costs for commercially available pumps, valves and ducting, and on system estimates for the Compact Ignition Tokamak (CIT) study. The power supply and electric plant equipment estimate is based on equipment lists developed in the design study and on unit prices for commercial equipment, with a 30% allowance for installation. The water heat transport system costs are scaled from similar commercial systems. The cryogenic system cost is based on the Mirror Fusion Test Facility (MFTF) data, scaled for capacity. The allowances for maintenance equipment and instrumentation and control (including diagnostics) are based on other reactor design studies and preliminary equipment lists for TIBER II. The miscellaneous plant equipment allowance is based on commercial plant experience, and the heat rejection system cost is scaled from existing systems. All of the cost estimates for the general plant systems included here can be improved by more detailed designs, equipment lists and system layouts.

1.5.3 INDIRECT COST AND CONTINGENCY

The indirect cost allowance for TIBER II is taken as 30% of the direct cost and process contingency (defined below). This allowance includes the construction services and equipment, the project management, the project systems engineering and related technical services, and the construction management. These services are provided to the entire project and are not charged directly to any system.

The process contingency allowance accounts for the uncertainties in the design and provides for increases in quantities as the design becomes more developed. Allowances ranging from 0% (commercial process) to 30% (new concept) were applied to the major accounts according to their level of development. In addition, a project contingency of 10% is applied to all costs to allow for cost increases during the construction.

1.5.4 OPERATING COSTS

Operating costs were estimated for TIBER II using guidance from the Nuclear Energy Cost Data Base (NECDB, Ref. [4]). An operating staff of 400 is estimated, with annual salary rates corresponding to those for nuclear plants. This estimate does not include the scientific and nuclear testing experimental

staff, expected to be comparable in size to the operations staff. The supplies and material expenses include first wall and divertor replacements, as well as allowances for miscellaneous items. Further definition of the plant equipment is required to improve this estimate. The electric power requirements for the plant are assumed to be provided at an average rate of \$0.05 per kilowatt-hour. The only fuel cost of significance is that of tritium at an estimated production cost of \$10 million per kilogram. It is assumed that the breeding shield will provide total tritium self sufficiency, and that the initial tritium inventory will be recovered at the end of the plant life, so that no operating costs are allowed for fuel. At the end of the plant life, the decommissioning expense is estimated to be \$100 million. The annual operating costs levelized over the plant life of 12 years with no interest or inflation effects are listed on Table 1.5-2.

The worksheet used to estimate all costs is included as Table 1.5-3. These data can be used to study cost sensitivities and perform trade studies near the TIBER II design point, but many of the factors are design-specific and should not be applied to other designs.

Table 1.5-2. TIBER II operating cost (millions of dollars per year).

ANNUAL OPERATING COST	86.15
Personnel	18.16
Misc. Supplies	5.00
Maint. Matls	12.61
First Wall	1.01
Divertor	6.60
Other	5.00
Electric Power	42.05
Fuel	0.00
Decommissioning	8.33

Table 1.5-3. TIBER II cost worksheet.

Account	Description	equipmt	install	eng'g	total	scaling	quantity	algorithm
21.	STRUCTURES			0.00	164.71			
211.	SITE IMPROVEMENTS			0.00	12.00	fixed		
212.	REACTOR BUILDING			0.00	37.37	volume,m^3	89056	4100*V^0.8/1e6
213.	TURBINE BUILDING			0.00	0.00	volume,m^3	0	
214.	MAINTENANCE BUILDING			0.00	50.02			
214.1	Hot Cells/Decon			0.00	40.86	volume,m^3	99572	4100*V^0.8/1e6
214.2	Support Areas			0.00	9.16	volume,m^3	64554	1300*V^0.8/1e6
215.	TRITIUM BUILDING			0.00	26.10			
215.1	Process Area			0.00	23.96	volume,m^3	51100	4100*V^0.8/1e6
215.2	Support Area			0.00	2.14	volume,m^3	10500	1300*V^0.8/1e6
216.	ELECTRICAL BLDG			0.00	7.96	volume,m^3	41800	1600*V^0.8/1e6
217.	OTHER BUILDINGS			0.00	31.25			
217.1	Administration				9.19	volume,m^3	50000	1600*V^0.8/1e6
217.2	Shops & Stores			0.00	9.09	volume,m^3	88700	1000*V^0.8/1e6
217.3	Cryogenic			0.00	2.98	volume,m^3	22000	1000*V^0.8/1e6
217.4	Miscellaneous			0.00	10.00	allowance		
22.	REACTOR EQUIPMENT				1217.88			
221.	REACTOR SYSTEM				255.73			
221.1	FIRST WALL	12.08	3.11	2.83	18.02			
221.1.1	Inboard	6.22	1.68	1.40	9.30	area,m^2	56	tfcx \$/cm^2:mat 11.1,i3,e2.5
221.1.2	Outboard	5.86	1.43	1.43	8.72	area,m^2	143	tfcx \$/cm^2:mat 4.1,i 1,e 1
221.2	BLANKET				0.00			not included
221.3	SHIELD	118.95	11.89	29.74	160.58			10% install, 25% engg
221.3.1	Bulk Materials	38.20						
	Structural Steel	0.56				wt,kg	214938	2.6\$/kg bulk
	Steel Balls	1.58				wt,kg	478204	3.3\$/kg bulk
	Tungsten Block	12.60				wt,kg	273970	46\$/kg bulk
	Tungsten Balls	4.59				wt,kg	77797	59\$/kg bulk
	Beryllium Plate	0.38				wt,kg	663	570 \$/kg bulk
	Beryllium Balls	18.49				wt,kg	32437	570 \$/kg bulk
	Fabrication	65.00						Shop Estimate
221.3.2	Port Plugs	15.75				wt,kg	250000	63\$/kg
221.4	STRUCTURE	23.03	2.30	2.30	27.63			
221.4.1	Shield Support	4.62	0.46	0.46	5.54	wt,kg	165000	28\$/kg+10%inst+10%engg
221.4.2	IF Intercoil	7.00	0.70	0.70	8.40	wt,kg	250000	28\$/kg+10%inst+10%engg
221.4.3	Outer PF Coil	2.10	0.21	0.21	2.52	wt,kg	75000	28\$/kg+10%inst+10%engg
221.4.4	Center Post	0.41	0.04	0.04	0.50	wt,kg	14900	28\$/kg+10%inst+10%engg
221.4.5	Cold Mass Support	2.52	0.25	0.25	3.02	wt,kg	90000	28\$/kg+10%inst+10%engg
221.4.6	Vacuum Vessel	6.37	0.64	0.64	7.65	mass,kg	227600	28\$/kg
221.5	DIVERTOR	39.60	0.00	9.90	49.50	area,m^2	72	tfcx \$/cm^2:mat 55,i 0,e25%

Table 1.5-3. TIBER II cost worksheet (continued).

Account	Description	equipmt	install	eng'g	total	scaling	quantity	algorithm
222.	MAGNETS	198.30	9.92	19.83	228.05			
222.1	TF COILS (16)	92.22	4.61	9.22	106.05			5% installation, 10% engg
222.1.1	Conductor	51.48				ILB, R _{max} T	1.72E+10	3\$/kR _{max} T
222.1.2	Winding	29.64				pack wt, kg	247000	120\$/kg
222.1.3	Case	11.10				case wt, kg	148000	75\$/kg
222.2	PF COILS	106.08	5.30	10.61	122.00			
222.2.1	Outer PF (4)	50.63	2.53	5.06	58.22			5% installation, 10% engg
	Conductor	12.78				ILB, R _{max} T	4.26E+09	3\$/kR _{max} T
	Winding	19.15				pack wt, kg	159600	120\$/kg
	Case	18.70				case wt, kg	249300	75\$/kg
222.2.2	Top & Bottom (4)	29.92	1.50	2.99	34.40			5% installation, 10% engg
	Conductor	11.82				ILB, R _{max} T	3.94E+09	3\$/kR _{max} T
	Winding	13.44				pack wt, kg	112000	120\$/kg
	Case	4.66				case wt, kg	62070	75\$/kg
222.2.3	Solenoid	25.54	1.28	2.55	29.37			5% installation, 10% engg
	Conductor	12.75				ILB, R _{max} T	4.25E+09	3\$/kR _{max} T
	Winding	8.69				pack wt, kg	72400	120\$/kg
	Case	4.10				case wt, kg	54690	75\$/kg
223.	POWER INJECTION				270.00			
223.1	ECRH System				30.00	injected w	1E+07	3\$/w allowance
223.2	Lower Hybrid				60.00	injected w	2E+07	3\$/w allowance
223.3	Neutral Beam				180.00	injected w	5E+07	3.6\$/w allowance
224.	VACUUM SYSTEM	19.74	5.92	0.00	25.66			30% installation
224.1	VACUUM PUMPS	9.04				no. of pumps	64	85k\$ per 1.8m ³ /s dt turbo
224.1.1	High Vacuum Pumps	5.44				no. of pumps	16	225k\$ ea.
224.1.2	Backing Pumps	3.60						
224.2	DUCTING	9.70						
224.2.1	Vacuum Duct	7.20				length, m	480	15k\$ per m, .6 m diam.
224.2.2	Valves	2.50				no. of valves	32	78k\$ ea., .6 m diam.
224.2.3	Duct Shield	0.00				wt, kg	0	20 \$/kg
224.3	SUPPORT SYSTEMS	1.00						
224.3.1	Instr & Control	1.00				allowance		
225.	POWER CONDITIONING	77.20	23.16	0.00	100.36			30% installation
225.1	TF Coil Power Supply	10.60	3.18	0.00	13.78			equipment list from sc
225.2	PF Coil Power Supply	33.00	9.90	0.00	42.90			equipment list
225.3	Energy Storage	33.60	10.08	0.00	43.68			equipment list

Table 1.5-3. TIBER II cost worksheet (continued).

Account	Description	equipmt	install	eng'g	total	scaling	quantity	algorithm
226.	HEAT TRANSPORT				64.86			
226.1	REACTOR COOLING	11.36	3.41	0.00	14.76			30% installation
226.1.1	Heat Exchangers	4.87				pth,u	3.50E+08	.62*nhx*(u/nhx)^.8, 200MW/hx
226.1.2	Pumps & Pipe	6.48				pth,u		.68*SUM(u/loop)^0.8*nloops
	First Wall	1.80					9.00E+07	2 loops
	Shield	3.55					2.10E+08	2 loops
	Diverter	1.13					5.00E+07	2 loops
226.2	CRYOGENIC SYSTEM				34.03			
226.2.1	Helium				32.03	watts @ 4.2K	70000	13000*watts^0.7
226.2.2	Nitrogen				2.00	allowance		
226.3	AUXILIARY COOLING	12.36	3.71	0.00	16.07			30% installation
226.3.1	Heat Exchangers	7.64				pth,u	5.70E+08	.45*nhx*(u/nhx)^.8, 50MW/hx
226.3.2	Pumps & Pipe	4.72				pth,u		.68*SUM(u/loop)^.8*nloops
	Neutral Beams	1.96					1.00E+08	2 loops
	ECH	0.47					2.00E+07	1 loop
	LHRF	0.47					2.00E+07	1 loop
	Cryogenic	0.74					3.50E+07	1 loop
	Vacuum	0.16					5.00E+06	1 loop
	Tritium	0.27					1.00E+07	1 loop
	Facilities	0.65					3.00E+07	1 loop
227.	FUEL HANDLING				103.22			
227.1	Fuel Injection				37.10			
227.1.1	Gas Puffer				0.10	lot	1	100 k\$ lot
227.1.2	Mechanical Gun				7.00	no.	1	7 M\$ ea.
227.1.3	Laser Jet				30.00	no.	1	30 M\$ ea.
227.2	Fuel Processing				7.27	D+T f,kg/s	3.34E-05	M=2.31+92.6*f^0.3+363*f^0.6
227.3	Storage and Receiving				6.74	inventory,g	3000	M=3.67+12.4*f^0.3+.77e-3*g
227.4	Atmospheric Processing				27.10			RTR m3/hr=V/t/eff*ln(Ci/m3)
227.4.1	ATR Units				17.11			t=120h, eff=.9, cf=5ue-bu1/ms
	Reactor Bldg				3.90	m^3/hr	5417	m^3/hr=U/108*ln1.93e9/V;
	Rx Bldg Bsat(MTS area)				1.76	m^3/hr	1733	\$=9500*(m3/hr)^.7
	Hot Cell				3.94	m^3/hr	5487	
	Hot Cell Bsat(RadMste)				2.54	m^3/hr	2926	
	Tritium Building				3.08	m^3/hr	3658	
	NBI Building				1.90	m^3/hr	1930	
227.4.2	Monitors				10.00			680\$/m^2; 1 monitor / 8 m^2
	Reactor Building				1.09	floor m^2	1600	
	Rx Bldg Bsat(MTS area)				1.05	floor m^2	1600	
	Hot Cell				1.93	floor m^2	2840	
	Hot Cell Bsat(RadMste)				1.93	floor m^2	2840	
	Tritium Building				3.14	floor m^2	4620	
	NBI Building				0.82	floor m^2	1200	
227.5	Water Recovery				25.00			ontario hydro est=30 Ci/1

Table 1.5-3. TIBER II cost worksheet (continued).

Account	Description	equipat	install	eng'g	total	scaling	quantity	algorithm
228.	INSTR. & CONTROL				100.00	allowance		
229.	MAINTENANCE EQMT				70.00	allowance		
23.	TURBINE PLANT				0.00			
24.	ELECTRIC PLANT	35.00	10.50	0.00	45.50			30% installation
241.	Pulsed Power System					power,u	3.50E+08	equipment list
242.	Facility Power					power,u	3.50E+07	equipment list
243.	Standby Power					fixed		1.6M# d gen, .4M# no-break
244.	Distribution					power,u		
25.	MISCELLANEOUS			0.00	26.00	allowance		
251.	Cranes/Transporters			0.00	4.00	allowance		
252.	Air/Water/Steam/Fire			0.00	15.00	allowance		
253.	Communications			0.00	2.00	allowance		
254.	Furnishings			0.00	2.00	allowance		
255.	Waste Treat			0.00	3.00	allowance		
26.	HEAT REJECTION			0.00	15.16	rej heat,u	5.70E+08	$1.5 \times 10^8 / 1e6$
2.	TOTAL DIRECT COST				1469			

Table 1.5-3. TIBER II cost worksheet (continued).

Account	Description	equipmt	install	eng'g	total	scaling	quantity	algorithm
9.	INDIRECT COST				510			30% of direct & proc cont.
91.	Constr Serv & Eqt				85			5%
92.	Engineering & Services				170			10%
93.	Field Office				170			10%
94.	Project Mngmt & Spares				85			5%
	CONTINGENCY				451			
	PROCESS CONTINGENCY				230			
	21. Buildings				0			0%Commercial Process - 0%
	221.1 First Wall				4			20%Full Scale Prototype - 10%
	221.2 Blanket				0			20%Bench Scale/Detail Des- 20%
	221.3 Shield				32			20%New Concept - 30%
	221.4 Structure				6			20%
	221.5 Divertor				10			20%
	222. Magnets				57			25%
	223.1 ECRH				6			20%
	223.2 LHRH				12			20%
	223.3 NBI				45			25%
	224. Vacuum				3			10%
	225. Power Cond.				10			10%
	226. HTS				3			5%
	227. Fuel Proc				15			15%
	228. IAC				10			10%
	229. Maint. Eqt.				18			25%
	24. Electric Plant				0			0%
	25. Misc.				0			0%
	26. Heat Rej.				0			0%
	PROJECT CONTINGENCY				221			10% direct, ind & proc con
	TOTAL CONSTRUCTED COST				2430			
	OPERATING COST							
	ANNUAL OPERATING COST				86.15			
	Personnel				18.16	staff, no.	400	45,400 \$/py needb
	Misc. Supplies				5.00	allowance		
	Maint. Matls				12.61			
	First Wall				1.01	\$/plant life	1	12 year life; 8.3% repl/yr
	Divertor				6.60	\$/plant life	2	16.6 % replaced/yr
	Other				5.00	eqt. cost		allow 5M\$
	Electric Power				42.05	kwhr/yr	8.41E+08	220MWh@30%; 30MWh@100%; .05\$/kw
	Fuel				0.00	net t use, kg	0.00	tritium @ 10M\$/kg; D, Li negl
	Decommissioning				8.33	allowance		100M\$ at end of life

1.6 WORK BREAKDOWN STRUCTURE (Doggett)

The WBS that follows is included to show the reader how this study was organized and to allow him to contact specific individuals about specific topics. The list of names does not include everyone who worked on each topic but should aid the reader in tracking down the right individual(s). A complete list of the participants can be found at the beginning of the report.

ETR/TIBER WORK BREAKDOWN STRUCTURE

1.0.0.0.	Project Management	Henning, Logan
1.1.0.0.	Configuration	Bulmer
1.2.0.0.	Design Integration	Doggett
1.2.1.0.	Systems Requirements	Lee
1.3.0.0.	Cost and Schedule	Thomson, Doggett
1.4.0.0.	Operations	Perkins
2.0.0.0.	Plasma Engineering	Perkins, Logan
2.1.0.0.	Zero-D Plasma Model	Fenstermacher
2.2.0.0.	Systems Code - See 7.0.0.0.)	
2.3.0.0.	Scoping Studies	Peng, Perkins
2.4.0.0.	Time-Dep. Zero-D Code	Blackfield
2.4.1.0.	Code Development	
2.4.2.0.	Startup	
2.5.0.0.	Current Drive	Devoto
2.5.1.0.	ECH	Smith, Nevins, Cohen
2.5.2.0.	Neutral Beam	Devoto, Stacey
2.5.3.0.	Fast Wave	Ehst, Maub
2.6.0.0.	Edge Physics	Barr
2.6.1.1.	Divertor Design	Barr
2.6.1.2.	Divertor Physics	Brooks
2.6.1.3.	Transport	Prinja
2.7.0.0.	Impur. & Ash Removal	Miley, Mendrekis
2.8.0.0.	Alpha Physics	Miley
2.8.1.0.	Transport	
2.8.2.0.	Alpha Driven Instab.	
2.8.3.0.	Burn Control	
2.9.0.0.	Fueling	Perkins, Ho
2.10.0.0.	Disruption Modeling	Merril, Perkins, Bruhn
2.11.0.0.	Disruption Control	Smith
2.12.0.0.	Magnetics and MHD	Bulmer
2.13.0.0.	PF Magnetics	Schultz, J.
2.14.0.0.	Vertical Stability	Lever

3.0.0.0.	Systems Engineering	Doggett
3.1.0.0.	Mechanical Systems	Neef
3.1.1.0.	Configuration	Lousteau, Neef
3.1.2.0.	Divertor	Schultz, K.
3.1.2.1.	Studies	Haines
3.1.2.2.	Design	Gallix
3.1.2.3.	Materials	Whitley, Wilson
3.1.2.4.	Analysis	Listvinski
3.1.2.5.	Shield Integration	Sviatoslavsky, Berwald
3.1.3.0.	Wall Armor & Limiter	Schultz, K.
3.1.3.1.	Studies	Haines
3.1.3.2.	Design	Gallix
3.1.3.3.	Materials	Whitley, Wilson
3.1.3.4.	Analysis	Listvinski
3.1.3.5.	Shield Integration	Sviatoslavsky, Berwald
3.1.4.0.	Vacuum Vessel	Lousteau, Neef
3.1.5.0.	Reactor Structure	Myall, Johnston
3.1.6.0.	Magnets	Miller
3.1.6.1.	PF Systems Analysis	Schultz, J.
3.1.6.2.	Coil Design	Miller
3.1.6.3.	Control Coils	Neef
3.1.6.4.	Coil Structure	Myall
3.1.6.5.	Conductors	Kerns
3.1.6.6.	Materials	Summers
3.1.7.0.	Cryogenics	Slack
3.1.9.0.	Fueler	Gorker, Haines
3.2.0.0.	Electrical Systems	Nelson, Flanagan
3.2.1.0.	ECH	Boehmer
3.2.2.0.	LH	Nelson/Flanagan, Yugo
3.2.3.0.	Negative Ion Beams	Fink, Berwald
3.2.4.0.	Diagnostics	Nelson/Flanagan, Gorker
3.2.5.0.	Power Supplies	Nelson/Flanagan, Gorker
3.2.6.0.	Controls	Nelson/Flanagan, Gorker
3.3.0.0.	Nuclear Anal. & Design	Kulcinski
3.3.1.0.	Concepts	Lee, Sawan, Steiner
3.3.2.0.	Design & Analysis	Sawan, Steiner
3.3.2.1.	Breeding Options	Sawan
3.3.2.2.	Mechanical Design	Sviatoslavsky, Deutsch
3.3.2.3.	Thermal Hydraulics	Sviatoslavsky, Deutsch
3.3.2.4.	Neutronics	Sawan, Steiner
3.3.2.5.	Activation	Attaya
3.3.2.6.	Inboard Shield Optim.	El-Guebaly
3.3.2.7.	Trit./Matls Interaction	Wittenberg
3.3.3.0.	Tritium Processing	Gierszewski
3.3.4.0.	Tritium Supply	Gierszewski
3.3.5.0.	First Wall/Armor Intfc	Sviatoslavsky, Schultz, K., Deutsch
3.4.0.0.	Facilities	Thomson
3.4.1.0.	Reactor Building	
3.4.2.0.	Hot Cells	
3.4.3.0.	Maintenance Equip.	
3.4.4.0.	Plant Design	
3.4.4.1.	Electrical Distrib.	
3.4.4.2.	Heat Rejection	
3.4.4.3.	Vacuum System	

3.4.4.4.	Tritium	
3.4.4.5.	Cryogenics	
3.5.0.0.	Tritium System	Anderson
3.5.1.0.	Fuel Processing	Anderson
3.5.2.0.	Cleanup	Anderson
4.0.0.0.	Nuclear Testing Performance	Baker
4.1.0.0.	Test Requirements	Abdou
4.2.0.0.	First Wall/Blnkt Tests	Mattas
4.2.1.0.	Liquid Metal-Cooled	Mattas
4.2.2.0.	Helium-Cooled	Schultz, K.
4.2.3.0.	Water-cooled	Abdou
4.2.4.0.	Flibe-Cooled	Sze
4.3.0.0.	Materials Tests	Powell
4.4.0.0.	High Heat Flux Comp Test	Gauster, Mattas
4.5.0.0.	Tritium System Tests	Finn, Anderson
4.6.0.0.	Neutronics Tests	Abdou
4.7.0.0.	Safety Tests	Piet, Sze
4.8.0.0.	R&D Assessment	Baker
5.0.0.0.	Rad. Safe. & Environ.	Piet, Stasko
5.1.0.0.	Design Requirements	Piet
5.2.0.0.	Accident Analysis	Piet
5.2.1.0.	Tritium Accidents	Stasko
5.2.1.1.	Sum. of T2 in Comp	Gierszewski
5.2.1.2.	T2 Release Potential	Stasko
5.2.1.3.	Air Cleanup	Anderson
5.2.2.0.	Activ. Prod. Accidents	Piet
5.2.2.1.	Product Inventory	Attaya
5.2.2.2.	Decay Heat Source Term	Attaya
5.2.2.3.	Release Potential	Piet
5.2.3.0.	Transient Analysis	Piet
5.2.3.1.	Plasma Disrup. Effects	Merrill
5.2.3.2.	Plasma Cntrl Failures	Perkins
5.2.3.3.	LOCA/LOFA Analysis B/S	Sviatoslavsky
5.2.3.4.	Liquid Metal Testing	Piet
5.3.0.0.	Personnel Safety/Maint	Stasko
5.4.0.0.	Effluent Control	Stasko
5.4.1.0.	Tritium	Stasko
5.4.2.0.	Air Activation	Attaya
5.4.3.0.	Coolant Activation	Embrechts
5.4.3.1.	Products in Coolant	Attaya
5.4.3.2.	Processing & Leakage	
5.5.0.0.	Waste Mgmt Decommiss.	Herring
5.5.1.0.	Activation Inventory	Attaya
5.5.2.0.	Waste Classification	Attaya
5.5.3.0.	Decommiss. Approach	Carney
5.6.0.0.	Facility Considerations	Piet
5.6.1.0.	Objectives for Bldgs	Piet
5.6.2.0.	Underground Siting	Stasko
5.6.3.0.	Apprch for Confinement	Piet
5.6.4.0.	Vented Filter System	Piet
5.7.0.0.	Conclusions & Summary	Piet, Stasko
5.7.1.0.	Future Work	Piet, Stasko

5.7.2.0.	Comparison to Targets	Piet, Stasko
5.7.3.0.	Potential Improvements	Piet, Stasko
6.0.0.0.	Development Program	Henning
7.0.0.0.	Systems Code Intro	Peng, Perkins
7.1.0.0.	Code Architecture	Reid
7.1.1.0.	Numer. Optimiz. Driver	Galambos
7.1.2.0.	Code Integration	Reid
7.1.3.0.	The Basis Shell	Dubois
7.1.4.0.	Post Processor (FY88)	Gladd
7.2.0.0.	Plasma Performance and Requirements	Fenstermacher
7.2.1.0.	Plasma Performance	Fenstermacher, Galambos, Uckan
7.2.2.0.	Plasma Current Drive and Heating	Fenstermacher, Ehst, Mau
7.2.3.0.	MHD Equilibrium Magnetics	Strickler, Bulmer, Miller
7.2.4.0.	Flux Linkage & Inductance	Kalsi, Galambos, Perkins
7.3.0.0.	Engineering Systems	Haines
7.3.1.0.	Mechanical Systems	Haines
7.3.1.1.	Torus Configuration	Galambos
7.3.1.2.	Torus Vacuum Systems	Haines
7.3.1.3.	Fueling Systems	Perkins, Ho
7.3.1.4.	Torus Support Struct.	Perkins, Myall
7.3.2.0.	Electrical Systems	Hicks
7.3.2.1.	ECH Systems	Blackfield, Wagner
7.3.2.2.	LH Systems	Yugo
7.3.2.3.	NB Systems	Perkins, Fink
7.3.2.4.	ICRH Systems	Yugo
7.3.2.5.	AC Power Systems	Gorker, Hicks
7.3.2.6.	Instrumentation and Control	Gorker, Hicks
7.3.3.0.	Magnetic Systems	Miller, Kalsi
7.3.3.1.	TF Magnets	Kerns, Kalsi, Galambos
7.3.3.2.	PF Magnets	Kerns, Kalsi, Galambos
7.3.3.3.	Magnet Power Supplies	Gorker, Hicks
7.3.4.0.	Facilities	Thomson
7.3.4.1.	Buildings and Plant Systems	Thomson
7.3.4.2.	Heat Transport Systems	Thomson
7.3.4.3.	Maintenance Equipment	Spampinato
7.4.0.0.	Nuclear Systems	Gohar
7.4.1.0.	First Wall	Gohar, Hassanein, Haines
7.4.2.0.	Blanket	Gohar
7.4.3.0.	Bulk Shield	Gohar
7.4.4.0.	Penetration Shield	Lee
7.4.5.0.	Reactor Vault Shield	Gohar
7.4.6.0.	Tritium Systems	Finn, Willms, Busigin
7.4.7.0.	Particle Control Syst	Brooks, Barr, Miller, Post
7.5.0.0.	Safety Systems (FY88)	Piet, Brereton
7.6.0.0.	Cost Accounting	Thomson
7.6.1.0.	Direct Cost	
7.6.2.0.	Operating Cost (FY88)	
7.6.3.0.	Research and Development Cost (in future)	

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2.0 PLASMA ENGINEERING

2.1 INTRODUCTION -- THE PHYSICS OBJECTIVES FOR TIBER

The general objective of TIBER--from a physics design viewpoint--is to seek a low-cost, innovative configuration to stimulate and influence national and international physics activities for the prospective Engineering Test Reactor (ETR). The result is a compact (3-m major radius), steady-state, current-driven tokamak ETR. Our philosophy in the extrapolation of advanced physics concepts in TIBER is to allow the application of promising concepts which could, in principle, be qualified in the worldwide fusion program (i.e., existing or planned test stands or large tokamaks) before the expected ETR design decision date (after 1993). In this way, one very important product of the TIBER design process is to drive R&D programs in the qualification of attractive ETR design concepts.

In Table 2.1-1, we list the major physics performance objectives for TIBER. In Table 2.1-2, we define the projected operating phases of the machine. Note, in Table 2.1-1, that the plasma performance objectives are designated as "nominal" operating requirements and "maximum" operating requirements. The former are assumed to apply during the nuclear testing phase (Phase III in Table 2.1-2) to reduce the possibility of disruptions and to improve testing reliability; these requirements form the basis of the baseline TIBER design discussed in detail in Section 2.4. However, the reactor systems of TIBER should be capable of also sustaining the "maximum" performance objectives in Table 2.1-2, corresponding to operation at combinations of full critical beta, maximum density limits, and minimum edge-safety factor ($q_{min}(a) \sim 2.2$). Operation at maximum performance is intended for limited periods during the physics testing phase (Phase II in Table 2.1-2). The characteristics of selected high performance, steady-state operating scenarios are discussed in Section 2.5 below.

Table 2.1-1. Performance requirement.

Requirements for Nominal Operation

- $R = 3 \text{ m}$
- $Q \geq 5^a$ (i.e., $P_a \geq P_{CD}$)
- Minimum average wall loading $\langle \Gamma \rangle = 1 \text{ MW m}^{-2} \text{ }^b$
- Minimum plasma current $I_p = 10 \text{ MA}^c$
- Non-inductive current-drive
- Steady-state burn with equilibrium impurity and alpha ash removal
- Minimum peak fluence at end-of-life $\phi t = 3 \text{ MW yr/m}^2 \text{ }^d$
- Minimum machine testing lifetime = 10 years^d
- $\langle n_e \rangle \leq 0.75 n_{e,\text{max}}^{e,f}$

Requirements for Maximum Operation

As nominal operation with the addition of

- $q_{\text{edge}} \geq 2.2^g$
- $\beta = \beta_{\text{crit}}$ (full Yamazaki ballooning limit)
- $\langle n_e \rangle \leq n_{e,\text{max}}$
- Maximum peak wall loading = $2 \text{ MW m}^{-2} \text{ }^h$

^aQ is defined by $P_{\text{fusion}}/P_{\text{current drive}}$.

^bIn the current design, this translates to a peak wall load of $\sim 1.54 \text{ MW m}^{-2}$.

^cFor the steady-state current drive baseline; this may be less under inductive operation (see Section 2.6 below).

^dThe definition of peak fluence and operating lifetime translates to a requirement for a full power lifetime of $\sim 6.3 \times 10^7$ burn-seconds.

^eThis limit is important only for pulsed inductive operation. For steady-state, current-driven operation, the density is always well below this limit.

^fWhere $n_{e,\text{max}}$ is the Greenwald density limit (see Sec. 2.2 below).

^gIn the current design, this translates to a plasma current of $\sim 14.5 \text{ MA}$.

^hEngineering constraint set by shielding limitations in a 3-m major radius machine.

Table 2.1-2. Projected operating scenario.

Phase	Duration	Emphasis	Availability	Annual tritium ^a consumption (kg/yr)	Integrated peak ^b fluence per phase (MW yr/m ²)
I	1	H checkout	10%	----	----
	1	DT checkout	5%	0.88 (0.22)	0.1 (0.03)
II	3	DT physics testing	10%	1.75 (0.45)	0.6 (0.15)
III	7	DT engineering testing	30%	5.26 (1.34)	4.2 (1.07)

^aTotal T₂ consumption at end-of-life = 43 (11) kg. The parameters shown under tritium consumption and fluence are for the neutral-beam-driven, steady-state baseline (Section 2.4). The values in parentheses are for the contingency pulsed operation (Section 2.6).

^bTotal accumulated peak fluence of end-of-life = 4.9 (1.3) MW yr/m²

The operating scenario requirements for TIBER shown in Table 2.1-2 are composed of three phases:

Phase I. Checkout Phase - 1-year hydrogen/deuterium checkout followed by 1-year DT checkout, both at low machine availability (< 10% and < 5% respectively)

Phase II. Physics Mission Phase - 3-year DT burning for physics missions at < 10% availability. High performance operation will be attempted for short periods during this phase.

Phase III. Nuclear Testing Mission Phase - 7-year DT burning phase for nuclear testing mission at 30% availability. Capability for true steady-state (> 1 week) should be demonstrated during this phase.

Note that the total machine lifetime is defined as 12 years (2 years checkout + 10 years operation).

Steady-state operation via non-inductive current drive is a major objective of the TIBER design. Our rationale for this is based on four principal assertions:

1. By employing current drive, the ETR will be both promoting and evaluating a concept that will do much to enhance the attractiveness of the eventual tokamak power reactor.
2. At this stage of the tokamak development program it would appear to be of vital importance to have independent control of the current-profile and decouple it from the density and temperature profiles in the machine.
3. Our goal of attaining high fluence through high machine availability (~ 30% in final phase) is much more credible.
4. Steady-state operation is highly desirable for successful engineering testing of nuclear components; plasma cycling causes time-dependent changes in environmental testing conditions such as thermal characteristics, tritium-concentration profiles, eddy-current effects, and failure of fracture modes.

Until both a large tokamak current drive test at ~ 1 MA and tests of a prototype current-drive source technology on a test stand dictate a winning current-drive option, it was decided that the TIBER II/ETR study should continue to examine various current-drive methods with regard to their suitability for meeting the performance requirements in Table 2.1-1. Current-drive efficiencies were required to include the classical bootstrap amplification for self-consistent, MHD stable profiles, as well as other calculable non-ideal effects. At this stage, the strongest contender is neutral-beam current-drive via high-energy, negative-ion neutral beams in combination with lower-hybrid edge current-drive; this forms the basis of the TIBER baseline design and is considered fully in Section 2.4 below. Other possible current-drive methods are discussed in appendixes at the end of this chapter.

Although the baseline requirement for TIBER is a steady-state current-driven machine, we have investigated, as a contingency, limited pulsed-inductive operation of TIBER employing whatever limited volt-seconds capability is available from the machine configuration for the baseline steady-state design. The purpose is to stress the flexibility of the design and to demonstrate that TIBER could be chosen as a basis for the prospective

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ETR design before we have sufficient time to qualify all aspects of steady-state current-drive with sufficient reduction of uncertainty. Note, however, that ignition per se is not a basic objective of TIBER. Pulsed-inductive and ignited scenarios are discussed in detail in Sections 2.5 and 2.6 below. Also investigated is the addition of lower hybrid to these pulsed scenarios to assist start-up, to increase the available volt-seconds during burn, and to maintain the plasma current during transformer reset (thereby potentially increasing the number of allowed stress cycles).

2.2 PHYSICS BASIS

In this section we review the principal physics relationships used to predict the performance of TIBER. Chief ingredients in the design are the maximum attainable β , denoted by β_c ; the maximum permissible plasma density, denoted by $\langle n_e \rangle_{\max}$; the scaling of the energy confinement time, τ_E ; and the relationship for the current drive efficiency, γ . The latter is defined by $\gamma = \langle n_{e20} \rangle R_0 I_p / P_{cd}$, where $\langle n_{e20} \rangle$ is the volume-averaged electron density in units of 10^{20} m^{-3} , R_0 is the geometric plasma major radius, I_p is the driven plasma current, and P_{cd} is the required power. In this expression and throughout this section, MKS units are employed, except that current is given in MA, power in MW, and energy or temperature in keV. In this section we also discuss the basis for the choice of certain other quantities necessary for predicting TIBER performance. Many of the formulas presented here are incorporated in a 0-D power balance code [1] which is described briefly in Appendix 2A. Calculations of TIBER performance presented in Sections 2.4, 2.5, and 2.6 are carried out with this code.

Before proceeding to the main discussion, we list in Table 2.2-1 the basic parameters of the TIBER design including the major radius; the minor radius, a ; elongation, κ ; triangularity, δ ; plasma current, I_p ; and toroidal magnetic field, B_0 , at the major radius. Definitions of R_0 , a , κ , and δ are given in Fig. 2.2-1. These quantities should be selected by an optimization procedure which finds some parameter, such as minimum cost, subject to constraints such as minimum acceptable Q and neutron wall load. A complete systems code was not available at the time these quantities were chosen, but they probably yield close to the minimum cost consistent with the engineering choices described elsewhere in this report.

2.2.1 EDGE SAFETY FACTOR

Experimental results have demonstrated that there is a lower limit on the edge safety factor for stable operation of a tokamak. The safety factor can be written in terms of local properties of the magnetic field as

$$q(\psi) = \frac{R B_t}{2\pi} \int_{\psi} \frac{dl}{R^2 B_p}, \quad (2.2-1)$$

Table 2.2-1. Basic TIBER parameters.^a

Major radius, R_0 -m	3.00	Minor radius, a -m	0.83 ^b
Elongation, κ	2.4 ^b	Triangularity, δ	0.4 ^b
Magnetic field, B_0 -T	5.55	Plasma current, I_p -MA	10

^aGeometric quantities are defined in Fig. 2.2.1

^bAt the 95% flux surface. The corresponding quantities on the separatrix are $a = 0.855$, $\kappa = 2.74$, $\sigma = 0.77$.

where B_t (B_p) are the toroidal (poloidal) magnetic field, respectively; R is the major radius coordinate; $R B_t$ is a function only of ψ , the poloidal flux, which is defined by

$$\psi = \frac{1}{2\pi} \int B_p \, dS_p, \quad (2.2-1)$$

and the line integral is evaluated around the closed flux contour. The plasma edge can be considered to be at the separatrix since this boundary separates the confined from the unconfined plasma. However, since B_p vanishes at the x point, the integral in Eq. (2.2-1) is undefined on the separatrix. In agreement with many other workers, we take the effective boundary at the surface enclosing 95% of the total flux enclosed by the separatrix. The safety factor on this boundary is denoted by q_ψ and has the value 3.7 for the TIBER magnet design (see Section 2.3).

Various approximate representations of the safety factor have been used. We quote some of them here with the TIBER values for the 95% flux surface. The simplest is the form valid for a circular cylinder,

$$q_{cyl} = \frac{5a^2 B_0}{I_p R_0} \quad [= 0.64 \text{ for TIBER}]. \quad (2.2-3)$$

A simple extension to elongated plasmas is the INTOR expression,

$$q_I = q_{cyl} \frac{1 + \kappa^2}{2} \quad [= 2.15 \text{ for TIBER}]. \quad (2.2-4)$$

Figure 2.2-1. Definition of geometrical quantities describing plasma shape.

A variant employed by the NET team is

$$q_{\text{NET}} = q_{\text{cyl}} \left(\frac{1}{2\pi} \frac{dl}{a} \right)^2 \quad [= 2.06 \text{ for TIBER}]. \quad (2.2-5)$$

Another variant from ORNL which makes some allowance for the triangularity is

$$q_{\text{ORNL}} = q_{\text{cyl}} \frac{1 + \kappa^2(1 + 2\delta^2)}{2} \quad [= 2.74 \text{ for TIBER}]. \quad (2.2-6)$$

It should be noted that none of these approximate expressions adequately represents the value of q_y for TIBER.

It has been observed that disruptions occur if the edge safety factor drops to values approaching 2. Various rules of thumb have been proposed for choosing the operating q_y . One rule is to maintain the $q_y = 2$ surface at least 5 cm inside the plasma surface. TIBER easily satisfies this rule since the $q_y = 2$ surface lies more than 10 cm inside the plasma surface. Another suggested rule is that the design must maintain the edge q greater than 3.0. We see that TIBER also satisfies this rule.

2.2.2 BETA LIMITS

It has become conventional to represent the maximum or critical β by the Troyon formula,

$$\beta_c (\%) = g I_N, \quad I_N = \frac{I_0}{a B_0}, \quad (2.2-7)$$

where various authors have selected the "constant" g in the range from 2.2 to 4.0%. In their original theoretical work Troyon et al. [2] found that their calculations were best fit with $g = 2.8\%$. However, they were mostly concerned with INTOR and JET and considered only a limited range of elongation, triangularity, and aspect ratio, $A = R_0/a$. They examined ballooning, including the Mercier criterion, and external $n = 1$ kinks with no conducting wall. Kinks were found to impose the limit on β for most of the range of currents they considered. More recent work on kinks shows that g should depend on the edge safety factor, q_y .

Yamazaki et al. [3] have carried out calculations of β_c due to infinite- n ballooning and the Mercier criterion. They examine a large range of

elongation, triangularity, and aspect ratio and show that the simple scaling of Eq. (2.2-7) cannot be used to fit their results. Instead, they suggest a different formula,

$$\beta_c(\%) = 4.7 I_N \left\{ 1 - 0.065 \frac{A}{\kappa^{1/2}} \left[1 - \frac{(\kappa - 1 + 0.05 \delta)}{\kappa^{1/2}} \left(\frac{1 + 1.5 \delta}{\kappa} \right)^{5/2} \right] \right\} I_N, \quad (2.2-8)$$

For studies of the baseline steady-state operation of TIBER we have selected the operating β from

$$\beta = \text{Min} \begin{cases} \beta_c \text{ from Eq. (2.2-7) with } g = 3.0\% \\ 0.75 \times \beta_c \text{ from Eq. (2.2-8)} \end{cases}$$

The factor 0.75 multiplying the Yamazaki formula is inserted to allow for a possible degradation of the ballooning β_c because of other instabilities such as external kinks. For the baseline parameters presented later in Section 2.4, we find that this modified Yamazaki limit is more severe. Coincidentally, the β so computed corresponds to a g coefficient of 2.8% in Eq. (2.2-7), just the value originally recommended by Troyon et al. [2]. β_c from Eqs. (2.2-7) and (2.2-8) is plotted for TIBER parameters vs I_p in Fig. 2.2-2.

There is considerable theoretical evidence that external kinks will be stabilized by a close-fitting conducting wall. Tsunematsu et al. [4], for example, find that a conducting wall at $a_w = 1.5a$ causes the kink and ballooning limits to coalesce. Their results are shown in Fig. 2.2-3 for a plasma with $\kappa = 1.6$ and $\delta = 0.3$. From the figure we see that the wall does not affect the loss of kink stability when $q_p \rightarrow 2$. The dip in the kink stability limit at $q_p = 3$ when the wall is at ∞ , which leads to the suggested restriction of $q_p > 3$ (see Section 2.2.1), is also evident in Fig 2.2-2. However, we see that this dip vanishes when the wall is brought into $a_w = 1.5a$, suggesting that we might be able to operate TIBER at a lower edge safety factor and thus higher plasma current than called for in the baseline design.

Some (but not all) experimental results support the absence of external kinks as a limit on the achievable β . The ASDEX workers [5], for example, do

Figure 2.2-2. β_c from Eq. (2.2-7) (Troyon) and Eq. (2.2-8) (Yamazaki).

Figure 2.2-3. Ballooning and external kink stability limits for $\kappa = 1.6$ and $\sigma = 0.3$ vs. q_ψ (from Ref. 4).

not observe low- n kinks at the maximum registered β values. For parameters and profiles of an ASDEX discharge at β values 5% below their maximum observed β_c , they find that the plasma is both ideal kink- and ideal ballooning-stable [5]. Nonetheless, they report a maximum Troyon q -value of 2.8% [see Eq. (2.2-7)] for their experiments. It has been suggested [6] that this stability limit in ASDEX is due to high- n resistive ballooning modes. In a higher temperature reactor plasma, one might expect that the resistive modes would approach the ideal modes. In this case, if the wall stabilizes external kinks, we might be able to operate at the ideal ballooning limit. For this reason we have also performed calculations of possible advanced TIBER performance at the full Yamazaki limit of Eq. (2.2-8). The results are described in Section 2.5 on alternate operating points.

2.2.3 DENSITY LIMITS

In experiments, it is found that there is a maximum density $\langle n_e \rangle_{\max}$ above which disruption-free operation is impossible. Murakami et al. [7] found that the limit could be related to the toroidal magnetic field and major radius by

$$\langle n_{e20} \rangle_{\max} = \nu \frac{B_0}{q_{\text{NET}} R_0}, \quad (2.2-9)$$

where q_{NET} is the safety factor defined in Eq. (2.2-5), and ν is a constant. It should be noted that most presentations of Eq. (2.2-9) use q_{cyl} , since the relation is applied to near circular tokamaks. For elongated plasmas q_{cyl} drops to very low values and we have replaced it by a more realistic q_{NET} . For most ohmic discharges, $\nu \approx 1.5$ seems to fit the data quite well, but data from pellet-fueled and neutral-beam-heated discharges suggest a higher value of ν for these conditions. More recently, Greenwald [8] has examined data on density limits and concluded that Eq. (2.2-9) yields a pessimistic limit for elongated plasmas such as TIBER. He suggests that the limit on line-averaged density, $\bar{n}_e$, should scale as $\propto I_p/A_p$, where A_p is the plasma cross-sectional area, with coefficient near unity. To be a bit more conservative than Greenwald suggested, we take a coefficient of 0.75 and write the "Greenwald limit" in the form

$$\langle n_{e20} \rangle_{\max} = 0.75 \frac{\langle n_e \rangle_{\text{e}} I_p}{\bar{n}_e A_p}. \quad (2.2-10)$$

The quantity $\langle n \rangle / \bar{n}$ is a numerical factor which depends on the assumed density profile. The Greenwald limit is used in this study but is generally of importance only for the alternate operating points with pulsed operation. For the baseline case described in Section 2.4 the density restriction is easily satisfied.

2.2.4 ENERGY CONFINEMENT SCALING

A number of scaling laws have been proposed for the total conductive/convective energy confinement time, τ_E . We have used the Kaye-Goldston relation [9] for prediction of baseline operation,

$$\tau_{KG} = 0.0553 H_f I_p^{1.24} P^{-0.58} R_0^{1.55} a^{-0.49} n_{e20}^{0.28} B_0^{0.26} A_1^{-0.09} \quad (2.2-11)$$

where the equivalent atomic mass A_1 is taken as $2(3)/2 = 3$. The quantity H_f is inserted to represent H-mode operation, wherein enhancements up to twice Kaye-Goldston scaling have been observed. In choosing our operating points we allow H_f to vary between 1.0, corresponding to L-mode operation, and 2.0, corresponding to H-mode operation. Although Eq. (2.2-11) has been deduced from data on discharges where P represents externally applied power, we choose to include in P both the internal power from α particles as well as the external, current-drive power, but we subtract out the power lost via radiation.

The confinement time given by Eq. (2.2-11) is thought to represent principally losses by electrons. In some experiments, energy losses through ions have been found to be much lower, and the equivalent ion heat diffusivity can be represented by a constant of order 3-10 times the neoclassical heat conductivity. In other experiments, the ion losses are found to be much worse, comparable to the electron losses. For this reason, in predicting TIBER baseline performance, we have adopted a conservative assumption that the ion energy confinement time, $\tau_{E,i}$, is the same as the electron energy time, $\tau_{E,e}$, and thus equal to τ_E given by Eq. (2.2-11).

In computing operating points for pulsed-mode cases in Sections 2.5 and 2.6, where there is ohmic current drive, we choose to use the two-regime model of Goldston [10],

$$\tau_E = [\tau_{NA}^{-2} + (H_f \tau_{KG})^{-2}]^{-1/2}, \quad (2.2-12)$$

where the ohmic regime neo-Alcator scaling is given by [11]

$$\tau_{NA} = 0.071 \langle n_{e20} \rangle R_o^2 (A_p/\pi)^{1/2} q_{NET} . \quad (2.2-13)$$

The original presentation of Eq. (2.2-13) [11] used q_{cyl} since the relation was derived from data on circular tokamaks. For elongated plasmas q_{cyl} drops to very low values, and therefore we have replaced it by the more realistic q_{NET} .

2.2.5 DENSITY AND TEMPERATURE PROFILES

Studies of radial transport, and thus of the temperature and density profiles, have not yet been carried out for TIBER. For the present, we assume that the radial density and temperature profiles are of the form

$$x(r) = x_0 [1 - (r/a)^2]^{\alpha_x}, \quad (2.2-14)$$

where x represents either density n or temperature T , and $x_0 = x(r=0)$. The same profile shapes are assumed for both electrons and ions. The exponents α_n and α_T are adjusted so that the resulting pressure profile $p(r) \propto n(r) T(r)$ matches the pressure profile assumed in the MHD equilibrium. The particular density and temperature profiles used are not unique; other combinations are certainly consistent with the equilibrium pressure. To be consistent with recent neutral-beam experiments on TFTR and JT-60, we have chosen the somewhat peaked density profiles $\alpha_n = 1.02$. The corresponding temperature profiles consistent with the equilibrium pressure yield $\alpha_T = 0.58$.

2.2.6 LOWER HYBRID CURRENT DRIVE

The current-drive efficiency γ is calculated for lower hybrid slow waves (hereafter denoted by LH) using the relativistic formulation of Fisch [12]. The normalized current density j per unit power density p is given by

$$j/p = \frac{G(Q)}{\beta_w} + \frac{\gamma_r X}{Q} \frac{dG}{dQ}, \quad (2.2-15)$$

where $\beta_w = \omega/k_y c = 1/N_{\parallel}$, $N_{\parallel}$ is the parallel index of refraction for the waves, $Q = p/mc$, $X = p_{\parallel}/mc$, and γ_r is the relativistic factor given by

$\gamma_r = (1 + Q^2)^{1/2}$. Here, p and $p_{\parallel}$ are the total and parallel momenta, respectively, of the current-carrying electrons. The function $G(Q)$ is given by

$$G(Q) = \frac{2}{Q} \left[\frac{\gamma_r + 1}{\gamma_r - 1} \right]^Z \int_0^Q d\psi \left[\frac{\psi}{\gamma_r(\psi)} \right]^3 \left[\frac{\gamma_r(\psi) - 1}{\gamma_r(\psi) + 1} \right]^Z, \quad (2.2-16)$$

where $Z = (1 + Z_{\text{eff}})/2$ and Z_{eff} is the effective charge of the plasma ions. The nonrelativistic limit ($\gamma_r \rightarrow 1$) of Eq. (2.2-15) gives the more familiar form for j/p , namely

$$(j/p)_{\text{nr}} = \frac{6Q}{5 + Z_{\text{eff}}} \left(X + \frac{1}{3} N_{\parallel} Q^2 \right). \quad (2.2-17)$$

Equations (2.2-16) and (2.2-17) are plotted in Fig. 2.2-4. There is about a factor of 2 difference between the relativistic and non-relativistic j/p under TIBER conditions ($X \approx 0.8$).

For the normalizations used in Eqs. (2.2-15) and (2.2-17) the current-drive efficiency is given by

$$\gamma = 31.2 (j/p) / \ln \Lambda_e, \quad (2.2-18)$$

where $\ln \Lambda_e$ is the Coulomb logarithm for electrons.

For LH current drive in TIBER, we assume that an electron tail with parallel energy $E \approx 5$ to $8T_e \approx 125$ to 200 keV, and perpendicular energy $E_{\perp} = T_e$ is formed by the waves. The parallel phase velocity of the waves is chosen to match the parallel velocity of these fast tail electrons. The parallel wave number must satisfy the accessibility condition

$$N_{\parallel} > N_{\parallel, \text{acc}} = \omega_{pe}/\omega_{ce} + (1 + \omega_{pe}^2/\omega_{ce}^2 - \sum \omega_{pi}^2/\omega_{ci}^2)^{1/2}, \quad (2.2-19)$$

where ω_{pa} and ω_{ca} are the plasma and cyclotron frequencies of species a . It is also necessary that the waves not suffer strong electron Landau damping (ELD) before they reach the desired penetration. This implies an upper bound on the usable $N_{\parallel}$ for current drive by LH slow waves given by [15]

Figure 2.2-4. Relativistic and non-relativistic j/p vs. $X = p/mc$ for lower hybrid current drive (from [12]).

$$N_{\parallel} < N_{\parallel, \text{ELD}} = 7.0 / T_e^{1/2}. \quad (2.2-20)$$

We use waves with $N_{\parallel}$ satisfying both the accessibility and strong ELD conditions, which allows penetration only to about $r/a \geq 0.6$. This is illustrated in Fig. 2.2-5 which shows Eqs. (2.2-19) and (2.2-20) under typical TIBER conditions. For this reason, we can use LH to drive current only in the outer regions of the plasma and must use another method for current drive in the core.

The relativistic efficiency formula Eq. (2.2-15) assumes that the bulk temperature is negligibly small ($T_e/mc \rightarrow 0$). The improvement in current-drive efficiency due to a warm background plasma has been calculated for LH waves by Karney and Fisch [13]. We use a parametrization of these results in the form of an enhancement factor for the efficiency given by

$$f_{\text{warm}} = 1 + (T_e^r / 25)^{1.16}. \quad (2.2-21)$$

where T_e^r is the electron temperature at the location of the LH absorption. Typically $T_e^r \leq 10$ keV, so that $f_{\text{warm}} \leq 1.3$ for TIBER parameters.

The current-drive efficiency for LH waves is degraded by electron trapping effects. The computation of the degradation factor f_t is carried out with the method by Cohen [14]. An additional loss is due to the launch of a fraction of the wave energy in the backward direction. For the baseline operation we assume that 20% of the power is lost in this way, but since typically $N_{\parallel} \geq 5$ in the backward direction, the energy is absorbed on the bulk electrons, heating them but not driving reverse current.

The method described in this section is incorporated in the computations with the zero-dimensional power balance code MUMAK [1] described in Appendix 2A. More detailed calculations including ray tracing of the waves and absorption with either (1) quasi-linear absorption in a cylindrical approximation to the plasma or (2) linear Landau damping with a two-dimensional MHD equilibrium, are covered in Appendix 2B.

Current drive with LH slow waves has been extensively studied in many tokamaks. Rather than reviewing these experimental results here, we note a recent review article [15] and summarize some important results. More recent work on LH current drive, including start-up, is covered in Refs. [16]-[19]. Experiments have demonstrated that:

Figure 2.2-5. Plot of $N_{//,acc}$ [Eq. (2.2-19)] and $N_{//,ELD}$ [Eq. (2.2-20)] vs. r/a in TIBER.

1. A large, steady-state current can be generated (>400 kA in PLT, 2 MA in JT-60 [16]).
2. Current can be ramped up with LH current drive (many experiments).
3. Current can be driven at high density provided the frequency is sufficiently high (to $n \approx 10^{20} \text{ m}^{-3}$ in Alcator C with $f = 4.6$ GHz [15]).
4. Current-drive efficiency is within a factor of 2 of the theoretical efficiency.
5. Current can be started with LH alone [17] or with ECH preionization followed by LH [18]-[19].

Special mention should be made of the very recent results in JT-60 [16]. This device has been used to drive a record 2 MA of current with 3.1 MW of LH power without OH assist. The current drive efficiency is in the range of 0.08 to 0.17 with LH alone and improves to 0.17 to 0.30 with simultaneous, near perpendicular injection of neutral beams. The increase has been attributed to electron heating by the neutral beams, but alternative explanations are possible. One is a decrease in Z_{eff} due to the deposition of hydrogen in the core of the plasma which, from Eq. (2.2-17) will raise the efficiency without a change in T_e . Another possibility is the generation of bootstrap currents by the enhanced density gradients which come from the neutral beams. Ushigusa et al. [20] report that the density profile obeys Eq. (2.2-14) with $\alpha_n = 1.9$, quite a steep profile, and that $\beta_p + I_i/2$ increases substantially during the neutral-beam injection. Part of the increase can be attributed to rearrangement of the plasma current leading to an enhanced internal plasma inductance, I_i , but there must also be a substantial increase in β_p . Since bootstrap current scales as β_p (see Section 2.2.8), it could be contributing to the observed increased current and improvement in efficiency with neutral-beam injection.

Because it may be necessary to use LH in TIBER to ramp up the current and plasma properties before applying neutral beams for current drive, we should look more closely at the results of the recent startup experiments. In the WT-2 experiment [17], ECH at the fundamental (35.6 GHz) was used to preionize the plasma, following which LH waves at 915 MHz were injected. Curiously, the ECH was found to drive a plasma current of 0.5 kA, although it was not intended to do so. After application of the LH radiation, current increased

to 5 kA. In the JIPP T-IIU experiments [18], preionization at the same ECH frequency was used, followed by injection of LH at 800 MHz. This experiment differed from WT-2 since the current of up to 20 kA was generated for much longer than the characteristic L/R time of the plasma. In the PLT experiments [19], two different LH launchers were used: one with $N_H = 3.5$, and one with $N_H = 2$. Initially, the waveguide phasing was 0° for breakdown and was switched to 90° or 135° for current drive. With the $N_H = 2$ antenna, special programming of the vertical-field coils was needed to furnish about 0.01 V-s to aid in the start-up. The current-drive efficiency was much lower (by a factor of 1/5 to 1/6) than previously observed at higher currents in inductively initiated discharges in PLT.

2.2.8 NEUTRAL-BEAM CURRENT DRIVE

To drive current in the core of the plasma, we need a method which will penetrate to the high densities and temperatures. Electron cyclotron heating was considered earlier but discarded because of low predicted efficiencies. Neutral beams of high energy were chosen because of good predicted efficiency and because experiments in TFTR and JET had obtained substantial current with neutral beams.

To obtain the current drive efficiency for neutral beams, we need to know the distribution function for the injected fast ions. Two methods have been used to obtain the distribution function: method (1) a numerical solution of the Fokker-Planck equation on radial zones in a circular approximation to TIBER [21]; this method is sketched in Appendix 2C. Method (2) is an approximate solution to the Fokker-Planck equation by Gaffey [22]. The distribution function is obtained under the assumption that hot-ion (beam) Coulomb collisions only take place with the background Maxwellian ions and electrons, and that the beam speed, v_b , satisfies the inequality $v_i \ll v_b \ll v_e$, where v_a is the thermal speed of species a . Mikkelsen and Singer [23] have integrated the distribution function to obtain an expression for the current-drive efficiency in a circular cross section,

$$\gamma = \frac{85 \langle T_e \rangle}{\ln A_e} \int_0^a \frac{(T_e / \langle T_e \rangle) v_H H J(x, y) F_e r dr}{(n_e / \langle n_e \rangle) v_b a^2}, \quad (2.2-22)$$

where v is the local component of the beam velocity parallel to the B field, H is the fraction of the total beam power deposited at radius r , and $J[x \approx (E_b/10A_b T_e), y \approx 0.8 Z_{eff}/A_b]$ is obtained from a quadrature that must be evaluated numerically. It has a maximum numerical value of about 0.2. An approximation valid for $0 < y < 4$ is

$$J(x,y) = x^2 / [4 + 3y + x^2(x + 1.39 + 0.61 y^{0.7})]. \quad (2.2-23)$$

The factor F_e accounts for the reverse current carried by the electrons in response to drag on the fast ions. It is a function of Z_b/Z_{eff} and $\epsilon = r/R$ in a toroidal device with trapped electrons. Values of F_e have been computed by Start and Cordey [24]. In the absence of trapped electrons, i.e., at the center, $F_e = 1 - Z_b/Z_{eff}$. Unlike wave-drive schemes, as electron trapping increases, the current-drive efficiency increases. For TIBER conditions, an effective $\langle F_e \rangle$ is about 0.76 at $Z_{eff} = 2.0$.

Although we could use Eq. (2.2-22) together with a computation of H to determine γ , we choose instead to use it as a scaling formula for computations carried out with numerical solutions of the Fokker-Planck code described in Appendix 2C. We rewrite Eq. (2.2- 22) as

$$\gamma = f_D \langle T_e \rangle J(\langle x \rangle, \langle y \rangle) F_e / \ln A_e. \quad (2.2-24)$$

Full expressions for the quantities x , y , and F_e can be found in Refs. [22,23]. It should be noted that a number of effects have been incorporated into the "constant" f_D . The neutral-beam energy, aiming and width, and the plasma density and electron temperature profiles are the chief quantities which cause variation of f_D . As long as these quantities are the same between the numerical Fokker-Planck solutions and the zero-dimensional MUMAK computations, this scaling should be valid. For conditions typical of the TIBER baseline we find $f_D \approx 2.65$.

Current-drive by neutral beams has been demonstrated in several experiments: the Culham superconducting Levitron [25,26], DITE [27,28], TFTR [29], and JET [30]. In the Levitron, a hydrogen neutral beam generated from a positive-ion source with mean energy $E_b = 8.5$ keV and current $I_b = 0.2$ A was injected into a $Z_{eff} = 1$ plasma, which had T_e varying from 1 to 4.7 eV. Under these conditions, the electron thermal speed varies from about 1/2 to 1 times

the beam speed. At the lower value, the electrons are moving too slowly to cancel the ion current, and nearly the full ion current was observed. At the other limit, the measured current was found to be opposite to the ion current. In a theoretical paper that treated the electrons without assuming $v_b < v_e$, Cordey, et al.[26] were able to predict the entire dependence of the measured current on T_e , including the observed reversal of current direction at $T_e = 4.7$ eV.

The first application of neutral-beam current drive to a tokamak occurred in DITE [27,28]. Approximately 1 MW of 24-keV neutral-beam power was injected into a plasma with peak density from 3.1 to $7.8 \times 10^{19} \text{ m}^{-3}$ and peak electron temperature from 0.67 to 1.1 keV. There was insufficient power to drive the entire plasma current with neutral beams, thus the reduction in loop voltage was measured while the plasma current was held fixed with the ohmic transformer. The authors verified that co-injection lowered the voltage while counter-injection raised it. Here co- and counter-injection refer to the direction of the toroidal current. Because heating of the electrons by the beam ions could cause a change in the loop voltage, the drop of the loop voltage cannot be taken as definitive evidence that beam-injected ions are carrying plasma current. They verified the existence of hot-ion current by comparing the theoretical loop voltage, as computed from the experimental temperature and density profiles, with the measured loop voltage. The good agreement demonstrated the presence of neutral-beam-driven current. The maximum driven current was 38 kA out of a total plasma current of 80 kA. Although the authors do not report a current drive efficiency, we can estimate it for their case (A) as $(38)(1.17)(0.15)/900 = 0.007$. The authors report that the effects of electron trapping, neoclassical current, and plasma rotation are less than 10% for this experiment.

More recently, results from neutral-beam experiments on TFTR [29] have demonstrated substantial current drive. The driven current is deduced in the same manner as for DITE; the neutral-beam and bootstrap currents are computed theoretically from experimental data. Theoretical and experimental loop voltages are then compared and good agreement is found. Some results from the data analysis are collected in Table 2.2.2. For these cases, in contrast to DITE, neutral-beam power is injected both co and counter to the toroidal current. We note good agreement between measured and predicted loop voltages in Table. 2.2.2.

Table 2.2-2. Results from current-drive experiments on TFTR [29].

Shot	25510	22872	22871	26627	22372	23164
P_{inj} -MW	9.8	11.5	12.5	10.6	11.5	15.6
P_{co}/P_{inj}	0.49	0.50	0.56	0.57	0.61	0.71
I_p -kA	800	900	900	850	900	1000
I_{BS}	290	375	440	350	400	325
I_{NB}	230	225	370	280	370	490
$V_{predict}$	-0.13	-0.10	-0.10	-0.05	-0.23	-0.20
$V_{measured}$	-0.13	-0.15	-0.18	-0.03	-0.20	-0.15

In the analysis of the JET data [30], it is assumed that steady state has been attained and that the inductive field strength, E , and Z_{eff} are constant over the cross section. Since very small bootstrap current is predicted for the flat density profiles in JET, it is possible to deduce the amount of the beam-driven current from the plasma current with the aid of Ohm's law and the computed neoclassical conductivity. Interestingly, the current interpreted from the experiment is somewhat higher than that predicted theoretically.

2.2.9 NEOCLASSICAL (BOOTSTRAP) CURRENT

The existence of a unidirectional toroidal current driven by radial density and temperature gradients in collisionless toroidal plasmas was predicted many years ago by Galeev and Sagdeev [31]. Bickerton, Connor, and Taylor [32] considered the possibility of using this effect to drive a large portion of the plasma current and called such a device a "Bootstrap Tokamak." Until the recent TFTR experiments [29], this bootstrap current had never been observed in a tokamak and was not used in predicting performance of larger devices. The expression for this so-called bootstrap current valid for arbitrary Z_{eff} can be found in Ref. [33]. We will not reproduce it here but quote only the limit for a large-aspect ratio circular tokamak with $Z_{eff} = 1$,

$$j_{BS}(r) = \frac{\epsilon}{B_p} \left[-2.44 \left(\frac{\partial p}{\partial r} - 1.17 n_e \frac{\partial T_i}{\partial r} \right) + 1.75 n_e \frac{\partial T_e}{\partial r} \right], \quad (2.2-25)$$

where p is the pressure of the thermal species and $\epsilon \equiv r / R_0$ is the local inverse aspect ratio. The $Z_{eff} \neq 1$ version of this expression has been used in connection with the numerical computations of the neutral-beam current to determine the fraction of the current from bootstrap effects [21]. Further details are given in Appendix 2C.

We can obtain a simple scaling law by noting that $B_p \propto I_p / a$, $\partial p / \partial r \propto p / a$, and $\epsilon \propto a / R_0$. An additional factor representing the scaling with Z_{eff} for the density and temperature profiles of Section 2.2.5 has been obtained from the calculations in [21]. The scaling expression reads

$$I_{bs} / I_p \approx 0.4 \beta_{pth} A^{-1/2} (1 + 0.94/Z_{eff} + 0.43/Z_{eff}^2), \quad (2.2-26)$$

where β_{pth} is based on the pressure of the thermal components and $A = R_0 / a$.

To increase the contribution of the bootstrap current, we need to increase β_p . For the baseline design, the total β_p is about 1.05, and the thermal component is about 0.8. To see the parameters which affect β_p we make use of Eq. (2.2-7) to write

$$\beta_p = \beta_t B_0^2 / B_p^2 \propto g a B_0 / I_p. \quad (2.2-26)$$

Thus a larger plasma radius, higher toroidal field and smaller plasma current favor larger β_p and larger bootstrap current. There appears, however, to be a limit on the maximum β_p . Both TFTR [29] and ASDEX [5] find the limit for a disruption-free plasma at about $\beta_p \approx 2$.

Figure 2.2.6 shows the projected split of the radial current profile for the TIBER baseline. The neutral-beam and bootstrap currents were computed as described in Appendix 2C. The LH current was scaled from calculations described in Appendix 2B. Although the calculations were not done self-consistently, they should indicate the distribution which might be expected from full calculations.

Many researchers have searched for experimental evidence of the bootstrap current, but only recently has its existence been confirmed in an Octupole [34,35]. It has also been necessary to invoke bootstrap current in order to

Figure 2.2-6. Radial distribution of current for the baseline case.

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explain the TFTR experiments with neutral beam injection [29] and high β_p . As can be seen from Table 2.2-2, a substantial fraction of the current is apparently due to neoclassical effects. There is also strong evidence that bootstrap current is generated in ASDEX in their high- β_p shots [36].

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2.3 MAGNETICS

2.3.1 INTRODUCTION

The magnetic configuration for TIBER-II takes advantage of high-performance superconductor technology (high field and current densities in a radiation environment) coupled with a highly stressed structure to yield a minimum sized confinement system. Coil sizes and placement are based on limiting conductor fields of 12 T in the TF coils and 14 T in the inner PF coils, and peak current densities of 4000 A/cm². The PF coil system is determined with the aid of the FEDC equilibrium code [ref. 1] which performs both fixed and free boundary toroidal equilibria analyses. The plasma size for the TIBER-II reference case is based on a major radius of 3 m, aspect ratio of 3.6, and a toroidal field of 5.55 T.

2.3.2 PLASMA BOUNDARY MODEL AND STABILITY CRITERION

We have adopted the crescent-shaped boundary model² as the preferred shape for the TIBER plasma cross-section. The critical beta scaling for the crescent boundary shape exhibits very close correlation with numerous stability calculations over a wide range of shape parameters and beta values.

The boundary model is:

$$\left[\frac{R - R_0}{a} + \delta \left(\frac{Z}{\kappa a} \right)^2 \right]^2 + \left(\frac{Z}{\kappa a} \right)^{2\lambda} = 1,$$

where

$$R - R_0 = a (\cos \theta - \delta \sqrt{\lambda \sin^2 \theta})$$

and

$$Z = \pm \kappa a \sqrt{\lambda \sin^2 \theta}.$$

The parameter λ is set to 1 for the crescent shape ("sharp-bean") and set to 2 for "round-bean." The quantities R_0 , a , c , and d are evaluated at the plasma midplane as shown in Fig. 2.3-1.

Major radius	R_0
Minor radius	a
Elongation	$\kappa = b/a$
Triangularity	$\delta = c/a$
Indentation	$i = d/a$
Aspect ratio	$A = R_0/a$

Figure 2.3-1. Plasma boundary model.

A comparison of three boundary models--Crescent ($\lambda = 1$), bean ($\lambda = 2$), and the INTOR-type "Dee" boundary^{*} is given in Fig. 2.3-2 for the two cases of mild triangularity (such as TIBER) and strong triangularity. Note that for mild triangularity, the differences in these boundary models is quite small.

The beta scaling against infinite-n ballooning mode instabilities is given in Ref. 1, as follows:

$$B(\%) = 4.7 I_n (1 - b_2 I_n),$$

where $I_n = I_p \text{ [MA]} / a \text{ [m]} B_0 \text{ [T]}$

and

$$b_2 = 0.065 \frac{A}{\sqrt{\kappa}} \left\{ 1 - \frac{(\kappa - 1 + 0.05\delta)}{\sqrt{\kappa}} \left(\frac{1 + 1.5\delta}{\kappa} \right)^{2.5} \right\}.$$

We have chosen a factor of 4/3 between the operating beta and the critical beta.

2.3.3 BASELINE CONFIGURATION

The TIBER-II baseline configuration is based on a plasma current of 10 MA. The size and shape requirements for the plasma are summarized below:

o	Plasma current	I_p	10 MA
o	Toroidal field	B_0	5.55 T
o	Major radius	R_0	3 m
o	Aspect ratio	A	3.6
o	Elongation	κ	2.4
o	Triangularity	δ	0.4
o	Safety factor on-axis	q_0	1
o	Boundary shape		"Crescent"
o	Operating beta	$\langle \beta \rangle$	6%
o	Critical beta	$\langle \beta \rangle_{crit}$	8%

^{*} $R - R_0 = a \cos(\theta + \delta \sin \theta).$

$$Z = \kappa a \sin \theta.$$

Figure 2.3-2. Comparison of boundary models.

These plasma parameters reflect changes from the previous baseline to accommodate: (1) a reduction in the peak field on the TF coil conductor (which required reducing the plasma field from 6-T to 5.55 T), (2) a reduction in the triangularity from 0.55 to 0.40 to allow more space for neutron shielding, and (3) the elongation was increased from 2.2 to 2.4 to regain both safety factor and beta.

The PF coil system is more complex than the previous baseline. We found that more coils were necessary to provide shaping and null-point control over a wide range of plasma parameters. In addition, the number of coils in the central stack was increased for more OH capability. We abandoned the original TIBER concept of "pusher" and "puller" coils as the emphasis shifted from an indented plasma to one with weaker triangularity.

The PF model for equilibrium calculations now uses a total of 24 coil elements (loops), symmetric about the horizontal midplane. The coil positions are given below.

<u>No.</u>	<u>r(m)</u>	<u>z(m)</u>
1	6.20	1.50
2	6.20	3.20
3	4.10	4.25
4	1.80	4.25
5	0.80	3.00
6	"	2.60
7	"	2.20
8	"	1.80
9	"	1.40
10	"	1.00
11	"	0.60
12	"	0.20

The results from the FEDC equilibrium code are shown in Figs. 2.3-3 through 2.3-5.

2.3.4 COIL PARAMETERIZATION FOR EFFI

We use the EFFI code³ for computing the magnetic field within the conductor as well as the magnetic force distribution and inductance matrix. EFFI can model circular loops, arcs, and straight conductor segments having rectangular cross sections with uniform current density. All of the PF coils are co-axial solenoids that can be modeled with the EFFI loop elements, but the TF coils require a more complex model.

Figure 2.3-3. Baseline equilibrium results: configuration.

Figure 2.3-4. Baseline equilibrium results: flux surfaces

Figure 2.3-5. Baseline equilibrium results: profiles.

We have chosen to model TF coils with four arc segments, as illustrated in Fig. 2.3-6. The distance to the inner and outer legs (r_1 and r_2), the half-height (h), and the bend radii (a_1 and a_2) are specified; then the radius a_3 and arc parameters are computed. Fig. 2.3-7 shows a picture of the EFFI model for the TF and PF coils.

Figure 2.3-6. TF coil model for EFFI.

Figure 2.3-7. EFFI model for TF and PF coils.

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2.4 BASELINE OPERATING POINT

2.4.1 INTRODUCTION

Plasma parameters for the baseline operating point of TIBER-II are described in this section. The baseline operating mode is steady state with non-inductive current drive by neutral beam injection (NBI) and lower hybrid (LH) current drive. Bootstrap current is included. The baseline parameters are summarized in Table 2.4-1. The physics assumptions which are used in calculating these parameters are described in Section 2.2. Conservative physics assumptions have been used in the baseline calculations in order to achieve high machine availability for nuclear testing under baseline conditions. Alternate operating points with more optimistic physics assumptions are presented in Section 2.5.

2.4.2 BASELINE POWER BALANCE

The detailed power balance parameters for the baseline operating point are given in Table 2.4-2. All terms in the power balance take into account the radial profiles of the density and temperature by integrating numerically the expressions for the powers using the assumed profiles (see Section 2.2.5).

Kaye-Goldston [1] energy confinement scaling is used with the power degradation term taken to be the sum of the alpha power and the current drive power minus the total radiation power (see Section 2.2.4). The radiation power includes bremsstrahlung and synchrotron radiation. The impurity is assumed to be oxygen with concentration $n_{\text{oxygen}}/(n_D + n_T) = C_{\text{oxygen}} = 0.024$. The concentration of thermal alpha ash is taken to be $C_\alpha = 0.05$. This results in an effective plasma charge $Z_{\text{eff}} = 2.1$. The wall reflectivity for the synchrotron power calculation is assumed to be 90%.

The thermal diffusivity for ions is taken to be equal to that for electrons so that $\tau_{E,e} = \tau_{E,i} = \tau_E$ (see Section 2.2.4). The H-mode energy confinement enhancement factor, defined by $H \equiv \tau_E/\tau_E$ (L-mode), has been selected to be $H = 1.14$. For the baseline beta constraint [0.75 times the ballooning beta limit (see Section 2.2.2)] and the H-mode enhancement factor constraint, $1 \leq H \leq 2$, (see Section 2.2), this value optimizes both the wall loading and plasma Q value (fusion power/total input power) for the operating point. The average neutron wall loading is $\langle \Gamma \rangle = 1.3 \text{ MW/m}^2$, which produces nuclear heating in the inner leg of the toroidal field (TF) coils equal to the

Table 2.4-1. Baseline Parameters.

Parameter	Description	Value
R_0 (m)	Plasma major radius	3.0
a (m)	Plasma minor radius ^a	0.834
κ	Elongation ^a	2.4
δ	Triangularity ^a	0.4
V_p (m ³)	Plasma volume ^a	96
S_p (m ²)	Plasma surface area ^a	170
B_0 (T)	Toroidal field at R_0	5.55
P_{fus} (MW)	Fusion power	314
Q	Fusion power/input power	5.0
$\langle \Gamma \rangle (\hat{\Gamma})$ (MW/m ²)	Average (peak) wall loading	1.3(2.0)
P_{nb} (MW)	Neutral beam power	44
P_{lh} (MW)	Lower hybrid power	19
E_{nb} (MeV)	Neutral beam energy	0.5
I_p (MA)	Total plasma current	10.0
I_{bs} (MA)	Bootstrap current	2.7
I_{lh} (MA)	Lower hybrid driven current	3.9
I_{nb} (MA)	Neutral beam driven current	3.4
$\gamma_{lh}(10^{20}m^{-2}A/W)$	Lower hybrid efficiency	0.47
$\gamma_{nb}(10^{20}m^{-2}A/W)$	Neutral beam efficiency	0.25
$\gamma_{cd}(10^{20}m^{-2}A/W)$	Overall current drive efficiency (NBI + LH + bootstrap currents)	0.52
$\langle T_e \rangle (T_e)$ (keV)	Average (peak) electron temperature	19(25)
$\langle T_i \rangle (T_i)$ (keV)	Average (peak) ion temperature	20(26)
$\langle n_e \rangle (n_e)$ ($10^{20}m^{-3}$)	Average (peak) electron density	1.1(2.1)
$\langle n_{Greenwald} \rangle (10^{20}m^{-3})$	Greenwald density limit	2.57
$\langle n_{Murakami} \rangle (10^{20}m^{-3})$	Murakami density limit	1.35
C_α	Thermal alpha concentration	0.05

Table 2.4-1. Baseline Parameters (continued).

Parameter	Description	Value
C_{oxygen}	Thermal impurity concentration	0.024
Z_{eff}	Plasma effective charge	2.1
τ_E (s)	Gross energy confinement time	0.68

^aValues on the 95% flux surface (see Section 2.2.1).

maximum set by the baseline assumptions for the magnet design including safety factors. The plasma Q value is 5.0 in this case. Higher values of H produce higher plasma Q value at lower wall loading for fixed beta; lower values of H give higher wall loading and lower Q.

The plasma shape is evaluated using the parameters for the 95% flux surface (see Section 2.2.1). The average neutron wall loading is estimated using an average scrape-off thickness $\Delta r_s = 0.117$ m. The model for the first wall assumes that this scrape-off layer is uniform around the plasma periphery, although it is actually thinner on the inside than on the outside (see Section 2.7).

The partitioning of fast beam ion and fast alpha power to bulk electrons and ions during slowing down is calculated using a test particle formalism [2,3]. Following Ref. 2, the equation for the test particle energy, E, assuming $v_i \ll v_{\text{test}} \ll v_e$ where v_i and v_e are the thermal velocities of the background ions and electrons respectively, is

$$dE/dt = -2E [1 + (E_c/E)^{1.5}] / \tau_s, \quad (2.4-1)$$

where the critical energy is (see also Ref. 4)

$$E_c = \left[\frac{3\sqrt{\pi}}{4} \frac{M^{3/2}}{n_e \sqrt{m_e}} \sum_j \left(\frac{n_j Z_j^2}{m_j} \ln A_j \right) \frac{1}{\ln A_e} \right]^{2/3} kTe. \quad (2.4-2)$$

and the slowing-down time is

Table 2.4-2. Baseline power balance parameters.

Parameter	Description	Value
P_{fus} (MW)	Total fusion power	314
$P_{fus,beam}$ (MW)	Beam-bulk fusion power	44
P_α (MW)	Alpha power	63
$P_{\alpha,e}(P_{\alpha,i})$ (MW)	Alpha power to electrons (ions)	42(21)
P_{lh} (MW)	Lower hybrid power	19
P_{nb} (MW)	Neutral beam power	44
$P_{nb,e}(P_{nb,i})$ (MW)	Neutral beam power to electrons (ions)	12(32)
P_{brem} (MW)	Bremsstrahlung power	7.3
P_{sync} (MW)	Synchrotron power lost ^a	2.0
$P_{c,e}$ (MW)	Electron conduction power lost	64
$P_{c,i}$ (MW)	Ion conduction power lost	52
$P_{l,e}$ (MW)	Ion-electron drag power	0.91
τ_E (s)	Gross energy confinement time	0.63
$H \equiv \tau_E/\tau_E(L\text{-mode})$	H-mode energy confinement enhancement	1.14

^aAssuming 90% wall reflectivity

$$\tau_s = \frac{3}{4} \frac{(kTe)^{3/2}}{\sqrt{2\pi} m_e Z^2} \left(\frac{4\pi\epsilon_0}{e^2} \right)^2 \frac{M}{n_e \ln \Lambda_e} \quad (2.4-3)$$

The fraction of the test particle energy given to bulk ions then follows from

$$f_i = 2(v_c^2/v_o^2) \int_0^{x_c} dx \frac{x}{(1+x^3)} \quad (2.4-4)$$

where

$$v_c^2 = 2 E_c/M, \quad v_o^2 = 2E_o/M, \quad x_c = v_o/v_c. \quad (2.4-5)$$

E_0 is the birth energy, and M is the test particle mass. For the baseline operating point, all of the alpha power and neutral beam power is assumed to couple to the bulk plasma electrons and ions. Alpha and hot ion ripple losses have been computed and are found to be negligible [5]. All of the LH power is assumed to couple to the bulk plasma electrons since the fast tail electron slowing-down time is much shorter than the electron energy confinement time.

Finally, the beam hot ion - bulk plasma fusion power is determined using reaction rate parameters calculated by the SIGV code [6]. The fast ion density is obtained from [4]

$$n_{\text{hot}} = I_b \tau_{s,\text{eff}} / V_p , \quad (2.4-6)$$

where the effective slowing down time is

$$\tau_{s,\text{eff}} = (\tau_s/3) \ln [1 + (E_b/E_c)^{1.5}] . \quad (2.4-7)$$

τ_s is the classical slowing-down time, and V_p is the plasma volume. The fast ion pressure, p_{hot} , is calculated from Eq. (2.4-12), and the average fast ion energy is

$$E_{\text{hot}} = 1.5 p_{\text{hot}} / n_{\text{hot}} . \quad (2.4-8)$$

This energy is used in the SIGV calculation to obtain the fusion reaction rate for hot ion-bulk Maxwellian reactions. The fusion power is calculated using average densities; no radial distribution of fast ion density has been calculated.

2.4.3 FAST ION BETA

The finite pressure of fast alphas and fast beam ions is calculated from a slowing-down distribution [2,3]. Following Ref. 2, the distribution function for fast ions, neglecting up-scattering, satisfies the equation

$$\partial/\partial E [g(E) dE/dt] = S(E) , \quad (2.4-9)$$

where dE/dt is given by Eq. 2.4-1. For a mono-energetic and isotropic source,

$$S(E) = S_0 \delta(E - E_0) , \quad (2.4-10)$$

the solution to Eq. 2.4-9 is

$$g(E) = S_0 \tau_s / \{ 2E [1 + (E_c/E)^{1.5}] \} \quad \text{for } E < E_0 \quad (2.4-11)$$

and $g(E > E_0) = 0$. The pressure of the hot ion distribution is obtained from

$$P_{\text{hot}} = MS_0 \tau_s \left(v_c^{2/3} \right) \int_0^{x_c} dx \frac{x^4}{(1 + x^3)} \quad (2.4-12)$$

where the definition of v_c is given in Eq. 2.4-5. The fast ion pressure is included in the MHD beta limit for the baseline operating point.

2.4.4 CURRENT DRIVE PARAMETERS

Detailed parameters for the NBI and LH current drive systems are given in Table 2.4-3. The total non-inductive current requirement is given by $I_{cd} = I_p - I_{bs}$ where I_{bs} is the bootstrap current. The scaling relation for the bootstrap current is given in Section 2.2.9. Neutral beams of sufficient energy to penetrate to the plasma axis are used for core current drive; lower hybrid slow waves are used on the outer flux surfaces. The radius within which the LH waves do not penetrate, r_{max} , is determined by combining the conditions for wave accessibility and avoidance of strong quasi-linear electron Landau damping (see Section 2.2.7).

The partitioning of the total non-inductive current between the NBI and LH systems is done using model current density profiles. The current density profile for the total plasma current, $j_p(r)$, is consistent with the equilibrium solution (see Section 2.3). A model current density profile for the bootstrap current, $j_{bs}(r)$, is obtained from Fokker-Planck analysis [7] (see Appendix 2C). The remaining current density profile, $j_{cd}(r) = j_p(r) - j_{bs}(r)$ is integrated out to the radius of maximum penetration for the LH waves, r_{max} , to give the core current fraction f_c . The neutral beam and lower hybrid currents are then $I_{nb} = f_c I_{cd}$ and $I_{lh} = I_{cd} - I_{nb}$ respectively. The current drive efficiency calculations are described in Sections 2.2.7 and 2.2.8. These models are consistent with results of detailed ray tracing and quasi-linear damping calculations (see Appendix 2B).

Table 2.4-3. Baseline current drive parameters.

Parameter	Description	Value
P_{nb} (MW)	Neutral beam power	43.7
P_{lh} (MW)	Lower hybrid power	18.7
I_{nb} (MA)	Neutral beam driven current	3.43
I_{lh} (MA)	Lower hybrid driven current	3.87
I_{bs} (MA)	Bootstrap current	2.70
E_{nb} (MeV)	Neutral beam energy	0.5
$N_{ ,lh}$	Lower hybrid parallel refractive index	1.60
$N_{ ,ac}$	Accessible parallel refractive index	1.55
$E_{ }/\langle T_e \rangle$	Normalized tail parallel energy	8.0
r_{max} (m)	Radius of maximum penetration (LH)	0.45
$\langle T_e \rangle_{lh}$ (keV)	Average electron temperature (LH)	5.5
$\langle n_{e20} \rangle_{lh2}$ ($10^{20} m^{-3}$)	Average electron density (LH)	0.75
$\omega_{pe}^2/\omega_{ce}$ r_{max}	Dielectric constant at r_{max}	0.35
E_{crit} (keV)	Hot deuterium critical energy	393
n_{hot} ($10^{20} m^{-3}$)	Hot ion density	0.031
E_{hot} (keV)	Hot ion average energy	2.64

2.4.5 FUELING REQUIREMENTS AND ALPHA ASH REMOVAL

The fueling requirements for steady state operation of TIBER are given in Table 2.4-4. The total deuterium (tritium, assuming a 50:50 mixture) current which must be replenished in steady state is 122 A. This includes 18 A of fuel ion burn-up and 104 A of direct particle loss. To estimate particle lifetimes for steady state operation, the alpha particle balance equation is solved for the thermal alpha confinement time, $\tau_{p,\alpha}$, assuming a thermal alpha concentration, $C_\alpha = 0.05$. The thermal fuel ion confinement time, $\tau_{p,i}$, is then assumed to be equal to the thermal alpha confinement time. The neutral beams supply 87 A of deuterium in steady state, assuming the fast ion confinement time is long compared with its slowing down time. The remainder of the required current must be supplied by the fueling system (see Section 2.8).

Table 2.4-4. Baseline Fueling Parameters

Parameter		Description	Value
$I_{c,d}=I_{c,t}$	(A)	Deuterium (tritium) loss current	104
I_{burn}	(A)	Deuterium (tritium) burn-up current	18
$I_{d,beam}$	(A)	Deuterium current from NBI	87
$I_{d,fuel}$	(A)	Deuterium fueling current required	35
$I_{t,fuel}$	(A)	Tritium fueling current required	122
$\tau_{p,\alpha}=\tau_{p,i}$	(s)	Thermal alpha (ion) confinement time	5.9
τ_E	(s)	Thermal energy confinement time	0.68

Note that for the baseline parameters, $\tau_{p,\alpha} \approx 10 \tau_E$ to be consistent with the assumption that $C_\alpha = 0.05$. Experimental results typically indicate that $\tau_{p,i} \approx 3 - 5 \tau_E$. If this were the case for TIBER in steady state, the thermal alpha concentration would be reduced but the fueling requirements would increase.

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2.5 ALTERNATE OPERATING POINTS

2.5.1. INTRODUCTION

Parameters for operating points of TIBER-II, different from the baseline operation described in Section 2.4, are described here. This section also serves as an introduction to the description of pulsed operating scenarios given in Section 2.6. All of the calculations described below were made with the MUMAK time-independent, zero-dimensional power balance code [1] (see Appendix 2A). Results of time dependent calculations, which produce scenarios for evolving the plasma from initiation through current rampup, densification, and heating, and ultimately to the baseline and alternate operating points calculated with MUMAK, are given in Section 2.6.

2.5.1.1 Motivation

The baseline operating mode for TIBER-II is steady state operation with neutral beam injection (NBI), lower hybrid (LH) current drive, and bootstrap current. Other operating modes are of interest, however, especially in the context of comparing the capabilities of the TIBER design with other national designs such as the Next European Torus (NET), the Fusion Engineering Reactor (FER), and the Soviet design (OTR). There are three reasons for exploring alternate operation points for TIBER-II. First, the potential of the design for reaching an ignited operating point must be assessed so that comparisons with other tokamak reactor studies, in which ignition is the driving factor in the design decisions, can be made. Second, the ranges of advanced steady state physics experiments and advanced engineering tests that are possible without making major modifications to the device must be assessed. These results have an impact on plans for the operation of the machine. Finally, results from studies of both pulsed ignited and advanced steady state operating points provide targets for time dependent analysis of TIBER-II, in which the time history of the plasma build-up to the steady state or ignition conditions must be determined.

The general framework for the studies presented in this section is summarized in Table 2.5-1. The scope of parameters for each study is described below, followed, by the results.

Table 2.5-1. Framework for studies of alternating operating points for TIBER-II.

	_____	Effect of toroidal current variations (Section 2.5.2.1)
Pulsed operation		
		_____
		Effect of Energy Confinement Scaling Law Assumptions (Section 2.5.2.2)
		_____
High performance steady state operation		High plasma $Q \equiv$ fusion power/input power (Section 2.5.3)
		_____
		High neutron wall loading (Section 2.5.3)

2.5.1.2 Scope of the Pulsed Operating Points

Pulsed operation, with the goal of achieving ignition, has been studied for the TIBER-II design using a range of toroidal currents and with a variety of energy confinement scaling laws. Cases will be presented using a set of assumptions which are consistent with those used in the CIT [2] studies (see Section 2.2 for discussion of the physics basis for TIBER). The currents used range from 8 to 12.5 MA. In addition, cases at toroidal current of 8 MA have been run with the four energy confinement scaling laws used by the ETR design teams in the U.S. [3], Europe [4], Japan [5], and the Soviet Union [6].

The results are presented in all cases in the form of Plasma Operating CONTOUR (POPCON) [7] plots. These are contour plots of the auxiliary power required to sustain the power balance at points in a density versus temperature space. Density limits, from either the Murakami [8] or Greenwald [9] formulas, and beta limits, from either the Troyon [10] or the Yamazaki [11] formulas (see Section 2.2), tend to produce boundaries for an "operating

window" in density vs temperature space in which the required auxiliary power is less than or equal to zero. Zero auxiliary power requirement means that the plasma power losses are balanced by the alpha heating and the Joule dissipation (ohmic heating) alone. Negative auxiliary power requirement indicates that some additional degradation of confinement or enhanced radiation power loss, which has not been accounted for in the modeling, can be tolerated during operation. This is often referred to as an "ignition margin".

The final step is to estimate the loop voltage for conditions in the operating window. Combining this with the available flux swing (volt-seconds) from the magnet set allows an estimate of the ignited burn time to be made. Combining this with the neutron wall loading during the burn allows estimates of the fluence from ignited operation to be determined. These calculations will be presented in Section 2.6.

2.5.1.3 Scope of the Steady State High Performance Operating Points

High performance steady state operation has been investigated with the assumption that NBI and LH would be used for current drive in the core and edge plasma regions, respectively, as in the baseline operating point. Cases operating at the limits set by physics constraints, without the margin incorporated in the baseline operating point, have been generated. High plasma Q value ($Q \equiv \text{fusion power/total power input to the plasma} \approx 20$) and high peak neutron wall loading $\hat{\Gamma} \approx 3.5 \text{ MW/m}^2$ have been calculated.

High performance operation of TIBER-II with NBI and LH current drive has been achieved at lower toroidal current than in the baseline scenario. Scoping studies have shown that the plasma Q value for NBI + LH scenarios tends to maximize for intermediate values of toroidal current with large bootstrap current fraction. The bootstrap current fraction I_{bs}/I_p , where I_p is the toroidal current, scales linearly with the thermal component of the poloidal beta (see Section 2.2, Eq. (2.2-26) and Appendix 2C). This implies that the toroidal current must be as low as possible (for fixed toroidal beta given by a beta limit) to give low poloidal magnetic field and hence high poloidal beta. However, the energy confinement time and the toroidal beta limit decrease with decreasing current (Section 2.2, [Eqs. 2.2-8 and 2.2-11]), so an intermediate current of 8 MA for TIBER-II seems to produce the best performance.

Other characteristics of the high performance cases include: 1) operation at the full ballooning beta limit given by the Yamazaki formula, 2) full H-mode energy confinement enhancement (factor of two increase in confinement time over Kaye-Goldston L-mode) for the high Q cases, 3) pure deuterium neutral beams with 1 MeV energy, and 4) lower hybrid waves with 90% of the power launched in the direction of the toroidal current. The corresponding baseline assumptions were (see Section 2.4): 1) operation at a more conservative beta limit, 2) partial H-mode enhancement (factor of 1.14 increase in confinement time), 3) 500-keV neutral beams, and 4) LH waves with only 80% of the power launched in the direction of the toroidal current.

The bootstrap current fraction scaling used in these cases includes the dependence on the thermal component of the poloidal beta, the inverse aspect ratio, and the plasma effective charge (see Section 2.2.9, [Eq. 2.2-26]). Finally, the assumption has been made that only 60% of the beam hot ion beta is included in the toroidal beta limit. This value has been used to reflect preliminary theoretical estimates of the bulk plasma stabilization due to hot beam ion finite Larmor radius effects [12].

2.5.2 PULSED OPERATING POINTS

Scoping studies have been made to determine the capability of the TIBER-II design for achieving ignition under various physics assumptions. The calculations are grouped into two study areas: 1) effects of various energy confinement scaling laws, and 2) consequences of operation at various toroidal currents, I_p . The results are presented in the form of POPCON plots. A typical POPCON plot is shown in Fig. 2.5-1. The capability for achieving ignition is evaluated from these plots by the existence and size of an operating window in which the plasma power losses are sustained by the alpha heating and the ohmic heating alone. The operating window boundaries (outlined in Fig. 2.5-1) are the contour of zero auxiliary power, the density limit, and the MHD beta limit.

2.5.2.1 Effect of Toroidal Current Variations

Cases which have been run in the study of toroidal current effects include plasmas with 8, 10, and 12.5 MA toroidal current. The assumptions made in all of these cases are summarized in Table 2.5-2. The energy

confinement scaling law used will be denoted the Kaye-Goldston plus Neo-Alcator (KG + NA) scaling. For this scaling, the gross energy confinement time τ_E is the quadrature sum (as proposed by Goldston [13]) of the Kaye-Goldston time [14], τ_{KG} , and the neo-Alcator time [15], τ_{NA} , (Section 2.2.5, [Eq. 2.2-12]) as

$$\tau_E = [(H_f \tau_{KG})^{-2} + \tau_{NA}^{-2}]^{-1/2} \quad (2.5-1)$$

where H_f is the H-mode enhancement factor taken to be in the range $1 \leq H_f \leq 2$. Cases with full H-mode ($H_f = 2$) are shown in Figs. 2.5-1 - 2.5-3 for TIBER operated with 8, 10, and 12.5 MA, respectively. A case at $I_p = 12.5$ MA using L-mode confinement ($H_f = 1.0$) is shown in Fig. 2.5-4. The density limits shown in these figures correspond to the Murakami formula [8] with leading coefficient 1.5 (hereafter referred to as the Murakami limit) and 0.75 times the Greenwald formula [9] (hereafter referred to as the Greenwald limit). The beta limit which is highlighted in the figures corresponds to the Troyon formula [10] with leading coefficient 2.8%. As expected, the increase in energy confinement and Troyon beta limit as the current increases tends to open the ignition operating window. The following conclusions can be drawn from these results:

- $I_p = 8$ MA, $H_f = 2$: No ignition window exists for densities below the Murakami limit. A narrow window exists for densities below the Greenwald limit for temperatures in the range $7 < T < 12$ keV. The power required to get through the saddle point is 9 MW.
- $I_p = 10$ MA, $H_f = 2$: Ignition is marginal below the Murakami limit, and a window exists only for high temperatures $T > 12$ keV. A substantial ignition window exists for densities below the Greenwald limit for temperatures in the range $6 < T < 15$ keV. The power required to get through the saddle point is 9 MW.

Figure 2.5-1. POPCON for TIBER using KG + NA scaling, $I_p = 8$ MA, $H = 2.0$. The Murakami density limit is $n_{Mu20} = 1.1$; the Greenwald limit is $n_{Gr20} = 2.3$. The beta limit is $\beta_{lim} = 4.8\%$.

Figure 2.5-2. POPCON for TIBER using KG + NA scaling, $I_p = 10 \text{ MA}$, $H = 2.0$. The Murakami density limit is $n_{Mu20} = 1.35$; the Greenwald limit is $n_{Gr20} = 2.9$. The beta limit is $\beta_{lim} = 6.0\%$.

Figure 2.5-3. POPCON for TIBER using KG + NA scaling, $I_p = 12.5 \text{ MA}$, $H_f = 2.9$. The Murakami density limit is $n_{Mu20} = 1.7$; the Greenwald limit is $n_{Gr20} = 3.6$. The beta limit is $\beta_{lim} = 7.6\%$.

Figure 2.5-4. POPCON for TIBER using KG + NA scaling, $I_p = 12.5 \text{ MA}$, $H = 1.0$. The Murakami density limit is $n_{Mu20} = 1.7$; the Greenwald limit is $n_{Gr20} = 3.6$. The beta limit is $\beta_{lim} = 7.6\%$.

Table 2.5-2. Assumptions used for ignition POPCONS with various toroidal currents.

Geometry	$R_0 = 3\text{m}, a = 0.83\text{m}, a_{\text{wall}} = 0.95\text{m}, \kappa = 2.4, \delta = 0.4, B_0 = 5.55 \text{ T}$
Impurities	$n_a/(n_D + n_T) = 0.05, Z_{\text{eff}} = 1.5, \text{Oxygen impurity}, n_D = n_T$
Profiles	$X(r) = X(0)(1-r^2/a^2)^{\alpha_X}, X = n, T; \alpha_n = 0.5, \alpha_T = 1.0$
Safety factor	$q_{\text{NET}} = 5a^2 B_0 K^2 / R_0 I_p; K = \int dl / 2\pi a$
Power degradation	Power P used in the Kaye-Goldston energy confinement scaling $P = P_\alpha + P_{\text{aux}} - P_{\text{brem}} - P_{\text{sync}}$
Neo-alcator scaling	$\tau_{\text{NA}} = 0.071 \langle n_{e20} \rangle R_0^2 (A_p/\pi)^{1/2} q_{\text{NET}}$

$I_p = 12.5 \text{ MA}, H_f = 2$: An ignition window exists below the Murakami limit for temperatures in the range $8 < T < 15 \text{ keV}$. The beta limit is not important in this case. The power required to get through the saddle point is 9 MW.

$H_f = 1$: A narrow ignition window exists below the Greenwald limit for temperatures in the range $8 < T < 10 \text{ keV}$. The power required to get through the saddle point is 18 MW.

These results are summarized in Table 2.5-3.

2.5.2.2 Effect of Energy Confinement Scaling Law

Cases which have been run in the study of energy confinement scaling include calculations with KG + NA (Eq. 2.5-1), Mirnov [16] scaling

$$\tau_{\text{Mir}} = 0.15 I_p (A_p/\pi)^{0.5} \quad (2.5-2)$$

and Mereshkin scaling [17]

Table 2.5-3. Summary of results from ignition studies at various toroidal currents.

Current, k_f	Ignition Murakami	Ignition Greenwald	Beta limited	Power required
8 MA, 2.0	No	Marginal	Yes	9 MW
10 MA, 2.0	No	Yes	Yes	9 MW
12.5 MA, 2.0	Yes	Yes	Marginal	9 MW
12.5 MA, 1.0	No	Marginal	Yes	18 MW

$$\tau_{mer} = 4.43 \times 10^{-3} n_{e20} (a/R_0)^{0.25} R_0^{3.0} q_{NET} A_i^{0.5} T_{e10}^{-0.5} \quad (2.5-3)$$

using the assumptions listed in Table 2.5-2. In Eqs. (2.5-2) and (2.5-3), A_p is the plasma cross sectional area (m^2), and the equivalent atomic mass is taken to be $A_i = A_D A_T / 2 = 3$. Additional calculations have been done with the assumptions given in Table 2.5-2 modified as shown in Table 2.5-4 and using the Asdex H-mode energy confinement scaling [18]

$$\tau_{As} = 0.065 I_p R_0 (A_D + A_T)^{0.5} \quad (2.5-4)$$

Table 2.5-4. Assumptions used to simulate NET design constraints.

Geometry	$R_0 = 3m$, $a = 0.855m$, $a_{wall} = 0.95m$, $\kappa = 2.74$, $\delta = 0.77$, $B_0 = 5.55T$
Impurities	$n_a / (n_D + n_T) = 0.08$, $Z_{eff} = 1.5$, oxygen impurity, $n_D = n_T$
Profiles	$X(r) = X(0)(1-r^2/a^2)^{\alpha_X}$, $X = n, T$; $\alpha_n = 0.3$, $\alpha_T = 0.3$

This last set of calculations shows the performance of TIBER-II under the physics assumptions used to evaluate the ignition capability of NET [14]. In all cases, the toroidal current is $I_p = 8$ MA. Results are shown in Figs. 2.5-5 - 2.5-7. The beta limit outlined in these figures corresponds to the Troyon value with leading coefficient $g = 2.8\%$. Note that the assumption used for NET is $g = 3.5\%$, corresponding to the dashed curve labelled 0.059 in Fig. 2.5-7. The operating window for ignition is substantially larger with this assumption. Conclusions which can be drawn from these figures are the following:

- KG + NA scaling: No ignition window exists below the Murakami density limit. A narrow ignition window exists for densities below the Greenwald limit for temperatures in the range $7 < T < 12$ keV. The power required to get through the saddle point is 9 MW.
- H-mode
- Mereshkin scaling: No ignition window exists below the Murakami limit. A narrow ignition window exists for densities below the Greenwald limit for temperatures in the range $7 < T < 10$ keV. The power required to get through the saddle point is 14 MW.
- H-mode
- Mirnov scaling: No ignition window exists below the Murakami limit. A narrow ignition window exists for densities below the Greenwald limit for temperatures in the range $8 < T < 13$ keV. The power required to get through the saddle point is 11 MW. High temperature, $T > 10$ keV, is required at the saddle point.
- NET assumptions: Ignition is marginal for densities below the Murakami limit at high temperatures, $T > 15$ keV. A narrow ignition window exists for densities below the Greenwald limit for temperatures in the range $9 < T < 17$ keV. The power required to get through the saddle point is 8 MW. High temperature, $T > 15$ keV, is required at the saddle point.
- (with $g = 2.8\%$)

These conclusions are summarized in Table 2.5-5.

Figure 2.5-5. POPCON for TIBER using Mereshkin scaling, $I_p = 8$ MA, $H = 2.0$. The Murakami density limit is $n_{Mu20} = 1.1$; the Greenwald limit is $n_{Gr20} = 2.3$. The beta limit is $\beta_{lim} = 4.8\%$.

Figure 2.5-6. POPCON for TIBER using Mirnov scaling, $I_p = 8$ MA, $H = 1.0$. The Murakami density limit is $n_{Mu20} = 1.1$; the Greenwald limit is $n_{Gr20} = 2.3$. The beta limit is $\beta_{lim} = 4.8\%$.

Figure 2.5-7. POPCON for TIBER using NET assumptions including Asdex H-mode scaling, $I_p = 8$ MA. The Murakami density limit is $n_{Mu20} = 1.1$; the Greenwald limit is $n_{Gr20} = 2.5$. Note that the beta limit outlined ($\beta_{11m} = 4.8\%$) corresponds to the Troyon value with $g = 2.8\%$. The NET assumption, $g = 3.5\%$ corresponds to the dashed curve labelled 0.059.

Table 2.5-5. Summary of results from ignition studies with various energy confinement scaling laws.

Scaling	Ignition Murakami	Ignition Greenwald	Window temperature	Saddle point temperature	Power required
KG + NA, H	No	Marginal	7 - 12 keV	5 keV	9 MW
Mereshkin, H	No	Marginal	7 - 10 keV	5 keV	14 MW
Mirnov	No	Marginal	8 - 13 keV	10 keV	11 MW
NET Assumptions (Asdex H-mode) (T>15 keV)	Marginal	Marginal	9 - 17 keV	13 keV	8 MW

2.5.3 STEADY STATE HIGH PERFORMANCE OPERATING POINTS

The baseline steady state operating point for TIBER-II incorporates conservative physics and engineering assumptions in a number of areas. In this section, alternate steady state operating points are presented, using NBI and LH current drive systems as in the baseline device, in which several of the conservative physics and engineering limits are removed.

The most important of the physics limits is the MHD beta limit. In the baseline scenario the operating beta is taken to be a conservative limit corresponding to the Troyon formula with coefficient $g = 2.8\%$. For the cases below, the full ballooning beta limit given by the Yamazaki formula is used. For TIBER-II parameters $B_0 = 5.55T$, $\kappa = 2.4$, $\delta = 0.4$, using a toroidal current of 8 MA, and the full minor radius to the separatrix $a = 0.855$ m, this means that the operating beta increases from $\beta_{op} = 4.7\%$ (Troyon) to $\beta_{op} = 6.6\%$ (Yamazaki).

Other physics parameters have been changed for the high performance calculations. The plasma effective charge Z_{eff} has been reduced from 2.1 to 1.5 because the scaling of the bootstrap current fraction with Z_{eff} produces higher values as Z_{eff} decreases (see Section 2.2.9). The beam energy has been increased to $E_b = 1$ MeV for higher current drive efficiency. This would require an upgrade of the accelerator technology in the neutral beam system (see Section 3). Finally, it has been assumed that improvements to launching

grill designs for the LH waves could increase the fraction of the LH power launched in the direction of the toroidal current from the 80% assumed in the baseline to 90%. These characteristics are summarized in Table 2.5-6.

Another physics constraint which becomes important when maximum wall loading cases are calculated is the density limit. High wall loading cases with densities up to the full Greenwald value have been explored. In the baseline calculation it is assumed that the density limit is 75% of the Greenwald value. The baseline operating point is actually well below this limit even with the 25% margin because the maximum achievable wall loading (and therefore the required density to achieve this wall loading) is limited by conservative limits on neutron heating in the inner leg of the toroidal field (TF) coils (see Section 3). The TF coil system has been designed to operate with a factor of nearly two higher heat load than is produced by the baseline operating point. High performance operating points have been investigated which produce up to a factor of two higher wall loading ($\langle \Gamma \rangle \approx 2.6 \text{ MW/m}^2$) compared with the baseline ($\langle \Gamma \rangle = 1.3 \text{ MW/m}^2$). In principle these operating points could be achieved without modifications to the coils or shield. With improvements in the forced flow cooling capabilities in the TF coils, even higher wall loading might be possible in TIBER-II.

To identify optimum operating points for both high Q and high neutron wall loading, scans of power balance solutions have been made using values of the H-mode enhancement factor in the range $1.0 < H_f < 2.0$ and using an operating beta fraction normalized to the full Yamazaki ballooning beta limit, $f_b = \beta_{op}/\beta$ (Yamazaki), in the range $0.7 < f_b < 1.1$. The results are given in Fig. 2.5-8. Contours of the plasma Q value (solid lines) and the average neutron wall loading (MW/m^2 - dashed lines) are shown. For fixed beta equal to the ballooning limit ($f_b = 1.0$), the Q value maximizes at $Q = 21$ for $H_f = 2.0$. The wall loading is maximum for $f_b = 1.0$ at $H_f = 1.45$ where $\langle \Gamma \rangle = 2.4 \text{ MW/m}^2$. The peak-to-average wall loading ratio for TIBER II geometry is approximately 1.5, so this corresponds to peak wall loading of $\hat{\Gamma} = 3.6 \text{ MW/m}^2$.

Parameters from the individual cases corresponding to the maximum Q and Γ values in Fig. 2.5-8 are given in Table 2.5-7. These cases have been refined somewhat compared with the corresponding points in Fig. 2.5-8, particularly in

Figure 2.5-8. Contours of Q (solid lines) and $\langle r \rangle$ (MW/m^2 - dashed lines) for TIBER using the high performance assumptions (see Table 2.5-6).

Table 2.5-6. Parameters and physics assumptions used for the high performance cases.

Geometry	$R_0 = 3\text{m}$, $a = 0.855\text{m}$, $a_{\text{wall}} = 0.971\text{m}$, $\kappa = 2.4$, $\delta = 0.4$, $B_0 = 5.55\text{T}$
Impurities	$n_\alpha/(n_D + n_T) = 0.05$, $Z_{\text{eff}} = 1.5$, oxygen impurity, $n_D = n_T$
Profiles	$\chi(r) = \chi(0)(1-r^2/a^2)^{\alpha_\chi}$, $\chi = n, T$; $\alpha_n = 1.17$, $\alpha_T = 0.66$
Energy confinement	Kaye-Goldston, H-mode ($1 < H_f < 2$)
Power degradation	Power P used in the Kaye-Goldston energy confinement scaling $P = P_\alpha + P_{\text{aux}} - P_{\text{brem}} - P_{\text{sync}}$
LH forward launch	90%
Beam energy	1.0 MeV

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the calculation of the radius of maximum penetration for the LH waves. This refinement also affects the partitioning of the neutral beam and lower hybrid driven currents.

2.5.4 SUMMARY

Two classes of alternative operating points have been examined for TIBER-II. Pulsed ignited operating points have been studied with POPCON analysis. High performance steady state operating points have been found which give high plasma Q value or alternately, high neutron wall loading.

The consequences of operation at various toroidal currents have been determined for pulsed scenarios. As expected, the ignition window becomes larger as the current increases because both the density limit and the beta limit increase with current. In addition, the ohmic power contribution to the power balance increases with current (at each point in density vs temperature space), which further widens the operating window for ignition. These results

Table 2.5-7. Plasma parameters for the high-Q and high-wall loading cases.

Parameter	Description	High Q	High Γ
P_{fus} (MW)	Fusion power	340	590
Q	$P_{fus}/\text{input power}$	21.6	13.1
$\langle \Gamma \rangle (\hat{\Gamma})$ (MW/m ²)	Average (peak) wall loading	1.4(2.1)	2.4(3.7)
P_{nb} (MW)	NBI power	11.5	31.3
P_{lh} (MW)	LH power	4.1	13.9
$T_e (T_i)$ (keV)	e ⁻ (ion) temperature	22.2(20.8)	11.4(10.5)
$\langle n_e \rangle (\langle n_{DT} \rangle)$ (10^{20} m^{-3})	e ⁻ (DT) density	0.96(0.80)	2.2(1.88)
$\langle n_{Gr20} \rangle$ (10^{20} m^{-3})	Density limit	2.51	2.51
Z_{eff}	Effective charge	1.50	1.50
$\beta_{fast \alpha}$ (%)	$H_f \tau_E / \tau_E$ (L-mode)	1.32	0.48
H_f	τ_E / τ_E (L-mode)	2.0	1.45
τ_E (s)	Confinement time	1.1	0.70
I_p (MA)	Plasma current	8.0	8.0
I_{nb} (MA)	NBI current	2.0	1.5
I_{lh} (MA)	LH current	1.4	1.0
I_{bs} (MA)	Bootstrap current	4.6	5.5
E_{beam} (MeV)	NBI energy	1.0	1.0
n_{lh}	Refractive index	1.53	2.1
γ_{lh} ($10^{20} \text{ m}^{-2} \text{ A/W}$)	LH efficiency	0.61	0.35
γ_{nb} ($10^{20} \text{ m}^{-2} \text{ A/W}$)	NBI efficiency	0.50	0.31

are summarized in Table 2.5-3. Also, the effects of various energy confinement scaling laws have been studied, including KG + NA, Mereshkin, and Mirnov scalings. Studies using the NET assumptions (including Asdex H-mode scaling) have also been done. Ignition is marginal for all four of these scalings for densities below the Greenwald limit and beta less than the Troyon kink limit with a leading coefficient of 2.8%. The temperatures in the ignition windows and at the saddle points of the POPCON plots are higher for the cases with Mirnov scaling and with the NET assumptions than for the KG + NA and Mereshkin scalings. For fixed inductive volt-seconds available during the burn, higher temperature means longer burn time because the loop voltage decreases with temperature. These results are summarized in Table 2.5-5.

It has been found that with NBI + LH current drive, optimum steady state performance occurs in TIBER-II at an intermediate value of toroidal current, $I_p = 8$ MA. This results from a balance between the performance advantages of high bootstrap current fraction (i.e., low toroidal current, see Section 2.2.9) and high beta and confinement time (i.e., high toroidal current, see Sections 2.2.2 and 2.2.4). Calculations using this current and physics assumptions which are less conservative than those incorporated in the baseline operating point have found operating points with high $Q \approx 20$ and high neutron wall loading $\hat{\Gamma} = 3.7$ MW/m². The most important assumptions used in these scenarios are: 1) operation at the ballooning beta limit and 2) use of 1 MeV neutral beams for good penetration to the axis at higher densities than in the baseline operating point. These high performance points can be sustained in steady state with no more current drive power than is specified in the baseline scenario.

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2.6 PULSED INDUCTIVE OPERATION

2.6.1 INTRODUCTION

As discussed in Section 2.5 above, the baseline TIBER design is a 3-m major radius, steady-state machine with non-inductive current-drive. In this section, we will investigate, as a contingency, pulsed inductive operation of TIBER employing whatever limited volt-seconds capability is available from the machine configuration for the steady-state design. In addition to neutral beam current-drive, the baseline employs 20 MW of lower hybrid (LH) power for current-drive in the outer flux surfaces. Consequently, effective use of this LH power can be made to enhance pulsed inductive performance. Our purpose in this section is to stress the flexibility of the design and demonstrate that TIBER could be chosen as a basis for the prospective ETR design by the expected design qualification date (early 1990s), before we have sufficient time to fully qualify all aspects of steady-state current drive with sufficient reduction of uncertainty. Therefore, besides being capable of full non-inductive, current-driven, steady-state operation, the following pulsed modes of operation will be investigated for the same 3-m major radius machine (i.e., same magnet set and reactor configuration):

1. Inductive ramp-up to an operating plasma current sufficient to sustain a limited flat-top ignited burn for sufficient number of cycles to permit accumulation of a limited neutron fluence.
2. Application of lower hybrid power to supplement the inductive volt-seconds capability and enhance the pulsed mode of operation through:
(a) startup assist, (b) reduction of the plasma loop voltage during burn, and (c) maintenance of the full plasma current at low density during transformer reset.

As might be expected, because our compact 3 m major-radius machine was designed for low cost and aggressive current-driven performance, we shall see that it has a limited volt-seconds capability (~58 V-s--see below). Also, we should emphasize that ignition of itself is not a basic requirement of the

TIBER design rationale. Accordingly, we investigate the pulsed, inductive performance here with the emphasis of contingency operation, and not as a fundamental machine performance requirement.

Figure 2.6-1 illustrates the evolution of TIBER alternate operating modes. Note that as we progress from the contingency modes of inductive, ignited operation and LH-assisted operation, through to the full steady-state current-driven baseline, the overall machine performance and attractiveness (e.g., reactor relevance) increases significantly. We will attempt to quantify this increased performance in terms of end-of-life fluence capability.

We mention here that two other contingency operating modes, namely (1) relaxed shaping and (2) increased plasma dimensions, have not been considered here. In the first case, relaxation of the TIBER triangularity requirement of $\delta = 0.4$ would probably provide more volt-seconds capability for pulsed operation due to relaxed PF coil conditions. However, because of our high elongation, high triangularity is certainly a requisite for maintenance of stability. In the second case, we note that TIBER's compact radial build has already been minimized for the baseline plasma geometry (i.e., $R = 3$ m, $a = 0.83$ m, $\kappa = 2.4$) and could not easily accommodate appreciable changes in plasma dimensions.

2.6.2 MAGNET OPTIONS

Figure 2.6-2 shows a cross-section view of the 3-m major-radius baseline magnet set. The poloidal field (PF) coil set, consisting of the central solenoid stack plus the eight external equilibrium field (EF) coils, represents an upgrade over the original TIBER-II baseline [1] design and is designed to give the desired plasma shape for steady-state operation (i.e., $\kappa = 2.4$, $\delta = 0.4$, $I = 10$ MA) while improving the available volt-seconds available for pulsed inductive operation ($\kappa = 2.4$, $\delta = 0.4$, $I = 8$ MA). Figure 2.6-3 illustrates the PF coil conditions at equilibrium, approximately at the mid-point of the flat top inductive burn. The area of each PF coil cross section in Fig. 2.6-3 is proportional to the current in the coil at that location, and the direction of the current is indicated. The flux surface in Fig. 2.6-2 and the outermost flux surface in Fig. 2.6-3 is that of the 95% flux surface where plasma shaping parameters are defined. The two crosses denote the position of the null points located on the separatrix. We have limited the plasma current in this 3-m pulsed machine to 8 MA where ignition

Figure 2.6-1. Evolution of operating scenarios.

Figure 2.6-2. A cross-section view of the upgraded 3 m major radius magnet set. The principal plasma parameters are indicated on the figure. This represents an upgrade over the original 1986 TIBER design, in that the central solenoid stack has been increased to give more volt-seconds while still providing the desired plasma shape for steady-state operation.

Figure 2.6-3. Equilibrium conditions for the 3 m major radius magnet set for a plasma current of 8 MA at approximately the mid-point of the flat top inductive burn. The area of the cross-section of each PF coil is proportional to the current in the coil at that location.

is achievable but where sufficient volt-seconds are available for limited flat-top burn. At 10 MA (as in the steady-state current-driven baseline), the ignition margin would be higher but virtually no volt-seconds are available for burn due to the larger amount of volt-seconds consumed in starting up to 10 MA (see below).

Table 2.6-1 summarizes the major features of the plasma/magnetic configuration under pulsed, inductive operation.

2.6.3 PF MAGNETICS AND VOLT-SECOND CAPABILITY

The pulsed performance of TIBER is dominated by the volt-seconds capability of the device. This transformer capability is necessarily limited for a compact steady-state machine which did not seek inductive, ignited operation as a fundamental requirement. In this section, we briefly assess the transformer capability as determined by the operating constraints on the PF coil set. Further engineering details may be found in Chapter 3.

Poloidal field scenarios have been developed for several modes of pulsed operation through use of the PFTRADE [2] and PFSCENARIO [3] codes. Options for which scenarios exist include ohmic operation, low-density, LH startup, followed by ohmic burn and LH current-maintenance during transformer reset. Engineering constraints on PF operation were determined from a review of the current and projected superconducting magnet database. This is presented in Chapter 3. The resulting set of physics and engineering constraints which we apply to the TIBER PF set is shown in Table 2.6-2. The poloidal field currents also provide high-beta MHD equilibria at the beginning and end of burn and a broad field null at the beginning of initiation. The coil positions, turns, and peak ampere-turn requirements are shown in Table 2.6-3.

In Table 2.6-4, we summarize the volt-seconds performance for TIBER at a plasma current of 8 MA. For comparison purposes, we show the effect on volt-seconds if a 10 MA plasma current is required. Note from the Table that application of our constraints above requires that, when cocked, the maximum initial flux bias of the PF coil set cannot exceed +22 Wb (i.e., +22 V-s). This is because both the $B_{\max}$ and $JB_{\max}$ constraints are reached at this flux linkage, the former in PF coil #6, the latter in PF coil #5, in the central stack. Similarly, at 8 MA, the minimum value of the final flux bias is -15 Wb due to a stress limit in PF coil #6 near the top of the stack, while the maximum final flux bias allowed is -35.5 Wb due to a maximum field limit of

Table 2.6-1. Features of the plasma/magnetic configuration for pulsed, inductive operation.

Plasma major radius (m)	3.0 ^a
Plasma minor radius (m)	0.834 ^a
Plasma aspect ratio	3.6 ^a
Plasma elongation ^e	2.4 ^a
Plasma triangularity ^e	0.4 ^a
Field on axis (T)	5.55
q (o)	1.0
q _{edge} ^e	5.4 ^b /3.7 ^c
Plasma current (MA)	8 ^b /10 ^c
Beta ^d	0.05 ^b /0.06 ^c
TF coil:	
Mean inner radius (m)	1.312
Mean outer radius (m)	5.50
Mean height (m)	3.55
Pack current density (kA cm ⁻²)	4000
Peak field (T)	12

^aTaken to be the same for both the steady-state baseline and pulsed inductive operation.

^bFor inductive pulsed operation.

^cFor steady-state, current-driven operation for comparison.

^dTaken as 75% of the lower of either Troyon with $g = 0.04$ or Yamazaki (see above).

^eDefined at the 95% flux surface.

Table 2.6-2. Poloidal field system constraints.

<u>Physics</u>	
Startup volt-seconds	48 Wb
Initiation voltage	21 V
<u>Engineering</u>	
$B_{\max}$	14 T
$JB_{\max}$	500 MA-T/m ²
$\sigma_{\text{Tresca membrane}}$	800 MPa
V_{terminal}	20 kV
I_{cond}	25 kA

14 T in PF coil #8 in the middle of the central stack. (We note that at any value of the final flux bias, the MHD equilibrium at a plasma current of 8 MA and at full plasma beta is maintained.) Accordingly, with a maximum initial bias of +22 Wb and a maximum final bias of -35.5 Wb, we obtain a total volt-seconds capability of 57.5 V-s. However, note from Table 2.6-4 that the maximum volt-seconds available during the burn at full equilibrium conditions is only $(-15) - (-35.5) = 20.5$ V-s. The equivalent values at a plasma current of 10 MA are seen to be a total capability of ~59 V-s and burn equilibrium capability of ~9 V-s.

We now note that 20.5 V-s (at 8 MA) and 9 V-s (at 10 MA) is the maximum allowable for burn due to coil constraints; the actual value will be less than this and is determined by the volt-seconds consumed during plasma initiation and ramp-up to full current,

$$\Delta\phi_{\text{startup}} = \int_0^{\tau_{\text{startup}}} L(t)I(t)dt + \int_0^{\tau_{\text{startup}}} I(t)R_p(t)dt$$

where the first and second terms represent the inductive and resistive volt-seconds consumed, respectively. Scaling from the startup scenario for the Compact Ignition Tokamak (CIT), we have determined an approximate requirement

Table 2.6-3. TIBER PF system winding pack dimension.

Coil	R (m)	Z (m)	R ₁ (m)	R ₂ (m)	Z ₁ (m)	Z ₂ (m)	NI (MAT)	n _{turns}
PF1,U	6.2	1.5	6.0	6.4	1.3	1.7	6.5	300
PF1,L	6.2	-1.5	6.0	6.4	-1.7	-1.3	6.5	300
PF2,U	6.2	3.2	6.0	6.4	3.0	3.4	2.536	150
PF2,L	6.2	-3.2	6.0	6.4	-3.4	3.0	2.536	150
PF3,U	4.1	4.25	3.9	4.3	4.05	4.45	2.353	120
PF3,L	4.1	-4.25	3.9	4.3	-4.45	-4.05	2.353	120
PF4,U	1.8	4.25	1.6	2.0	4.05	4.45	5.026	300
PF4,L	1.8	-4.25	1.6	2.0	-4.45	-4.05	5.026	300
PF5+6,U	0.8	2.8	0.6	1.0	2.412	3.188	11.445	600
PF5+6,L	0.8	-2.8	0.6	1.0	-3.188	-2.412	11.445	600
PF7+8,U	0.8	2.0	0.6	1.0	1.613	2.388	9.729	600
PF7+8,L	0.8	-2.0	0.6	1.0	-2.388	-1.613	9.729	600
PF9+10,U	0.8	1.2	0.6	1.0	0.813	1.588	5.88	600
PF9+10,L	0.8	-1.2	0.6	1.0	-1.588	-0.813	5.88	600
PF11+12,U	0.8	0.4	0.6	1.0	0.013	0.788	11.07	600
PF11+12,L	0.8	-0.4	0.6	1.0	-0.788	-0.013	11.07	600

for $\Delta\phi_{\text{startup}}$ of 48 V-s for the 3-m magnet set for pure OH inductive startup; we will use this for subsequent analysis. A more complete analysis with a 0-D time dependent code is currently being conducted (see Section 2.12); preliminary results show a startup requirement for a constant boundary plasma of ~49 V-s, i.e., close to the above scaled value. Use of an expanding boundary scenario, however, can result in several more resistive volt-seconds being consumed during startup--this is currently under investigation. Subtracting 48 V-s from the total available of 57.5 V-s leaves ~9.5 V-s for the equilibrium burn phase at 8 MA.

Note from Table 2.6-4, by contrast, that at 10 MA virtually all of the volt-seconds capability is required for startup, leaving essentially no transformer capability to drive an ignited burn. For this reason, purely

Table 2.6-4. Inductive volt-seconds capability of TIBER.^a

	<u>I = 8 MA</u>	<u>I = 10 MA</u>
Capability		
Maximum initial flux bias (V-s)	+ 22	+ 22
Maximum final flux bias (V-s)	<u>- 35.5</u>	<u>- 37</u>
Total V-s capability (V-s)	57.5	59
Inductive startup		
Volt-seconds consumed during inductive startup (V-s)	48.0	58.5
Volt-seconds available for burn (V-s)	9.5	~ 0
Non-inductive startup		
Minimum final flux bias (V-s)	>- 15	>- 28
Maximum final flux bias (V-s)	<u>- 35.5</u>	<u>- 37</u>
Total equilibrium swing capability	< 20.5	< 9
Transformer reset at low beta		
Minimum final flux bias (V-s)	- 22.3	--
Maximum final flux bias (V-s)	<u>- 32.5</u>	<u>--</u>
Total equilibrium swing capability	10.2	~ 0

^aAt $I \sim 12.5$ MA, there is no flat-top equilibrium swing capability regardless of the startup method.

inductive scenarios for TIBER are investigated at 8 MA only. A plasma current of 10 MA can, therefore, only be utilized in inductive mode given some form of assisted, current-driven operation. As we will see in the next section, application of lower hybrid for startup assist allows us to considerably reduce the inductive volt-seconds consumed during startup and permits fuller use of the equilibrium volt-seconds swing of 20.5 V-s (for the 8 MA case) and 9 V-s (for the 10 MA case) to be applied for the equilibrium burn phase; i.e., the volt-seconds for the equilibrium burn phase are then determined only by coil constraints.

We have also investigated the behavior of the PF coil set at 8 MA for the case of current maintenance by LH during transformer reset; this LH operating scenario is discussed in detail below. The important point to notice from Table 2.6-4 is that under this mode of operation, the plasma beta is essentially zero during reset due to the low density requirements of LH current drive, even though the current is maintained at its full value. Consequently, the PF set is required to work harder to maintain the full current equilibrium during reset due to the absence of plasma diamagnetism. Consequently, whereas 20.5 V-s of full equilibrium swing are potentially available in the forward direction (i.e., full current, full beta), only ~10.2 V-s are available if LH-assisted reset is used due to the constraints of the full current, but low beta, reset. The consequences will be discussed below. We have not assessed the transformer reset capabilities of the 10 MA case. However, we would expect the equilibrium volt-second capability in the case of low beta reset to be very limited, certainly considerably less than the 9 V-s available for the forward scenario.

The contributions of the various PF coils to the plasma volt-seconds are shown in Fig. 2.6-4. The scenario shown here has an artificial burn time of 20 s for display purposes only. The actual burn times are discussed later. The largest individual contribution is from PF1, the outside vertical field coil. The loop voltage provided by the PF coils to the plasma is shown in Fig. 2.6-5. The loop voltage provided to the plasma for initiation is 21 V for 50 ms, which is somewhat aggressive in comparison with the operating experience on tokamaks, not using any rf assist. However, it is felt to be conservative, given the large amounts of rf power available for initiation assistance, if needed. The detailed electrical and mechanical requirements of this scenario are described later in Section 3.

2.6.4 THE BENEFITS OF PULSED OPERATION WITH LOWER HYBRID ASSIST

The TIBER current-driven baseline is configured with approximately 20 MW of lower hybrid (LH) capability for startup purposes and for combination NB/LH or current-drive during steady-state burn. In this section we address the question: to what extent can contingency use of this 20 MW of LH power enhance our pulsed inductive capability?

RF injection at close to the LH frequency has shown potential and flexibility for plasma heating, non-inductive startup, current-drive, and

Figure 2.6-4. Volt-second contribution of PF coils to the 8 MA inductive startup and burn scenario.

Figure 2.6-5. Loop voltage contribution of PF coils to the 8 MA inductive startup and burn scenario.

current maintenance during transformer recharge (see, for example, Ref. 4). Clearly, under full burn conditions ($T \sim 10\text{--}20$ keV, $n_e \sim 10^{20} \text{ m}^{-3}$), current-drive by LH slow waves is restricted to the outer edge of the plasma (i.e., $r/a \gtrsim 0.6$) due to both wave accessibility and quasi-linear electron Landau damping. Accordingly, we can only drive a fraction of the total plasma current with LH during the burn phase. However, utilization of the 20 MW of available LH power can significantly increase the attractiveness of our pulsed inductive operation due to the following factors:

1. LH can be used for current ramp-up, thereby allowing more of the available PF volt-seconds to be applied to the burn phase.
2. LH can drive a fraction $I_{LH}/I \sim 0.1\text{--}0.4$ of the full current in the outer regions of the plasma during the burn, thereby lowering the plasma loop voltage, relative to that for the pure inductive scenario.
3. LH can maintain the full plasma current at low density during transformer recharge. The resulting advantages are:
  - (a) Only one stress cycle per burn cycle rather than two in the pure inductive case. In addition, the maximum stress amplitude of the cycle may be reduced because we are not ramping-up from zero current. The total number of fatigue stress cycles to end of life can, therefore, be increased.
  - (b) Because no current ramp-up is required, we have more equilibrium volt-seconds available for burn. However, as we saw in Table 2.6-4 above, the low beta operation during transformer reset (low density is a necessity for full LH penetration) puts more demand on the PF coil set, with the result that the reverse low-beta swing is constrained to a smaller number of volt seconds. This, therefore, determines the volt-seconds available for burn and not the forward, full beta, volt-seconds capability.
4. Injection of LH at the edge may effectively suppress central sawteeth in the core (see, for example, Ref. 5). Elimination of sawteeth removes the detrimental effects of increasing loop voltage due to flatter profiles and resulting depression of the electron temperature.

The potential advantages of the application of LH are therefore: a) a longer burn-pulse length due to both more volt-seconds available for burn and a lower loop voltage during the burn and (b) a higher number of allowed burn pulses due to the occurrence of only one magnet stress cycle per pulse and lower stress amplitudes. As we shall see, these may both result in an enhanced end-of-life fluence over the pure inductive case.

2.6.5 STRESS ANALYSIS AND CYCLE LIMITATIONS

For the prediction of fatigue life of the PF coil set, we apply the methodology discussed fully in the 1986 TIBER baseline report [1] for crack-growth propagation in the weld metal of the sheath for an initial flaw size of 10% of the sheath thickness. The results for our present sheath thickness of 2.3 mm are summarized in Fig. 2.6-6 where we plot the number of cycles to failure for weld metal as a function of the stress amplitude. Applying the worst case stress amplitude results of 374 MPa for the PF coil at the center of the central stack yields a limit of $\sim 1.04 \times 10^5$ cycles to failure. However, two further restrictions must be put on this cycle limit: First, i.e. apply a factor of 2 safety factor and allow operation to a maximum of only 50% of the number of cycles to failure. Second, we note that, for pure inductive operation, the PF coil set will experience two such stress cycles per burn pulse, i.e., once at the initial cocking of the coils and once at the end of the burn where the PF coil currents are in the opposite direction. Accordingly, we must apply a further factor of 2 reduction in the allowed number of cycles for this type of operation. The net result is a limit of $0.5 \times 0.5 \times 104,000 \sim 26,000$ cycles. As we outlined above, inductive operation with lower hybrid current maintenance during reset may allow this cycle limit to be increased by at least a factor of 2.

We should note here that a more sophisticated analysis which considers the actual time history of PF coil currents (and thus stresses) throughout the PF coil swing would allow us to optimize the performance of the coil set. In particular, because of the exponential dependence of the cycle limit N_{cycles} on stress amplitude (Fig. 2.6-6), reducing the amount of volt seconds $\Delta\phi_{\text{burn}}$ available during the burn, thereby restricting the coil stress amplitudes (i.e., lowering the limits on the permitted coil currents/stresses) can result in a higher product of $N_{\text{cycles}} \times \Delta\phi_{\text{burn}}$. Therefore, because, for a given loop voltage V_{loop} , the burn time per pulse τ_{burn} is proportional to $\Delta\phi_{\text{burn}}$, this

Figure 2.6-6. Number of cycles to failure due to fatigue crack growth in the PF coil conductor sheath as a function of the cyclic stress amplitude. Application of this figure requires a 50% reduction in the number of cycles for a safety allowance and, in the case of pure inductive operation, a further 50% reduction to accommodate two stress cycles per burn cycle.

can result in an optimum product of N_{cycles} and τ_{burn} , and, thus, a maximum neutron fluence at end-of-life. Future analyses of pulsed operation should include such a refinement.

Before concluding this section on cycle limits, we should offer the following important caveat: The cyclic stress analysis tacitly assumes that the PF coils are free standing. However, the latest magnet configuration for TIBER assumes that the inward centering loads of the TF coil set are reacted by the central PF coil stack without the use of additional bucking cylinders. In this case, it is unlikely that the central stack PF coils ever experience significant tension loads with the result that the stress amplitudes per burn cycle may be considerably reduced. This implies that machine cycle limits may be determined by other factors (e.g., first wall thermal cycling fatigue, etc.) and not by pulsed fatigue in the PF set.

2.6.6 IGNITED OPERATION AT 8 MA WITH INDUCTIVE STARTUP

In Table 2.6-5, we show the major parameters for TIBER under ignited operation with pure inductive startup and burn at 8 MA. The ion temperature is 10 keV. In the case of the 10 MA operation we will show that higher ion temperatures yield a higher product of the wall loading and the burn pulse length and thus high end-of-life fluences; however, remember that pure inductive startup and burn is precluded for 10 MA. In the case of the 8 MA operation here in Table 2.6-5, we note that higher ion temperatures result in poor, potentially sub-ignited ignition performance and consequently just show a representative burn at 10 keV. Further details of ignition performance are given in Section 2.5 above.

With the maximum volt-seconds available for burn of 9.5 V-s and a plasma loop voltage of 0.068 V (at $T_i = 10$ keV), we obtain a flat top burn pulse length of ~ 141 s. Accordingly, a peak wall loading of 2 MW/m^2 and a pulse fatigue limit of 2.6×10^4 cycles results in an end-of-life peak fluence of $\sim 0.23 \text{ MW yr/m}^2$ for pure inductive operation in the 3-m baseline machine.

2.6.7 IGNITED BURN AT 10 MA WITH LH STARTUP

In Table 2.6-6, we illustrate ignited burn at 10 MA assuming some form of LH-assisted startup, and assuming we can make full use of the maximum volt-seconds available at equilibrium, i.e., full current, full beta. This assumption may be reduced, however, by the fact that we have to somehow go

Table 2.6-5. Ignited inductive burn at 8 MA; inductive startup.

Magnetic field (T)	5.55
Plasma current (MA)	8
$\langle \beta \rangle$	0.05
Total volt-sec capability (V-s)	57.5
Startup volt-sec required (V-s)	48
Volt-sec for burn (V-s)	9.5
Ion temperature (keV)	10
Peak wall load (MW/m ²)	2.0
Fusion power (MW)	314
Ignition margin	1.12
Loop voltage (V)	0.068
Burn pulse length (s)	141
Max number of cycles	2.6×10^4
Fluence (MW yr/m ²)	0.23

from the low density, full current, ramp-up condition to the full beta burn conditions while conserving volt-seconds. As we discussed above in connection with transformer reset, coil constraints at low beta may similarly reduce volt-seconds at the beginning of burn. Clearly, a self-consistent scenario is required for LH-assisted startup. Note above that ignition at 10 MA was precluded if we required full inductive startup.

In Table 2.6-6, we illustrate the additional effect of varying the ion temperature from $T_i = 10$ keV to 18 keV.

Now the end-of-life fluence can be written:

$$\Phi t \text{ (MW yr/m}^2\text{)} = 3.17 \times 10^{-8} \tau_{\text{burn}} \Gamma_{\text{cycles}} = 3.17 \times 10^{-8} \frac{\Delta \phi_{\text{burn}} \Gamma_{\text{cycle}}}{V_{\text{loop}}}$$

where $\tau_{\text{burn}} = \Delta \phi_{\text{burn}} / V_{\text{loop}} = \text{burn pulse length (s)}$

Table 2.6-6. Ignited burn at 10 MA assuming LH startup.

Magnetic field (T)	5.55	
Plasma current (MA)	10	
Beta	0.06	
Total volt-secs capability (V-s)	59	
Startup volt-secs required (V-s)	~.59	
Equilibrium Volt-secs for LH startup (V-s)	9 (zero for inductive startup)	
<u>Ion temperature (keV)</u>	<u>10</u>	<u>18</u>
Peak wall load (MW/m^{-2})	3.1	2.0
Fusion power	482	309
Ignition margin	1.43	1.18
Loop voltage (V)	0.084	0.036
Burn pulse length (s)	107	251
Fluence (MW yr/m^2)	0.27	0.41

Γ = peak neutron wall load (MW/m^{-2})

$\Delta\phi_{\text{burn}}$ = Volt-seconds available for burn (V-s)

V_{loop} = Plasma loop voltage during burn (V).

As we saw the maximum volt-seconds for 10 MA plasma for burn was $\Delta\phi_{\text{burn}} \sim 9.5$ V-s. Accordingly, from the equation above, once we have determined the product of $\Delta\phi_{\text{burn}}$ and N_{cycles} from consideration of the PF magnet set, we can further optimize the end-of-life fluence by maximizing the product of τ_{burn} and Γ (i.e., Γ/V_{loop}) from consideration of the plasma performance. Reference to Table 2.6-6, as T_i increases from 10 to 18 keV, shows that although the wall loading decreases from 3.1 to 2.0 MW/m^{-2} due to the $\langle\sigma v\rangle/T^2$ cross-section dependence at constant beta, the loop voltage also decreases from 0.084 to 0.036 V, respectively, due to the $T_e^{-3/2}$ dependence of the Spitzer resistivity. Note, therefore, that Γ/V_{loop} (i.e., $\Gamma\tau_{\text{burn}}$) increases from 36.9 $\text{MW/m}^{-2}\text{V}^{-1}$ at 10 keV to 55.6 $\text{MW/m}^{-2}\text{V}^{-1}$ at 18 keV. At $T_i \sim 16$ -18 keV, this effect is

saturating, and ~ 18 keV could be considered an optimum temperature for operation where our peak fluence is ~ 0.4 MW yr/m⁻². Of course, our ignition margin performance decreases with increasing temperature.

2.6.8 PULSED OPERATION WITH LH BURN-ASSIST AND TRANSFORMER RESET

In Table 2.6-7, we illustrate the results of the application of LH for all three phases, i.e., startup-assist, burn-assist, and current-maintenance during transformer reset. Clearly, the benefits of startup-assist in this case are not as relevant as they were for the case in Table 2.6-6 because, once we achieve full operating current, we do not need to startup again. Of course, even though the plasma current remains constant, the fusion power goes to zero during the reset due to low density operation. Therefore, although magnet stress cycling may be alleviated, thermal cycling of first-wall and blanket components continues as in pure inductive operation.

As above, we parameterize the results in terms of ion temperature. Note that we set an upper limit on T_i of 13 keV because, at higher temperatures the average wall load falls below 1 MW/m² (corresponding to the peak falling below ~ 1.54 MW/m²), and we deem this region of operating space to be less interesting for general testing operation. Note also that we have a maximum of 20 MW of LH power available and that this drives successively higher fractions of the total plasma current as the temperature increases due, at constant beta, to the decreasing density.

Now relative to the pure inductive case, the number of allowed stress cycles here are in the range of 50,000, and the burn pulse lengths are considerably longer, both effects being due to the beneficial effects of the LH. Note from the table that, at an operating temperature of around 13 keV, we can achieve an end-of-life fluence of ~ 1.3 MW yr/m⁻², a considerable improvement over the 0.23 MW yr/m⁻² value attainable with pure inductive, ignited operation.

2.6.9 LOWER HYBRID CURRENT DRIVE DURING TRANSFORMER RECHARGE

2.6.9.1 Modeling of Current Ramp-Up and Transformer Recharge

This section reports the design logic used to select the lower hybrid system requirements during low-beta, high-current operation in which lower hybrid current drive (LHCD) alone sustains the toroidal current in the

Table 2.6-7. LH assisted burn at 8 MA; LH startup; LH current maintenance during transformer reset.

Magnetic field (T)	5.55	
Plasma current (MA)	8	
Beta	0.05	
Volt-secs:		
Total forward equilibrium burn ^a (V-s)	< 20.5	
With transformer recharge ^b (V-s)	10.3	
<u>Ion temperature (keV)</u>	<u>10</u>	<u>13</u>
Peak wall load (MW/m ⁻²)	2.0	1.54
Fusion power (MW)	308	242
Q	15.2	11.8
Loop voltage (V)	0.047	0.021
Burn pulse length (s)	216	493
Maximum number of cycles	~52,000	~52,000
Fluence (MW yr/m ²)	0.70	1.25

^aEquilibrium at full beta^bEquilibrium at low beta for reset

presence of a reversed electric field. The philosophy is to use the same LHCD system that is needed to provide surface current density for the baseline, high-beta current drive operating mode. This means the LHCD requirements must be met with ~20 MW of power at 3.7 GHz.

2.6.9.2 Low-Beta Equilibrium Sustained by Lower Hybrid Waves

The calculation of transformer recharge proceeds in two steps. The first, reported here, is to determine the LHCD power requirement in the absence of an electric field ($E = 0$). This is followed by a calculation of power requirements for $E \neq 0$, based on the Karney-Fisch theory. In a calculation of the self-consistent MHD equilibrium generated by LHCD we seek to employ the

appropriate wave power spectrum to generate a stable current profile. For the present we assume the low beta plasma phase has a central electron density and temperature of $n_{eo} = 0.11 \times 10^{20} \text{ m}^{-3}$ and $T_{eo} = 8.0 \text{ keV}$, with a peak ion temperature $T_{io} = 7.6 \text{ keV}$. The density and temperature profiles are

$$n_e = n_{eo} \tilde{\psi}^{0.7}$$

$$T_e = T_{eo} \tilde{\psi}^{0.7}$$

where $\tilde{\psi}$ is the normalized poloidal flux. With a central vacuum toroidal field $B_0 = 5.55 \text{ T}$ this corresponds to a paramagnetic, low beta plasma with $\beta = 8.04 \times 10^{-4}$. The ray tracing calculation is done on the flux surface geometry of this equilibrium in the noncircular boundary with $R_0 = 3.0 \text{ m}$, $A = 3.6$, $\kappa = 2.4$, and $d = 0.4$. Dynamic simulations of the recharging phase [6] show that the total current may be considerably reduced from the value needed for high-beta, full power operation. In our calculation we generate an equilibrium with $I_0 = 8.0 \text{ MA}$, $q_a = 1.0$, and $q_b = 4.92$. The ions are a 50% DT mix, and the effective charge is $Z_{\text{eff}} = 1.33$.

Multiple test rays at different n_e and launch locations, θ , all with $f = 3.7 \text{ GHz}$, were tried to determine the optimum combination. Figure 2.6-7 charts the progress of a ray with $n_e = 1.5$, launched at $\theta = -0.20$, the outboard region well below the midplane. Note that n_e increases along the trajectory, and strong Landau damping occurs at about 70% of the peak density. Ion damping is negligible. Figure 2.6-8 illustrates the current density, $G = \langle j_{\parallel} B \rangle$ generated by a series of rays, $n_e = 2.3, 2.2, \dots, 1.6$, and 1.5 , all launched at $\theta = -0.20$. The highest n_e generates current in the outermost region, peaking at $\tilde{\psi} = 0.3$, while the lowest n_e achieves the deepest current generation, peaked around $\tilde{\psi} \approx 0.6$. Figure 2.6-9(a) plots the fate of the same set of n_e values as in Fig. 2.6-8, except the waves are launched nearer to the midplane, at $\theta = -0.15$. It is possible for the waves to pass all the way through the plasma to reflect off the inboard wall. From the same launch position it is possible to generate current near the edge as well, and this is shown in Fig. 2.6-9(b) for $n_e = 10.0, 6.0, 4.0, 3.5, 3.0, 2.7, 2.6, 2.5$, and 2.4 .

As θ increases it becomes increasingly difficult to drive current near the interior. Figures 2.6-10(a) and -10(b) illustrate the same n_e sequence as in

Figure 2.6-7. (a) LH ray trajectory in low beta plasma during ramp up or transformer recharge: $n_e = 1.5$ and $f = 3.7$ GHz. (b) Ray damping due to electrons; ion damping is negligible.

Figure 2.6-8. Current density, G , and ray paths for spectrum of waves, $2.3 \geq n_{\parallel} \geq 1.5$, launched at the same location, $\theta = -0.20$.

Figure 2.6-9. (a) Ray spectrum launched at $\theta = -0.15$; note current generation near the magnetic axis ($\Psi = 1.0$). (b) Ray spectrum launched at $\theta = -0.15$, but larger $n_{||}$ values than in (a).

Figure 2.6-10. (a) Rays launched at $\theta = 0$; ray paths do not approach magnetic axis. (b) Rays launched at $\theta = 0.15$. For lowest n_r modes electron damping is so weak that waves pass completely through plasma.

Fig. 2.6-9(b), except that $\theta = 0.0$ and $\theta = 0.15$, respectively. The group velocity carries power along paths which do not approach the magnetic axis; in the extreme cases single-pass absorption is not achieved.

After considering the physical constraints on locating the LHCD grill near the bottom of the torus, as well as the weak absorption of rays launched above the midplane, it would appear most useful to locate the grill approximately spanning $-0.15 \lesssim \theta \lesssim 0$. The final, self-consistent LHCD generated MHD equilibrium thus was developed with a set of modes at $\theta = 0, -0.05, -0.10$, and -0.15 with values of $n_{\parallel}$ in the range 4.1 to 1.5. Ray tracing is done in the initial flux surface geometry, and $G = \langle j_{\parallel} B \rangle$ is found by summing contributions from all rays. From $G(\theta)$ the diamagnetism, $F(\theta) = RB_{\theta}$, is computed, and an updated equilibrium is then found from the Grad-Shafranov equation. Ray tracing and current drive is computed anew and $F(\theta)$ is updated. The process is iterated until the equilibrium converges [7]. The code optimizes the power spectrum to reproduce the original, desired equilibrium. The resulting, converged equilibrium is shown in Fig. 2.6-11, along with the seven rays' contributions to G . The final equilibrium has $\beta = 8.06 \times 10^{-4}$, $I_0 = 8.13$ MA, $q_a = 1.01$, and $q_b = 4.89$. It is paramagnetic and has the magnetic axis slightly shifted, to $R_{\text{mag}} = 3.02$ m. The total power distributed over the spectrum is $P_{\text{CD}} = 6.90$ MW, and, for $\bar{n}_e = 0.06 \times 10^{20} \text{ m}^{-3}$, this represents an LHCD efficiency of $\gamma = \bar{n}_e R_0 I_0 / P_{\text{CD}} = 0.23$.

2.6.9.3 Experimental Expectation for LHCD Efficiency

The calculation above predicts a centrally peaked current profile with a reasonably high γ value. It is encouraging that, in fact, similar results have been demonstrated in the laboratory. Figure 2.6-12 shows an empirical trend for γ to increase with T_{e0} . In fact, the highest reported γ values, ~ 0.15 to 0.30 , have been achieved on JT-60, which is prototypical of the low beta phase of TIBER II. JT-60 has $R_0 = 3.1$ m, $a = 0.9$ m and launches ~ 2 GHz waves with $n_{\parallel} \sim 1.7$; with operation at $\bar{n}_e = 0.09 \times 10^{20} \text{ m}^{-3}$ and auxiliary heating to raise $T_{e0} \gtrsim 5$ keV [8], values of $\gamma \sim 0.3$ are reported. Under the same conditions partial transformer recharge was also achieved on JT-60.

Increasing confidence that LHCD ramp up-and-recharging are credible comes from various tokamak experiments. Table 2.6-8 cites several demonstrations along with the respective low density and temperatures which characterized these experiments. All these experiments report $\gamma \sim 0.1$ when $T_{e0} \lesssim 2$ keV in

Figure 2.6-11. Converged, self-consistent MHD/LHCD equilibrium with $\beta = 8 \times 10^{-4}$, $I = 8.1$ MA; total LH power is 6.9 MW, distributed over $4.1 \leq 1.5$ $\lesssim n_n \lesssim 1.5$ and several poloidal locations. This plasma is typical of transformer recharge, $n \approx 6 \times 10^{18} \text{ m}^{-3}$ and $T_{eo} = 8$ KeV, with $\gamma = 0.23$.

Figure 2.6-12. Experimental values of γ vs T_{eo} for LHCD. Note, empirical value of γ refers to launched power rather than absorbed power.

the absence of an electric field. Also listed in the table is the normalized runaway speed, $v_r = \omega/k_{\parallel} v_r$, where $v_r = 9.33 \times 10^{-3} \sqrt{n_e/|E|}$ in SI units. The ratio u_r must be ≤ 1.5 to avoid reverse runaway generation [9]. The effect of the reversed electric field is detailed in the next section.

2.6.9.4 Current Drive Efficiency with Non-Zero Electric Field

Based on the presentations and discussions at the INTOR-Related Specialists' Workshop on Noninductive Current Drive (IPP-Garching, September 15-17, 1986), there appears to be a consensus that the Karney-Fisch theory [9] is a good model which reproduces the features of current ramp-up and transformer recharging seen in various experiments. For the purposes of scoping studies on TIBER II we can provide the following simplified algorithms to model current drive when $E \neq 0$.

The total toroidal current in the plasma is

$$I = (V_L/R_{Sp}) + I_{CD},$$

where the loop voltage is $V_L = 2\pi R_0 E$, and R_{Sp} is the Spitzer toroidal resistance. (Presumably a more detailed theory would modify R_{Sp} to include neoclassical effects.) The noninductive current, I_{CD} , is related to the value where $E = 0$ by the following approximate formula:

$$I_{CD} = \gamma \frac{P_{CD}}{\bar{n}_e R_0} \times N(u_r) .$$

Here, $\gamma P_{CD} \sqrt{n_e} R_0$ is the current driven by an absorbed power, P_{CD} , when $E = 0$; the value of γ has been discussed above. The factor N enhances the noninductive current with a forward emf ($u_r < 0$), such as occurs during current ramp-down, and N reduces the current with a reverse emf ($u_r > 0$) during ramp-up. The argument of $N(u_r)$ is the ratio of wave phase speed to the runaway speed,

$$u_r = (c/n_{\parallel})/v_r ,$$

Table 2.6.8. Recharging experiments.

Experiment	$\bar{n}_e (10^{20} \text{ m}^{-3})$	$T_{eo} \text{ (keV)}$	u_r
Alcator	0.10 - 0.30	1.8	0 - 1.0
PLT	0.02 - 0.06	2	0 - 1.5
ASDEX	0.01 - 0.06	2.5	--
JT-60	0.04	2	--

where n_u includes any upshift in the plasma. The function N is displayed in Fig. 2.6-13 for $Z_{eff} = 5$. For our formulation to be valid there must be no generation or retention of runaway electrons, which implies the following constraint:

$$-1 \leq u_r \leq 1.5 .$$

Since u_r varies spatially throughout the plasma, it is problematic how to evaluate u_r or N . One possibility is to crudely evaluate u_r by using $n_u = 2.5$, putting n_e equal to the average value $\bar{n}_e$, and setting E to the value given at the major radius R_0 . The function N is only mildly dependent on Z_{eff} , but one should recall that γ has a Z_{eff} dependence.

Finally, it should be emphasized that the Karney-Fisch theory assumes waves at high phase speeds, $c/n_{||}/v_e \gg 1$, where v_e is the electron thermal speed, and this constraint is usually satisfied at low T_e ($\lesssim 2$ keV). If somehow slower phase speeds were used ($c/n_{||}/v_e \simeq 1$), a different formulation of the electric field effects would be needed, as, for example, given by V. Chan, et al. [10].

2.6.9.5 Current Drive Power Requirement

In calculating the reverse loop voltage during the recharge phase, we use the relation $V_L = -\Delta\phi/t_{dw}$, here $\Delta\phi$ is the flux swing in the transformer and t_{dw} is the recharging dwell period. Combining our expressions we obtain the general formula for LHCD power required for recharge:

Figure 2.6-13. Conductivity function in Karney-Fisch theory, for ion charge $Z = 5$. Abscissa is phase speed normalized to runaway speed.

$$P_{CD} = \left| \frac{\bar{n}_e R_0}{\gamma N} I + \frac{|\Delta\phi|}{R_{Sp} t_{dw}} \right|.$$

Generally P_{CD} is a function of six design variables, $\bar{n}_e$, I , $\Delta\phi$, and t_{dw} , and through their effect on γ and R_{Sp} , T_{eo} and Z_{eff} . The complete optimization in this parameter space is beyond the scope of our present study and would require a knowledge of cost details of various reactor subsystems (e.g., power supplies). However, we can achieve short recharge periods (and high reactor duty factor) with moderate LHCD power, as we shall show. To simplify matters, we assume recharge occurs with $T_{eo} = 8$ keV at $n_{eo} = 0.07 \times 10^{20} \text{ m}^{-3}$ and $Z_{eff} = 3$. Dynamic (transport) code calculations suggest these plasma parameters are consistent with $P_{CD} \simeq 20$ MW in the TIBER II device. If we prescribe the plasma resistance as $R_{Sp} = (0.103 \mu\Omega) Z_{eff} \bar{T}_e (\text{keV})^{-3/2}$, we find $R_{Sp} = 30 \mu\Omega$. Likewise, we decrease the γ value calculated in Section 2.6.9.2 (at $Z_{eff} = 1.33$) by the scaling factor [7] which accounts for the higher Z_{eff} , obtaining $\gamma = 0.157$. Also, we fix $I = 8.0 \text{ M}$ and specify $\Delta\phi = 20.5 \text{ V-s}$. The reference value of $\Delta\phi$ is actually 10.2 V-s (see Table 2.6-4), and the 20.5 V-s corresponds to an earlier design. What remains at this point is a simple relationship between P_{CD} and the dwell period. The argument of $N(u_r)$ is calculated from the electric field given by $E = (\Delta\phi/t_{dw} 2\pi R_0)$, and $n_* = 2.5$ is assumed, based on an analysis of the ray history in the low-beta MHD equilibrium.

Figure 2.6-14 illustrates for our assumed parameters that recharging with $P_{CD} = 20$ MW requires only 27 s. Since $u_r < 1.5$ there is little danger of generating reverse runaways. However, it may be more desirable to increase t_{dw} to ~ 80 to 100 s, which would reduce P_{CD} and the power supply rating needed to recharge the transformer. Note that any dwell period less than ~ 100 s will still result in a high duty factor since the burn period may exceed 800 s, especially with LHCD assist to the transformer. Operation with $t_{dw} \simeq 100$ s provides an additional safety margin. For example, if recharging requires higher $\bar{n}_e$ or if γ is less than our prediction, it would still be possible to recharge with $P_{CD} \leq 20$ MW at $t_{dw} \simeq 100$ s, provided the ratio $\bar{n}_e/\gamma$ has not increased by more than 50% over the values used in our calculation. On the other hand, if $\bar{n}_e/\gamma$ increases by a factor of 100% or more, then P_{CD} would exceed 20 MW, even for very long t_{dw} values.

Figure 2.6-14. LHCD power needed to recharge transformer vs recharge dwell period. As long as $u_r \leq 1.5$ reverse runaways are insignificant.

2.6.9.6 Conclusions on LH Recharging

At low beta (but $T_{eo} \approx 8$ keV) the efficiency γ for LHCD is adequate to allow transformer recharge at $P_{CD} \lesssim 20$ MW in a dwell period $t_{dw} < 100$ s for a reset flux of $\Delta\phi = 20.5$ V-s. There is a comfortable safety margin which encourages us that the technique will work even with somewhat degraded LHCD performance. Further reductions in P_{CD} or in t_{dw} are likely if continued studies of Z_{eff} and T_{eo} variations are carried out. The ultimate choice among t_{dw} and P_{CD} options will depend on additional considerations such as power supply costs and magnet fatigue.

2.6.10 CONCLUSIONS

In Table 2.6-9, we illustrate the conclusions from this section. We have directed our attention towards volt-seconds capability (the driving constraint on pulsed operation) and resulting end-of-life fluence (an important figure-of-merit). Clearly, as we suggested previously in Table 2.6-1, the transition from basic pulsed inductive ignited operation, through LH-assisted operation to full steady-state current-drive, brings increasing benefits in terms of larger burn times, higher fluences, reduced stress cycling, thermal burn stability and, particularly, increased reactor attractiveness.

Given this brief explorative study of the application of LH to pulsed operation, it appears, in principle, to have several potential advantages over pure inductive, ignited operation if used for startup, burn-assist, and transformer reset. We should, however, note that the use of LH transformer-reset alone with a conventional ignited burn in the forward direction may not result in a tangible improvement in volt-seconds for burn over pure inductive operation; this is because the burn volt-seconds in both cases may be limited to about the same value due to PF constraints under low beta in the first case and inductive startup requirements in the second (for example, from Table 2.6-4, the values for TIBER are ~ 10.2 and 9.5 V-s, respectively). Of course, the benefits of burn-assist and stress cycle reduction are additional potential advantages. In the analysis above, we have indicated that an improvement in end-of-life fluence by a factor of ~ 5 may be possible (the improvement is ~ 20 for full steady-state operation). It is our opinion at present that the "jury is still out" on the issue of LH, relative to a pure inductive ignited scenario. This is certainly a fruitful area for further research.

Table 2.6-9. A summary of TIBER operating modes.

	Ignited burn; inductive startup		Ignited burn; LH startup		LH-assisted burn, LH transformer reset		Steady-state current drive (TIBER base- line)
	<u>8 MA</u>	<u>10 MA</u>	<u>8 MA</u>	<u>10 MA</u>	<u>8 MA</u>	<u>10 MA</u>	
Volt-seconds for burn (V-s)	9.5	~0	20.5	9	10.2	~0	NA
Loop voltage (V)	0.068		0.068	0.036	0.021		NA
Burn time (s)	141		301	251	493		Steady-state
Startup or reset time (s)	~10-20		≥ 50	≥ 50	~100		NA
Peak wall load (MW/m ²)	2.0		2.0	2.0	1.54		2.0
N _{cycles}	26,000		26,000	~26,000	>52,000		NA
Fluence (MW yr/m ²)	0.23	~0	0.49	0.41	1.25	~0	5

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2.7 EDGE PHYSICS AND IMPURITY CONTROL

2.7.1 INTRODUCTION

The evolution of poloidal divertor designs [1] is shown in Fig. 2.7-1, starting with an early reactor concept and leading up to the INTOR design. In the early design the diverted flux channels are large and result in a large system. Experiments and analyses on ASDEX, PDX, and D-III led to a better understanding of the conditions that a divertor must produce to allow good confinement of the hot plasma. In particular, it was found that the divertor configuration and the divertor plasma must prevent the penetration of gas into the confined plasma region. Fig. 2.7-2 shows the DIII-D [2] and TIBER II configurations on roughly the same scale to allow a comparison of divertor dimensions. The most recent designs, including NET [3] and CIT [4], use open configurations and rely on a dense edge plasma to intercept the gas produced at the neutralizer plate. Experiments on ASDEX, PDX, D-III, and recently on DIII-D have shown that a dense edge plasma can be produced by re-ionizing most of the gas produced at the plate. This high recycling mode allows a large flux of particles to carry the heat to the plate so that the mean energy per particle can be low. Sputter erosion of the plate can be greatly reduced if the average impact energy of the ions at the plate can be reduced to near or below the threshold for sputtering of the plate material.

A divertor configuration similar to that in DIII-D has been adopted for TIBER. The main reason for the change from our original vented-plate concept was that the vented plate required a high neutral pressure on the plasma side in order to pump helium. Hydrogen would be pumped because of the reflux of charge exchange neutrals. The present configuration allows part of the flux of edge plasma ions to be neutralized at the entrance to the pumping duct so that helium is pumped as well as hydrogen. In this configuration, an effort was made to preserve the insensitivity to plasma motion and preserve as much as possible the compactness of the vented plate. The new configuration shown in Fig. 2.7-3 has the following features:

- Peak heat loads of about 3 MW/m^2 , and less than 4.5 MW/m^2 , even if the plasma shifts position by as much as 0.1 m in any direction.
- High gas recycling to reduce sputtering, erosion of the plates, and production of impurity ions. $n_t > 10^{20}/\text{m}^3$, $T_t < 10 \text{ eV}$ near the plates.

Figure 2.7-1. Evolution of poloidal divertor designs from early reactor concepts to INTOR divertor; from Ref. 1. Early divertors were large and required large magnets.

Figure 2.7-2. DIII-D and TIBER are shown on the same scale for comparison of the divertors.

Figure 2.7-3. The TIBER divertor configuration showing magnetic flux surfaces in the edge plasma, for coil configuration k67e, and 1 of 16 vacuum pipes.

- Small vacuum ducts made possible by the high pressure in the ducts due to the high recycling. Pumping speed is $S = 50 \text{ m}^3/\text{s}$, and pressure in the vacuum ducts is $p \approx 40 \text{ mTorr}$.
- At least 0.48 m of neutron shielding between plasma and the superconducting magnets.

2.7.2 THE DIVERTOR CONFIGURATION

The divertor configuration shown in Fig. 2.7-3 attempts to satisfy TIBER II's design goal of ignition in a minimum-sized device and to maintain the required spacing (about 0.5 m) for neutron shielding everywhere between the plasma and the superconducting magnets. For the divertor to act as a divertor and not as a limiter, it must be located outside the magnetic separatrix where it intercepts no closed field lines. The minimum distance that can be allowed between the neutralizer plates and the closed field lines is determined by the need to prevent gas and impurities from reaching the confined plasma. With high recycling at the plates, the plasma temperature and density near the plates will be about 10 eV and 10^{20} m^{-3} . In such a plasma, the mean free paths, mfp, for the recycled particles are short, with the longest being $\lambda_{01} \approx 20 \text{ mm}$ for the fast neutrals that result from charge exchange CX. Because of the large charge exchange cross section, a recycled D^0 or T^0 atom can CX several times before finally being ionized. These multiple CXs allow deeper penetration than the mfp would suggest. In TIBER's divertor configuration the minimum space between a plate and the confined plasma is 150 mm. This is much greater than the mfp for any neutral atom, when the recycle rate is high. Because the poloidal field decreases toward the null, the thickness of the SOL is greatest (see Fig. 2.7-3) and the power flux is least near the nulls. For this reason, we located the divertor plates as near to the null points as the mfps allowed.

Identical top and bottom divertors are each divided into 16 modules so that each module can be removed between 2 of the 16 TF coils. Each lower module contains an inner, a lower, and an outer divertor plate; upper modules are identical. This description is of a lower module. There is a slab-shaped vacuum duct under each lower plate, and a circular vacuum pipe connects the duct in each module to an external pumping manifold. The physics of the vacuum system is discussed in Section 2.7.10, and the design of the vacuum system is described in Section 3.1.4.0.

TIBER's first wall is located on a magnetic flux surface to minimize the particle and heat load to the wall. The separation of the first wall from the confined plasma is several radial decay lengths for both particle and power flow. The decay length for energy will be shown in Section 2.7.4 to be $\Delta_E \approx 15$ mm at the midplane in both the outer and inner scrape-off layers (SOLs). The decay length for particle flow will be shown in Section 2.7.3 to be about $\Delta_N \approx 20$ mm in the outer, and $\Delta_N \approx 17$ mm in the inner SOL. Particle fluxes and heat fluxes to the inner and outer walls are discussed in Section 2.7.5.

2.7.3 THE RADIAL SCALE LENGTH FOR PARTICLES IN THE EDGE PLASMA

The decay length for particle flow has been shown [5] to scale as $\sqrt{(L/B)}$, as would result from Bohm diffusion. Bohm cross-field diffusion $D \text{ (m}^2/\text{s)} = T_e \text{ (eV)}/16B \text{ (T)}$ results in the following expression for the e-fold length, Δ_N , for particles at the midplane and near the separatrix:

$$\Delta_N^2 \text{ (mm)} = 7.14 \cdot \frac{L \text{ (m)}}{B \text{ (T)}} \cdot \sqrt{T_e \text{ (eV)}} .$$

Here, L is half the connection length along a field line from divertor to divertor. Figure 2.7-4 taken from Ref. [5] shows some experimentally measured particle e-fold lengths, Δ_N , plotted against $\sqrt{(L/B)}$. We calculate $L \approx 26$ m and $L \approx 38$ m for TIBER's outer and inner scrape-off layer (SOL), respectively, by integrating along the field lines that start at $\Delta_E \cdot \ln(2) = 10.4$ mm out into the SOL at the midplane. These field lines are located midway out in the power flow. (Note that L becomes infinite on the separatrix because the poloidal field vanishes at a null point.) Table 2.7-1 lists the parameter values in each SOL which result in $\Delta_N \approx 20$ mm in the outer SOL and $\Delta_N \approx 17$ mm in the inner SOL, assuming Bohm diffusion.

The electron temperatures, T_e , shown in Table 2.7-1 were calculated by assuming that the power transport along the field direction is by classical thermal electron conduction all along the length, L , of a magnetic flux tube, whose cross section varies along the length to preserve magnetic flux. The T_e in Table 2.7-1 are for points located at the midplane and at distances of $\Delta_E \cdot \ln(2)$ out into the SOLs.

Figure 2.7-4. Comparison of experimentally measured density scrape-off lengths, Δ_N , with the theoretical value $\Delta_{Bohm} = (2D_{Bohm} L_c / C_s)^{1/2}$, where $D_{Bohm} = 0.06T[eV]/B[T]$, $[m^2/s]$. From P.C. Stangeby [5].

Table 2.7-1. Parameters used to calculate the radial, particle scale lengths in the outer and inner scrape-off layers (SOLs). L is half the connection length, B is the toroidal magnetic field, T_e is the average electron temperature, and Δ_N is the e-fold length for density in the SOL.

	$L(m)$	$B(T)$	$T_e(eV)$	$\Delta_N(mm)$
Outer	26	4.35	80	19.5
Inner	38	7.8	76	17.4

2.7.4 THE RADIAL SCALE LENGTH FOR ENERGY IN THE EDGE PLASMA

The heat load on the divertor plates is determined by three factors in addition to total power: The radial scale length for power flow in the edge plasma, the magnetic mapping of the power flux onto the divertor plates, and the angle that the plates make with the flow direction. The calculation of the radial scale length for power flow in the SOL is presented here, and the results indicate thicknesses of about 1.5 cm at the midplane in both outer and inner SOL when the radial thermal diffusivity has the expected value of $\chi_{\perp} \approx 4 \text{ m}^2/\text{s}$.

In Ref. [6], P.J. Harbour describes a simple method of estimating the average radial decay length for power density in the SOL. In that model, the time for the radial transport of heat is obtained from the ratio of stored energy to the power flow in the SOL. Because this gives the average radial heat transport, the thickness so derived is also an average value. In Ref. [7], P.C. Stangeby assumes uniform heat flow across the separatrix and obtains an expression for the thickness in terms of the connection length.

Here, we use a model that attempts to include the variation of thickness caused by variation in the magnetic field strength along a flux tube. Figure 2.7-3 shows the large variation in thickness of the tubes of magnetic flux in TIBER's SOLs. The primary simplifying assumption made here is that the plasma pressure is constant on a magnetic flux surface. This should be a good assumption everywhere except near the divertor plates where a pressure gradient develops to accelerate the new, recycled ions. The constant- p assumption has the effect of reducing the radial gradient in energy density near the null points where the poloidal field is weak and the flux tubes are

thick, as can be seen in Fig. 2.7-3. As a result, the average length for power flow in the SOL is increased because a larger fraction of the power enters the SOL near the midplane, and less enters near the divertor.

Radial heat conduction across the separatrix and into the SOL can be written as

$$Q_{\perp} = -n\chi_1 \frac{d(kT)}{dr} \approx \chi_1 nkT/\Delta_E,$$

where Δ_E is the radial scale length for power flow in the SOL. The conservation of magnetic flux requires that $2\pi R \Delta_E B_p$ remain constant along a flux tube (only the poloidal component B_p of the total field B produces flux through the area $2\pi R \Delta_E$) and, in particular:

$$2\pi R \Delta_E B_p = 2\pi R_o \Delta_{EO} B_{p0} = A_{\parallel} B = A_{\parallel t} B_t, \quad (2.7-1)$$

where subscript o refers to values at the midplane, and subscript t refers to values at the divertor target. The area $A_{\parallel}$ is the cross sectional area for power flow in the direction of total field B . We now assume that the plasma pressure is constant along a line of flux between the null points, and set $nkT = n_o kT_o$. (Outside of the null points the pressure decreases near the sheaths at the divertors.) Then, substituting for Δ_E from Eq. (1) gives:

$$Q = \chi_1 \left(\frac{n_o kT_o}{\Delta_{EO}} \right) \left(\frac{R B_p}{R_o B_{p0}} \right).$$

We obtain the total power flow P_{SOL} into the outer or inner SOL by integrating Q , over the part of the plasma surface that faces that SOL:

$$P_{SOL} = \chi_1 \frac{n_o kT_o}{\Delta_{EO}} 2\pi \int_{null-}^{null+} \frac{R \cdot B_p}{R_o \cdot B_{p0}} \cdot R \cdot ds_p,$$

where s_p is measured along the separatrix in the poloidal plane and the integral runs from the lower null point to the upper one (see Fig. 2.7-3). In steady state, half of this power must flow out across the sheath at either end of the SOL. Power balance at either sheath then gives

$$\gamma^c St n_t kT_{et} A_{\parallel t} = (1/2) P_{SOL}.$$

Here, $c_s \approx \sqrt{2kT_e/M_i}$ is the ion sound speed, $\gamma \approx 8$ is the sheath parameter, and subscript t indicates values near the target plate. γ gives the energy deposited on the divertor plate per incident ion, in units of kT_t . Although γ includes the presheath ion energy, the ion acceleration at the sheath, the electron energy, and recombination energy, γ , are relatively small here because of the rather high reflection coefficient for ion energy when energy is low as it is here (see Section 2.7.10.5).

Equating these two expressions for P_{SOL} and replacing A_{nt} with its equivalent from flux conservation Eq. (i) gives:

$$\chi \frac{n_0 e T_0}{R_0 B_{p0} \Delta_{EO}} \cdot 2\pi \int_{null-}^{null+} R^2 B_p \cdot ds_p = 2n_t e T_t \gamma c_s 2\pi R_0 \Delta_{EO} \left(\frac{B_{p0}}{B_t} \right).$$

This simplifies in the high recycling mode because when mean free paths are short and fluid equations are appropriate, they require that $(1 + M^2)p$ is constant along the flow direction. Here p is the plasma pressure and M is the Mach number for plasma flow. Since $M = 0$ at the midplane where the net flow is zero and $M \approx 1$ at the sheath, we can set $n_0 kT_0 = 2n_t kT_t$. This does not contradict the earlier assumption of constant pressure along the separatrix because $M \approx 0$ everywhere except near the divertors. Using this simplification and solving for the thickness at the midplane, Δ_{EO} , results in

$$\Delta_{EO} = \sqrt{\frac{\chi \cdot L_p}{\gamma \cdot c_s}},$$

where

$$L_p = \frac{B_t}{B_{p0}} \int_{null-}^{null+} \frac{R^2 \cdot B_p}{R_0^2 B_{p0}} \cdot ds_p$$

and the integral runs in the poloidal plane from lower to upper null point, in either the outer or inner SOL. Notice that L_p is roughly equal to the projection of the usual connection length onto the poloidal plane, but it is not exactly equal because it depends differently on B_p . The uncertainty in

Table 2.7-2. Thickness Δ_{EO} of TIBER's outer and inner scrape-off layers at the midplane, and parameters used in the calculation. L_p is an effective length for power flow, projected onto the poloidal plane.

Ion sound speed at the sheath:		$c_S = 2.3 \times 10^4$ m/s
Energy (kT) per ion at the plate:		$\gamma = 7$
$L_p = 9.68$ m (outer)	9.58 m (inner)	
$\Delta_{EO} = 15.4$ mm (outer)	15.4 mm (inner)	when $\chi_i = 4$ m ² /s
$\Delta_{EO} = 10.9$ cm (outer)	10.9 mm (inner)	when $\chi_i = 2$ m ² /s

Δ_{EO} is due to the uncertainty in χ_i whose value is expected to lie in the range from 1 to 4 m²/s [6,7]. TIBER's parameters and the calculated values of Δ_{EO} are shown in Table 2.7-2.

2.7.5 PARTICLE AND HEAT FLUXES TO THE WALLS AND LIMITERS

We chose midplane spacings between the wall and the last closed flux surface of 120 mm in the outer and 60 mm in the inner SOL, respectively. The inner SOL spacing allows less than four particle e-fold lengths at the midplane, but increases to more than twice that many near the divertor. This rather narrow spacing could be increased, but at the cost of either reducing the neutron shielding or of increasing the plasma major radius.

2.7.5.1 Fluxes to the Inner Wall and Armor

In normal operation and when the plasma is in its normal position, the particle flux across the inner SOL to the wall armor is determined by the cross-field diffusion rate. We assume Bohm diffusion (as in the previous section) and write the radial particle flux Γ_1 :

$$\Gamma_1 = -D_{\perp} \nabla_{\perp} n \approx \frac{T_e}{16 \cdot B} \cdot \frac{n}{\Delta_N}.$$

Near the inner wall, the plasma density is

$$n_w \approx (5 \times 10^{19}/\text{m}^3) \cdot \exp \{-60 \text{ mm}/17.4 \text{ mm}\} \approx 2.5 \times 10^{18}/\text{m}^3$$

and $T_e = 76\text{ eV}$ (average across the SOL) and $B = 7.8\text{ T}$ from Table 2.7-1. These values give $\Gamma_{\perp} \approx 6 \times 10^{19}\text{ ions/m}^2/\text{s}$ at the wall. Recycling of the resultant gas increases the wall flux by about 40% to $\Gamma_{\perp,w} \approx 1 \times 10^{20}/\text{m}^2/\text{s}$. (The 40% reflux is discussed in Section 2.7.10.1.) Ion impact energy is determined by the local ion temperature and the sheath drop at the wall. Since the electron flow is almost entirely along the field lines to the divertors, charge balance in the SOL is maintained by the end sheath. The sheath drop at the wall is simply the sum of the end sheath drop plus the drop along the length of the SOL. We estimate the sheath drop at the inner wall as about equal to $3T_e$, using the local, average T_e from Table 3.7-1. This gives an ion impact energy $E_i \approx 2T_i + 3T_e \approx 400\text{ eV}$, while the recycled atoms strike the wall with the energy $E_{at} \approx (3/2)T \approx 150\text{ eV}$. The inner wall heat due to particle impact is then:

$$Q_{\perp,w} \approx e\Gamma_{\perp,w} \cdot (E_i + 0.4E_{at}) \approx 4.4 \times 10^3\text{ W/m}^2 = 0.0044\text{ MW/m}^2$$

which is completely negligible as compared to the 0.25 MW/m^2 of heating of the inner wall armor by electromagnetic radiation. A possibly more serious problem is that of sputtering. The physical sputtering yield for 400-eV D^+ is $Y_{C,D} \approx 0.026$, and for 400-eV T^+ it is $Y_{C,T} \approx 0.035$ atoms per ion [9], for normal incidence on carbon. The off-normal incident angle here could double these yields. We therefore assume a value $Y_C \approx 0.05$ atoms per ion. For the 150-eV neutrals the yields are about 0.020 for D^0 and 0.035 for T^0 . To allow for the spread in incident angles we use the same value for the 150-eV atoms as for the 400-eV ions to estimate the rate of erosion of the graphite inner wall armor:

$$\frac{dr_w}{dt} = \frac{\Gamma_{\perp,w} \cdot Y_C \cdot M_C}{\rho_C} \approx 2.5 \times 10^{-11}\text{ m/s} = 0.8\text{ mm/yr}$$

where $M_C = 12\text{ amu}$ is the atomic mass of carbon, and $\rho_C = 2260\text{ kg/m}^3$ is the density. Although this erosion rate is serious, it is slow enough to allow convenient scheduling of armor replacement.

2.7.5.2 Fluxes to the Outer Wall and Poloidal Limiters

At the outer wall the situation is different because the spacing is larger and because of the poloidal limiters that protect the wall, instead of

armor as on the inner wall. The spacing between plasma and wall is 120 mm at the midplane, and each of the 16 limiters extends 20 mm out from the wall. Figure 2.7-5 shows a cross section through one of the poloidal limiters. Again using values from Table 2.7-1 we have $\Delta_N = 20$ mm, $B = 4.35$ T, and $T = 80$ eV (average across the SOL). A similar calculation for the inner SOL gives an ion flux $\Gamma_i = 2 \times 10^{19}/\text{m}^2/\text{s}$ at the radius of the limiters, not including the reflux atoms (since they are not scraped off on the limiters). The plasma surface area facing the outer SOL is about 100 m^2 , so that the total radial flow out to the radius of the limiters is about 2×10^{21} ions/s. The power deposited on the limiters by this flux of ions is determined by the ion temperature and the sheath potential. We assume a sheath potential of about $3T_e$, as for the inner wall, and obtain an ion impact energy of about $E_{\text{imp}} \approx 400$ eV. Since the wall is located on a magnetic flux surface, no field lines lead to the wall or limiters. Therefore, any heat generated on the limiters by electrons should be negligible.

The heat deposited on the limiters in steady, normal operation is

$$H_{\text{lim}} = (2 \times 10^{21} \text{ ions/s}) \cdot e \cdot (400 \text{ eV/ion}) = 1.3 \times 10^5 \text{ W} = 0.13 \text{ MW},$$

which results in an average of 8 kW on each of the 16 limiters. The bombarded area on each limiter is determined by the radial distance that the ions diffuse out during the time between limiters, as they spiral around the machine on their way to the divertor. Radial diffusion causes an average radial velocity v :

$$v_i = \frac{\Gamma_i}{n} \approx \frac{T_e}{16B\Delta_N} \approx \frac{80}{16.4 \cdot 4.35 \cdot 0.0195} \approx 59 \text{ ms}$$

and the time for radial drift between limiters is determined by the limiter spacing, $\Delta L = 2\pi R_w/16 = 1.55$ m, and the ion velocity $v_i = \sqrt{eT_e/M_i} \approx 5.5 \times 10^4$ m/s. The time is therefore $t = 2.3 \times 10^{-5}$ s, which allows a radial drift of $v t = 1.65$ mm. Since the limiters are about 3.8 m long, each limiter has an exposed area normal to the field direction of about $A_{\text{lim}} \approx 6.3 \times 10^{-3} \text{ m}^2$. Since the heat is divided roughly evenly between the two faces of each limiter, the heat load normal to the field is therefore $0.008 \text{ MW}/0.0063 \text{ m}^2/2 = 0.6 \text{ MW/m}^2$ on each outer edge of each limiter. In fact,

Figure 2.7-5. A section through one of the poloidal limiters on the outer wall. Ions moving nearly parallel to the wall can diffuse radially 1.65 mm between limiters, and erode the outer surface from each direction.

the area is larger because the rounded edges of the limiters cause them to intercept the ions at an angle. This heat does not appear to be a serious problem.

In normal steady operation the limiters will be eroded because of sputtering by this ion flux. The ion flux to each side of each limiter, flowing parallel to the field at the limiter, is

$$\Gamma_{\parallel} \approx (2 \times 10^{21} \text{ ions/s}) / (2 \cdot 16 \cdot 1.65 \times 10^{-3} \text{ m} \cdot 3.8 \text{ m}) = 1 \times 10^{22} \text{ ions/m}^2/\text{s}.$$

Impact energy of the ions is $E_{\text{imp}} \approx 2T_i + 3T_e \approx 400 \text{ eV}$, as for the inner wall, and the average physical sputtering yield is again about $Y_C \approx 0.05$ atoms per ion. Since the ion flow is parallel to the wall, the erosion due to sputtering will cut 1.65-mm deep rounded steps from each side of the outer faces of the limiters (see Fig. 2.7-5), and these steps will grow in width but not in depth. The rate of growth of the widths is

$$dx/dt = \Gamma_{\parallel} Y_C M_C \rho_C \approx 4.4 \times 10^{-9} \text{ m/s} = 140 \text{ mm/yr}$$

from each direction across the face of each limiter. When the steps meet at the center of a limiter, it will then expose the next 1.65 mm on the neighboring limiter. Therefore, the rate that the limiters erode radially is 1.65 mm/yr if they are each 140-mm wide, and the device operates at full power for 50% of the time. In that case, new steps will be started at a rate of about one per year. Also as at the inner wall, about 40% of the neutralized ions will return to the wall and to the limiters as reflux atoms. About half of them will strike the limiters, and the other half will strike the wall. The heat generated by these reflux neutrals is negligible, as it is for the inner wall.

Each SOL allows nearly one full alpha-particle orbit diameter between the burning plasma and the wall. The rate of production of alphas varies as $n^2 \langle \sigma v \rangle_{DT}$, and $\langle \sigma v \rangle_{DT}$ decreases rapidly as ion temperature decreases toward the edge. The result is that only about 100 W of alpha power is produced within 40 mm of the inner separatrix, and practically none is produced within reach of the inner wall. The outer wall spacing was made greater because the space was not needed for neutron shielding.

2.7.6 EFFECT ON THE SOL OF A LARGE PUFF OF GAS

The edge plasma is a very good vacuum pump for gas loads up to a rather high limit, but the peak local gas load during pellet injection can exceed that limit. We discuss here some of the effects of a large puff of gas on the outer edge plasma. The most important conclusion is that the reflux of fast atoms to the wall and the erosion due to sputtering can be much less than would result if the same amount of gas was admitted slowly. When the gas density is great enough that the rate of charge exchange is limited by the flux of ions into the gas cloud, and the rate of ionization is limited by the much greater flux of electrons into the cloud, then ionization can become much more probable than charge exchange.

2.7.6.1 Plasma Interaction with a Large Puff of Gas

Consider a puff of DT gas released at the outer first wall. It requires about 10^{-4} s for the gas cloud to expand across the 12-cm separation between the outer wall and the confined plasma, since thermal speed for DT molecules at room temperature is about $v_{mo} \approx 1300$ m/s. Another few 10^{-4} s is required for the remainder of the gas to penetrate to that depth. TIBER will require pellet fueling at a rate of about 14 g/hr of both tritium and deuterium. Present estimates are that the propellant gas will equal the net fueling, and that another 20% will diffuse in from the pellet port [10]. If the fueling rate is five pellets per second, the gas cloud that accompanies each pellet will contain about $N_0 \approx 4 \times 10^{20}$ atoms, and the average gas density will be about $n_0 = 4 \times 10^{15}$ atoms/cm³ when the cloud has expanded to the separatrix. The mean free path for the 100-eV edge-plasma electrons in this cloud is

$$\lambda_e = \frac{v_e}{n_0 \langle \sigma_i v_i \rangle} \approx \frac{6 \times 10^8 \text{ cm/s}}{(4 \times 10^{15} \text{ cm}^{-3})(3 \times 10^{-8} \text{ cm}^3/\text{s})} = 5 \text{ cm} ,$$

and the mean free path for 100 eV edge-plasma ions in the cloud is

$$\lambda_i = \frac{v_i}{n_0 \langle \sigma_{cx} v_e \rangle} \approx \frac{1 \times 10^7 \text{ cm/s}}{(4 \times 10^{15} \text{ cm}^{-3})(4 \times 10^{-8} \text{ cm}^3/\text{s})} = 0.06 \text{ cm} .$$

Neither the electrons nor the ions can penetrate through the cloud. This means that for a few 10^{-4} s the cloud will make a shadow along the magnetic field lines in the edge plasma. Of greater importance is the fact that the

total reflux of fast atoms to the wall is limited by the finite flux of ions into the cloud. Since the ion flux into the cloud is less than the electron flux by a factor of about $\sqrt{m_e/M_i}$, the probability of charge exchange before electron ionization is also reduced by that factor. Of course, multiple charge exchanges can occur on the new cold ions, but this does not produce fast neutrals and reflux to the wall. The situation with a dense gas cloud is very different from the usual one where sufficiently few gas atoms are present at any one time so that the probabilities are determined mainly by the reaction rate coefficients $\langle \sigma_{cx} v_i \rangle$ and $\langle \sigma_i v_e \rangle$.

As the electrons are cooled by the cloud, the gradient in T_e causes power to flow in along the magnetic field lines. The energy available for ionizing the gas is some fraction (about half) of the stored energy in the magnetic tube defined by the dimensions of the cloud, plus about half of the energy incident on the tube from the hot plasma during the life of the cloud. The stored energy in the entire outer scrape-off layer is $E_{sto} = P_{SOL} \tau_E$, where P_{SOL} is the power flow in the outer SOL and τ_E is the average transit time for energy in the SOL. We can estimate τ_E from [6]:

$$\Delta_E = \sqrt{\frac{\chi_1 \tau_E}{3}}$$

using the value $\Delta_E \approx 0.015$ m calculated above for the thickness for power flow in the SOL. Here, $\chi_1 \approx 4 \text{ m}^2/\text{s}$ is the perpendicular thermal diffusivity in the plasma near the separatrix. This gives $\tau_E \approx 1.7 \times 10^{-4}$ s, which is roughly equal to the transit time for the gas. This means that the stored energy is replaced by new energy at a rate about equal to the rate that the cloud consumes it.

In TIBER, about 50 MW of alpha power plus about 50 MW of heating and current-drive power must flow out to the inner and outer SOLs. About 30% of this is radiated to the walls, leaving 70%, or 70 MW, to flow in the two SOLs. Of this, about 50 MW flows in the outer SOL. Only a fraction of the SOL power flows on field lines that pass through the cloud. This fraction is the fraction, $\Delta\psi_{gas}/\psi_{SOL}$, of the magnetic flux in the SOL that threads the cloud

$$\frac{\Delta\psi_{gas}}{\psi_{SOL}} = \frac{\Delta_E^d \text{cloud} B}{2\pi(R+a)\Delta_E B_p} \approx \frac{(0.015 \text{ m})(0.15 \text{ m})(4.8 \text{ T})}{2\pi(3.83 \text{ m})(0.015 \text{ m})(2.0 \text{ T})} = 0.015 ,$$

where the values of the parameters are for the outer midplane in TIBER. Therefore, about 0.015 of the 50 MW flows in the flux tube, and about half of that flows to the cloud. The energy available to ionize the gas in the cloud is therefore about

$$E_{\text{ioniz}} = (\Delta\psi_{\text{gas}}/\psi_{\text{SOL}}) \cdot (1/2) P_{\text{SOL}} \cdot t \\ \approx (0.015)(1/2)(5 \times 10^7 \text{ W})(5 \times 10^{-4} \text{ s}) \approx 200 \text{ J} \approx 1 \times 10^{21} \text{ eV}$$

if the total interaction time is about $t \approx 5 \times 10^{-4} \text{ s}$. But, to dissociate the gas molecules and ionize all of the atoms requires at least 30 eV per atom, plus roughly an equal amount of energy that is radiated before the final ionization. Little energy is used in heating the new ions because the duration of this phase is shorter than electron-ion thermalization time ($\tau_{ie} \approx 1 \times 10^{-3} \text{ s}$). About 60 eV per atom, or $2.4 \times 10^{22} \text{ eV}$ for the entire cloud, is therefore used in ionizing the gas. The energy available from the SOL, $E_{\text{ioniz}} \approx 1 \times 10^{21} \text{ eV}$, is only enough to ionize about 4% of the gas as it passes through the SOL.

The remainder of the gas penetrates a short distance into the confined plasma. The penetration distance into the core plasma is only a few centimeters because the power flux is more than two orders of magnitude higher inside the separatrix than outside in the SOL. In most cases of hydrogen gas interacting with a hydrogen plasma, the rate of charge exchange (CX) exceeds that of ionization, and the fast atoms that result from CX can penetrate deeply. Here, the rate of CX is limited by the ion flux into the cloud rather than by the number of gas atoms. Even though the rate coefficient for CX is larger than that for ionization, the fact that the electron flux is much greater than the ion flux causes ionization to dominate over CX. This reduces the rate of CX and reduces the penetration of the gas and the reflux of fast atoms to the wall.

2.7.6.2 Reflux and Erosion at the Wall

The electron energy flux flowing into the cloud along the magnetic field from either side is

$$Q_{\parallel e} = (1/4) n_e v_e (3/2) kT_e.$$

This energy will produce ions at a rate $R_{ion} = Q_{we} / \Delta E_{ioniz}$ per unit of exposed area per second, if ΔE_{ioniz} is the average energy lost from the electrons per new ion. The corresponding rate of charge exchange is just equal to the flux of hot ions into the cloud: $R_{CX} = \Gamma_{CX} = (1/4)n_i v_i$. The ratio of these two rates gives the number of charge exchange events per ionization:

$$\frac{R_{CX}}{R_{ion}} = \frac{(1/4)n_i v_i}{(1/4)n_e v_e (3/2)kT_e / \Delta E_{ioniz}} = (2/3)(v_i/v_e)(\Delta E_{ioniz}/kT_e) .$$

Since all of the gas atoms either CX or become ionized, the fraction that CX is

$$\frac{N_{CX}}{N_0} = \frac{R_{CX}}{R_{ion} + R_{CX}} \approx \frac{R_{CX}}{R_{ion}} \approx \frac{m_e}{M_i} \cdot \frac{2\Delta E_{ioniz}}{3kT_e} .$$

Since the time is too short for thermalization, $\Delta E_{ioniz} \approx 50$ eV is just the sum of the average energy radiated per ionization plus dissociation and ionization energies. Near the separatrix where the reactions take place, $T_e \approx 100$ to 1000 eV, giving a ratio:

$$N_{CX}/N_0 \approx 5 \times 10^{-3} \text{ to } 5 \times 10^{-4} .$$

Compare this with the situation when CX is limited only by the supply of gas atoms. In that case, the fraction that CX would be $N_{CS}/N_0 = \langle \sigma_{CX} v_i \rangle / (\langle \sigma_{CX} v_i \rangle + \langle \sigma_i v_e \rangle) > 0.7$, giving two to three orders of magnitude greater reflux to the wall.

The rate of sputter erosion of the wall by fast CX atoms is therefore reduced when the gas is introduced in large puffs. Since these CX reactions take place near the plasma surface, about half of the CX atoms will strike the wall and the other half will penetrate into the plasma. The sputtering yield for 1000-eV D-atoms incident normally on stainless steel is about 0.028 atoms per D-atom [9], and for T-atoms it is about 2.1 times as great, or 0.058 atoms per T-atom. At 100 eV the yield is about ten times smaller, but N_{CX} would be ten times larger, so that the sputtering rate would be about the same. Using the average value of 0.043, we estimate that the number of metal atoms N_{SS} removed per pellet is $N_{SS} \approx (1/2)(0.043)N_{CX} \approx 5 \times 10^{15}$ per pellet. This is

about 76 g/y, or $10 \text{ cm}^3/\text{y}$ of metal, when the injection rate is 5 pellets per second. This erosion takes place over a wall area of about one square meter, so that the average erosion rate is negligibly small.

In order to estimate the rate of erosion as a function of distance from the pellet port, we model the wall as being a plane surface and assume that the CX atoms are launched uniformly in angle. This assumption results in a CX flux, Γ_{CX} , that decreases as $1/r^2$, where r is the distance from the source of CX. The sketch in Fig. 2.7-6 shows the notation used. After allowing for the angle of incidence, the CX flux to the wall is

$$\Gamma_{\text{CX},W} = \Gamma_{\text{CX}} \cos(\theta) \approx \frac{\nu N_{\text{CX}}}{4\pi h^2 [1 + (x/h)^2]^{3/2}},$$

where $\nu \approx 5$ pellets/s is the pellet rate, $h \approx 12$ cm is the normal distance from the source of CX to the wall, and x is the radial distance along the wall from the axis of the pellet port (see Fig. 2.7-6). This predicts a peak flux to the wall of

$$\Gamma_{\text{CX},W} \leq \frac{\nu N_{\text{CX}}}{4\pi h^2} \approx \frac{(5)(2 \times 10^{17})}{4\pi \cdot 12^2} = 6 \times 10^{14} \text{ atoms/cm}^2/\text{s}$$

and a corresponding peak erosion rate of 0.043 times this, or 3×10^{13} SS-atoms/cm²/s, which is 0.08 g/y/cm², or 0.1 mm/y.

2.7.7 HEAT LOAD ON THE DIVERTOR PLATES

The distribution of heat on the divertor plates is calculated by assuming that the energy density in the scrape-off layer (SOL) decreases exponentially with distance out from the plasma at the midplane. The e-fold distance at the midplane, Δ_E , is discussed in Section 2.7.4. This exponential distribution is then mapped by following magnetic field lines to the divertor plates to determine the power intercepted by each increment of area on the plates. One can choose a shape for the plates that results in a nearly uniform heat load, but such special shaping appears to make the peak heat load very sensitive to even a small displacement of the plasma. A computer code was written to allow interactive adjustment of the shape in order to search for a shape that both gives a reasonably low ratio of peak-to-average heat load and is also rather insensitive to plasma displacement.

Figure 2.7-6. Parameters used to calculate the flux profile at the wall, of atoms that charge exchange on gas from the pellet port.

Figure 2.7-7(a) shows the divertor region as modeled by the code. The divertor is specified as a series of coordinate points, and the computer then maps the power flow and plots the average heat loads on the segments. Fig. 2.7-7(b) shows the average heat loads plotted against the distance from the separatrix, measured along the surface of the plate. Total power flow in the SOL is set equal to 70% of the total power carried out by charged particles, which is equal to the sum of alpha power, heating power, and current-drive power. The other 30% is assumed to be radiated to other surfaces and does not reach the divertors. For the figures, it is further assumed that the SOL power divides with 2/3 of it flowing in the outer SOL and 1/3 flowing in the inner SOL. The data is also presented for an 80/20 division of power. Total charged particle power in TIBER is about 110 MW, with alpha power accounting for about half of it. The three parts of the divertor are labeled "inner," "lower," and "outer." Figure 2.7-8 is an array of figures; the central figure is a repeat of Fig. 2.7-7(a), but with finer segmenting of the plates. The surrounding figures show the plasma displaced by 10 cm in different directions, as labeled. Figure 2.7-9 shows the corresponding heat-load plots, arranged in the same order as in Fig. 2.7-8. It will be seen in the figure that the peak heat load (3.6 MW/m^2) occurs on the lower divertor plate when the plasma is displaced either upward (away from the plate) or outward. One can get the heat loads for other divisions of the power between outer and inner SOLs by scaling from the 2/3-to-1/3 division assumed in the figures.

Table 2.7-3 lists the highest heat loads on each of the three plates of the divertor: inner, lower, and outer (see Fig. 2.7-3) when the plasma is shifted by 10 cm up or down and/or in or out. The highest values in the table are 4.3 and 4.2 MW/m^2 on the lower plate when the plasma is shifted upward by 10 cm or radially outward by 10 cm, respectively, and when the outer SOL carries 80% of the divertor heat. A worst case would be when this condition exists, and when one null is partially lost so that the heat to the other lower plate increases by 50%. This would give a maximum of 6.6 MW/m^2 on that lower plate.

For normal steady operation we assume that the power divides equally between upper and lower divertors. A perfect 50:50 ratio requires exact symmetry about the midplane. Almost any small asymmetry make the path along the field lines to one divertor shorter than the other. This effect is

Figure 2.7-7. Heat loads on the divertor plates: (a) the coordinate points used in the flux mapping of the power flow. (b) average heat loads on segments of the plates between coordinate points. Distance s is measured along the surface away from the intersection with the separatrix.

Figure 2.7-8. Showing the location of the divertor plates relative to the plasma for the displacements of the plasma by 0 or 10 cm in each direction. Smaller segments are used here than in Fig. 2.7-7(a). The array is in the same order as in Fig. 2.7-9 and in Table 2.7-3.

Figure 2.7-9. Heat profiles as in Fig. 2.7-7, but with the plasma shifted as shown in Fig. 2.7-8. Here, the total heat is 50 MW to the outer divertor and 25 MW to the inner, and $\Delta E_{out} = \Delta E_{in} = 15$ mm. The dotted curve is for the inner plate, dashed curve is for the lower plate, and the solid curve is for the outer plate.

Table 2.7-3. The peak heat loads are shown on the three divertor plates (inner, lower, and outer) when the plasma position shifts by ± 10 cm in or out and/or up or down. Values are given in the same order as shown in Figs. 2.7-8 and 2.7-9. Results for two values of the inner/outer power ratio are shown.

Values are given in the order:

Inner divertor i/o power ratio 1:2 (20:80)^c

Lower divertor i/o power ratio 1:2 (20:80)

Outer divertor i/o power ratio 1:2 (20:80)

Heat in MW/m ²		+ ΔR (cm) +		
		-10	0	+10
	+10	1.6 (1.0)	2.7 (1.6)	3.5 (2.1)
		2.9 (3.5)	3.6 (4.3)	0 (0) ^a
		1.7 (2.0)	2.3 (2.8)	1.6 (1.9)
ΔZ (cm)	0	1.8 (1.1)	2.5 (1.5)	3.5 (2.1)
		2.6 (3.1)	3.1 (3.7)	3.5 (4.2)
		1.2 (1.4)	1.3 (1.6)	1.7 (2.0)
	-10	2.1 (1.3)	2.3 (1.4)	3.5 (2.1)
		2.8 (3.4)	2.8 (3.4)	3.0 (3.6)
		1.1 (1.3)	2.0 (2.4)	3.3 (4.0) ^b

^aThe power flow is moved completely off the lower plate.

^bPlasma is almost in contact with the outer plate.

amplified by the shallow pitch angle due to the weak poloidal field near the nulls. A large deviation from symmetry could destroy one null and allow the other divertor to act as a limiter. Assuming that we can avoid the complete loss of either null, the extreme ratio of power to the two divertors that we

might expect is about 25/75. The effect of this ratio can be seen from Table 2.7-3 by multiplying the outer heat loads by 1.5 and the inner loads by 0.5.

The maximum shift of a null point that this configuration allows is about 10 cm either downward, or in or out. Since the plasma will start small and grow to the final configuration during start-up, and since the plasma shape will tend to relax away from the crescent shape, it is unlikely that a null point will shift very far inward or downward. Also, the spacing between confined plasma and the inner wall at the midplane is only 6 cm, which means that the divertor configuration will be lost and the inner wall will become a limiter if the plasma shifts inward by more than 6 cm without changing shape.

2.7.8 HEAT AND PARTICLE FLUENCES IN A DISRUPTION

A plasma disruption results in two processes, the thermal quench which takes place in tens of μs , followed by the decay of the plasma current which takes place at a rate of about 1 MA/ms. The thermal quench apparently occurs when χ , suddenly increases and allows the stored kinetic energy of the plasma to be transported out radially. Since plasma β is not large, the sudden cooling of the plasma does not significantly affect the magnetic field. For this reason the energy flow in the thermal quench follows nearly the same path as in steady operation and is therefore deposited mainly on the divertor plates. The differences from steady operation are the fast time scale and the fact that the faster radial transport causes the heat to be spread over a somewhat larger area on the divertor plates. Also, a larger fraction of the power is radiated because of the higher density of impurities that results from heating the plates. Plasma current decays on an L/R time scale, where L is the mutual inductance and R is the electrical resistance for current in the plasma. Although R varies with time as the plasma temperature changes, the L/R time is much longer than the thermal time. Immediately after the thermal quench, the plasma current is left flowing in a cooler plasma. The resistive drag on the current heats the plasma and adds new kinetic energy, which finally flows out along the open field lines. These lines move as the current decays, and some may intercept the wall before reaching the divertor. However, most of the energy stored magnetically by the plasma current is either radiated to the walls or converted to heat in the various electrical conductors in which current is induced as the plasma current changes.

In TIBER, the stored kinetic energy of electrons plus ions is about 118 MJ, and the stored magnetic energy is 244 MJ. Since the stored kinetic energy flows out along the field lines to the divertors, the heat-load plots shown in Figs. 2.7-6 and 2.7-8 can be scaled to show the profiles of energy deposition. After the current has begun to decay and the magnetic field has become distorted, a separate calculation must be made to determine the history of the heat and particle fluxes. A calculation of the induced currents and of the resultant heat and magnetic forces must be done separately. See Section 2.10, "Disruption Modeling".

2.7.9 MICROWAVE REFLECTION BY CARBON-COATED WALLS

Graphite is one possible choice of material for TIBER's limiters, divertor plates, and any protective armor. Graphite has been chosen for TFTR, JET, D-III, and now JT-60 because it has been partly responsible for "super shots" and H-mode operation. The disastrous chemical sputtering that was originally predicted has not occurred, and by preconditioning, graphite can be made to temporarily pump hydrogen gas at the walls. Because graphite has worked better than originally expected, the other possible materials have not been carefully tested. The only other materials that appear to be candidates are beryllium (with its low Z and good neutronics but low melting point and low energy threshold for sputtering) and tungsten (with a high melting point and high threshold for sputtering, but with a high Z).

The concern about the use of graphite in TIBER is that sputtered carbon will coat the walls and possibly reduce the wall reflectivity for high-frequency rf power. Highly reflecting walls are required if ECH is used either for current drive or for plasma heating in TIBER. Even if no applied rf power is used, it would be better if the walls reflect the synchrotron radiation rather than allow it to be absorbed and used to heat the walls. Here we consider the reflectivity of metal walls coated with a layer of carbon.

Reflectivity takes place within about one "skin depth" into the wall. If the carbon coating on a wall is less than a skin depth, then the metal wall will contribute to the reflectivity. The skin depth, d , is given for frequency ω at a plane surface of a material with conductivity σ , by

$$d = \sqrt{\frac{2}{\mu\sigma\omega}}.$$

For non-magnetic materials, $\mu = \mu_0 = 4\pi \times 10^{-7}$ h/m. Kohl [11] gives the conductivity of carbon as $\sigma = 2.2 \times 10^4$ to 2.9×10^4 mho/m and of graphite as $\sigma = 7.7 \times 10^4$ to 12.5×10^4 mho/m. We assume that carbon sputtered onto the walls will be amorphous and therefore have an $\sigma \approx 2.5 \times 10^4$ mho/m. When the frequency is equal to the electron cyclotron frequency, $\omega_{ce} = eB/m_e = 1 \times 10^{12} \text{ s}^{-1}$ at $B = 6 \text{ T}$, the skin depth in carbon is $d \approx 7.8 \mu\text{m}$.

It appears reasonable to assume for reflectivity calculations that if TIBER's limiters or divertors use graphite, then TIBER's walls will become coated to such a thickness that they will reflect and absorb like solid carbon. Fowles [12] gives the Hagen-Rubens formula for the reflection coefficient, R , at normal incidence:

$$R = 1 - 2\sqrt{\frac{2\omega\epsilon_0}{\sigma}} = 1 - 0.055\sqrt{\frac{\omega}{1.05 \times 10^{12}/\text{s}} \cdot \frac{2.5 \times 10^4 \text{ mho/m}}{\sigma}},$$

where $\epsilon_0 = 8.85 \times 10^{-12} \text{ f/m}$ is the permittivity of vacuum, and we have used $\omega_{ce} = 1.05 \times 10^{12} \text{ s}^{-1}$ as a reference frequency. This predicts 5.5% absorption at the fundamental of the electron cyclotron frequency. The spectrum of Doppler-shifted synchrotron radiation extends out to high harmonics, and wall absorption would become large. For off-normal angles of incidence, R varies differently, depending on the polarization of the radiation. In general, $R_{\perp}$ increases and $R_{\parallel}$ decreases as the angle deviates from normal incidence (see Fig. 5.6 in Ref. 12). Since synchrotron radiation is incident over a wide range of angles and over all polarization directions, the effective reflectivity would be roughly equal to the R for normal incidence, given above.

2.7.10 PARTICLE AND ENERGY TRANSPORT IN THE EDGE PLASMA

The density and temperature in the SOL are roughly $n \approx 5 \times 10^{19}/\text{m}^3$ and $T_e \approx 100 \text{ eV}$. In this plasma the mean free path for electron-electron scattering is $\lambda_{ee}(\text{m}) \approx 1 \times 10^{16} T_e^2/n_e \approx 2 \text{ m}$, and $\lambda_{ii} \approx \lambda_{ee}$. Close to the divertor the mean free paths become shorter. Since these are smaller than the scale length for temperature gradients, we can assume to good approximation that ions and electrons have Maxwellian velocity distributions. However, the mean free path for ions to thermalize with electrons is about 10 m, but this decreases to a few centimeters where the density increases and the temperature decreases near the divertor plates. This means that the electron and ion

temperatures will be different in the hotter part of the SOL, but will approach equality near the divertor. Since we are not particularly concerned here with the ion temperature T_i , we assume that the SOL plasma flows like a fluid to the divertor, and we set $T_i = T_e$ at each point along the flow path. We further assume that electron thermal conduction dominates over convection, and use classical ("Spitzer") electron conductivity.

2.7.10.1 The 1-D Fluid Equations

Net particle and energy flow in the SOL is much faster along the magnetic field than across it. It is therefore a rather good approximation to ignore the cross-field transport when calculating the parallel fluxes. The omission of cross-field diffusion should result in somewhat higher peak heat loads and particle fluxes to the divertor plates than one would calculate if diffusion were included. Here, we use one-dimensional fluid equations, but include variations of the flux tube cross section, $A_{\parallel}$, along the length of the spiralling tube.

We measure distance s along the flow direction from the midplane and identify the particular flux tube being followed by its distance x , measured out from the separatrix at the midplane. As dependent variables we choose: particle flux written as a current density j (A/m^2), plasma temperature T (eV) ($T_i = T_e$ is assumed), and plasma pressure in the field direction p (Pa). Plasma density, n , and flow velocity, v , can be obtained from these dependent variables using the equations: $j = nev$ and $p = 2neT$, or $n = p/(2eT)$ and $v = 2jT/p$. In terms of j , T , and p , the continuity equation is

$$\frac{d(A_{\parallel}j)}{ds} = A_{\parallel}S,$$

where the ion source is $S(s,x) = en_e n_0 \langle \sigma_i v_e \rangle$. The units of S are A/m^3 to be consistent with those of j . The integral form of the continuity equation is

$$j(s,x) = \frac{1}{A_{\parallel}(s)} \int_0^s A_{\parallel}(s') S(s';x) ds'.$$

By symmetry, $j(0,x) = 0$ at the midplane. The one-dimensional momentum equation can be written as:

$$\left(\frac{2M_i}{e}\right) \frac{d}{ds} \left(\frac{A_{||} j^2 T}{p}\right) + A_{||} \frac{dp}{ds} = 0 .$$

The right-hand side is zero because we assume no momentum source term. When $A_{||}$ is constant, this equation can be easily integrated to give

$$(2M_i/e)[j^2 T/p] + p = p_0 ,$$

with p_0 being the constant pressure at the midplane where $j = 0$. The momentum equation appears especially simple when $A_{||}$ is constant and the equation is written in terms of Mach number, $M = v/c_s$, where $c_s = \sqrt{(2eT/M_i)}$ is the speed of sound in the SOL plasma. In terms of M and with constant $A_{||}$, the equation becomes

$$(1 + M^2) \cdot p = p_0 ,$$

which satisfies the $p = p_0$ condition at the midplane where $M = 0$ and shows that the pressure at the sheath $p_{sh} = p(s_L) = p_0/2$ where the Bohm sheath criterion requires $M > 1$ at the sheath [13]. The singularity in p at $M = 1$ complicates the solving of fluid equations.

We combine the energy equations for ions and for electrons to get

$$\frac{d}{ds} \left(- A_{||} \kappa T^{5/2} \frac{dT}{ds} + 5 A_{||} j T \right) = A_{||} q - E_r A_{||} S .$$

The jv^2 term is omitted here because it is much smaller than the $5jT$ term everywhere except near the divertor plate because flow speed, v , is much smaller than sound speed, c_s , except near the sheath. The coefficient for (Spitzer) electron thermal conductivity is $\kappa [W(eV)^{-7/2} m^{-1}] = 3.15 \times 10^4 / (Z_{eff} \ln \Lambda)$, where Z_{eff} is the effective charge state of the SOL plasma. On the right-hand side of the equation, $q = q(s, x)$ is the volumetric power source, and E_r is the average energy lost from the flow per ionization due to radiation. For hydrogen atoms, $E_r \approx 30$ eV per ionization [14].

We define two regions in the SOL, as shown in Fig. 2.7-10. Region I is the part of the SOL between the midplane where $s = 0$ and some point s_{dp} near the divertor where gas ionization begins to be important. Region I occupies the length $0 \leq s \leq s_{dp}$. Region II is the region near the plates where gas

Figure 2.7-10. Showing the geometry used for the 1-D fluid model of the SOL. Distance s is measured along the magnetic field from the midplane. Region I is adjacent to the confined plasma. Region II is near the divertor where gas recycling takes place.

ionization is important, and we identify this region by $s_{dp} \leq s \leq s_L$, where s_L is the value of s at the plate. Table 2.7-4 lists the values of these and other parameters for TIBER's edge plasma. We use the above fluid equations, with the neutral density n_0 in S calculated from the rate that ions are neutralized on the plate directly adjacent to each point s along the path. To include this in a 1-D calculation where the currents in adjacent flux tubes are not calculated, we use the boundary conditions at the sheath to determine the ion flux being neutralized, and then allow for attenuation of the neutrals due to ionization between the plate and the position of interest on the flux tube.

2.7.10.2. Power Flow in Region I

In this region ($0 \leq s \leq s_{dp}$) we assume that $S = 0$ and $j = 0$. In that case, the momentum equation shows that $p(s, x) = p(0, x)$ is constant along a field line. The energy equation is also simplified. The remaining terms after a first integration are

$$- A_{\parallel}(s) \kappa T^{5/2} \frac{dT}{ds} = \int_0^s A_{\parallel}(s') q(s', x) ds' ,$$

since $dT/ds|_{s=0} = 0$. This can be integrated again to give

$$T^{7/2}(s, x) - T^{7/2}(s_{dp}, x) = \frac{7}{2\kappa} \int_0^s Q_{\parallel}(s', x) ds' ,$$

where we have set

$$Q_{\parallel}(s, x) = \frac{1}{A_{\parallel}(s)} \int_0^s A_{\parallel}(s', x) q(s', x) ds' .$$

This expression gives the temperature variation along the field line in Region I and, in particular, determines the temperature difference between midplane ($s = 0$) and the beginning of Region II ($s = s_{dp}$). The solutions in Region I and Region II must join smoothly at $s = s_{dp}$.

Table 2.7-4. TIBER parameters. A computer code SEPTRIX generates the parameters needed to calculate the particle and energy flow in the SOL. For given distance x out from the separatrix at the midplane, the code integrates the various functions along a spiraling field line from midplane, $s = 0$, to either the null point, $s = s_x$, or to the divertor plate, $s = s_L$.

For TIBER, $q(s,x)$ is fit rather well by the function

$$q(s,x) \approx \frac{c \cdot \exp\{-x/\Delta_E\}}{[1 + (s/b)^{5/2}]} \quad \begin{array}{ll} c_{out} = 56 \text{ MW/m}^3 & c_{in} = 29.5 \text{ MW/m}^3 \\ b_{out} = 3.2 \text{ m} & b_{in} = 4.2 \text{ m} \end{array}$$

Also for TIBER, $A_{||}$ is nearly constant and independent of both s and x in both outer and inner SOLs. $A_{||,outer} \approx 0.12 \text{ m}^2$, and $A_{||,inner} \approx 0.057 \text{ m}^2$.

The null point occurs at $R_x = 2.437$, $Z_x = \pm 2.353 \text{ m}$. SEPTRIX gives the locations of s_x , s_{dp} , and s_L on the field line. Here $x = 1.0 \text{ cm} = \Delta_E \cdot \ln(2)$, which centers the field line in the power flow:

	Outer	Inner	
s_x , null point	17.8	34	(m)
s_{dp} , recycle starts	~11	~15	(m)
s_L , divertor plate	26	38	(m)
$Q_{ }(s_{dp}, 0)$	218	151	(MW/m ²)
$\int_0^{s_{dp}} Q_{ }(s) ds$	1792.	1710.	(MW/m)

2.7.10.3. Power Flow in Region II

Integrating the energy equation once in Region II where $s_{dp} \leq s \leq s_L$, and using our assumption that $j(s_{dp}, x) = 0$, gives

$$-A_{||} \frac{2\kappa}{7} \cdot \frac{dT^{7/2}}{ds} + 5A_{||} jT = -\frac{2\kappa}{7} A_{||} \frac{dT^{7/2}}{ds} s_{dp} + \int_{s_{dp}}^s A_{||} q \cdot ds - E_r \int_{s_{dp}}^s A_{||} S \cdot ds$$

The continuity equation shows that the last term can be replaced by $-E_r A_{\parallel}(s) j(s)$. We again use our definition of $Q_{\parallel}$, and set $U = T^{1/2}$, to rewrite this energy equation:

$$U'(s) = \frac{A_{\parallel dp}}{A_{\parallel}} \cdot U'(s_{dp}) - \frac{7}{2\kappa} \cdot \left(Q_{\parallel}(s) - Q_{\parallel}(s_{dp}) - [5T(s) + E_r] \cdot j(s) \right).$$

Here $U' = dU/ds$. When $q(s, x) = 0$ as it is in Region II, then $Q_{\parallel}(s) = Q_{\parallel}(s_{dp})$ in Region I, and the equation simplifies to:

$$U'(s) = (A_{\parallel dp}/A_{\parallel}) U'(s_{dp}) + (7/2\kappa) \cdot [5T(s) + E_r] \cdot j(s).$$

Using this form of the energy equation in Region II, the 1-D fluid equations are integrated numerically. The numerical integration is iterated until the value of $U(s_{dp})$ agrees with the power input to the SOL and there is power, pressure, and particle balance at the sheath. These sheath boundary conditions are discussed next.

2.7.10.4. Boundary Conditions at the Sheath

A sheath forms at the interface between the SOL plasma and the divertor plate, located at $s = s_L$. The sheath forms because the average speed of the electrons is much greater than that of the ions, and initially this allows the electrons to escape to the plate faster than the ions. But as electrons leave, the potential of the plasma rises positively until the potential drop across the sheath retards the electron escape rate and makes it equal to that of the ions. The thickness of the sheath is of the order on a Debye length, which is very small ($\ll 1$ mm) in the warm, dense SOL plasma near the plate. The boundary conditions at the sheath are determined by the requirements that particle, pressure, and energy balance exist there. Particle balance requires that:

$$j_{sh} = j(s_L) = \frac{1}{A_{\parallel}(s)} \int_{s_{dp}}^{s_L} A_{\parallel}(s) S[T(s), n_0(s), p(s)] ds,$$

where for convenience, the lower limit of integration can be taken as $s = s_{dp}$ since $S = 0$ in $0 \leq s \leq s_{dp}$. When $A_{\parallel}$ is constant, pressure balance at the sheath requires that:

$$p(s_L) = (1/2)p(s_{dp}) = (1/2)p_0 ,$$

as determined by the sheath criterion. This condition means that only half of the static pressure can be supported by momentum transfer to the ions produced by gas ionization without the production of a shock wave where the Mach number passes through $M = 1$. Finally, energy balance at the sheath requires that

$$\begin{aligned} Q_{sh} &= \gamma j(s_L, x) T(s_L, x) \\ &= [Q_{II}(s_{dp}, x) - E_r j(s_L, x)] \frac{A_{II}(s_{dp})}{A_{II}(s_L)} \\ &= - \frac{2\kappa}{7} \cdot \frac{dT^{7/2}}{ds} \Big|_{s_L, x} + 5j(s_L, x) T(s_L, x) . \end{aligned}$$

These various forms of the energy balance condition can be rewritten to show that

$$\left(\frac{dT^{7/2}}{ds} \right)_{sh} = - (\gamma - 5) \frac{7T_{sh} j_{sh}}{2\kappa} ,$$

which establishes the relation between heat conduction and convection into the sheath. The value of γ is considered next.

2.7.10.5. The Sheath Factor γ

The sheath factor γ gives the energy loss from the SOL plasma across the sheath per ion, including that of the associated electron. In units of T_e at the sheath, its value is [15]

$$\gamma k T_e = [2kT_i + Ze(\phi + \phi')] (1 - R_E) + \chi_i + \phi_w + \frac{2\kappa T_e}{1-\lambda} ,$$

Here, $Ze(\phi + \phi') \approx 3.5kT_e$ is due to the acceleration of the ion across the sheath and pre-sheath; R_E is the energy reflection coefficient which is a function of the ion impact energy $[2kT_i + Ze(\phi + \phi')]$; χ_i is the threshold energy for ionization; and ϕ_w is the work function for the plate material. λ is the secondary electron emission coefficient.

When λ is negligible, this gives $\gamma \approx 8$, depending on R_E . While λ determines the ion flux required to carry the power across the sheath, a more important quantity is the ion impact energy, since this determines the sputtering coefficient. Ion impact energy is determined partly by T_i , but mostly by the sheath potential ϕ . The sheath potential adjusts itself to restrict the electron flux across the sheath and therefore maintains charge balance. It should be possible to reduce the ϕ and the ion impact energy by deliberately increasing λ . When λ is large and secondary electrons are released into the plasma, then ϕ must decrease to allow more electrons out. The net effect is that the electrons then carry out a larger fraction of the power than before, and the ions can impact at lower energy.

2.7.11 GAS RECYCLING AT THE DIVERTOR

The divertor shape was also chosen to enhance the recycling of gas at the plates by sloping toward the plasma so that gas atoms from the plates are directed into increasing plasma density and temperature. A high recycle rate is desirable because it can reduce the average impact energy of the ions by allowing a much larger number of particles to share the power. When the impact energy is reduced to near the threshold for sputtering of plate material, the impurity ion production rate and the rate of erosion of the plates can be greatly reduced. To reduce the ion impact energy to some value E_{th} one must reduce the electron temperature to about $T_e \approx E_{th}/(1 + 3Z)$ because the drop in potential across the sheath accelerates the ions before they impact. Threshold energies for a few combinations of ions and plate material [8,9] are shown in Table 2.7-5. It is seen in the table that tungsten has the highest threshold, $E_{th} \approx 175, 140$, and 100 eV for D^+ , T^+ , and He^{++} . Even tungsten requires $T_e < 35$ eV to prevent sputtering by fuel ions, and $T_e < 14$ eV to prevent sputtering by helium ions. Other materials, such as Be or graphite, have much lower thresholds. Since low- Z impurities are less serious than high- Z impurities, a higher sputtering rate can be tolerated for Be and C than for W.

2.7.11.1. Gas Penetration into the SOL

The 1-D code GASREFLX was used to study the interaction of gas with the edge plasma at a divertor plate. Sophisticated 2-D and 3-D codes [16,17,18] will be used in the future to more accurately model the interaction and

Table 2.7-5. Threshold energy (eV) for sputtering of plate material by impacting ions. Data is from Ref. [8].

Target	Ion			
	H	D	T	He
Be	20	20	21 ^a	30
C	15	25	30 ^a	65
Fe	68	46	40 ^a	33
W	458	160	140 ^a	100

^aEstimated value.

determine the plasma parameters near the divertor. GASREFLX calculates the penetration of H^0 into an H^+ plasma, where "H" may be H, D, or T. The code launches molecules from a plate with a cosine distribution in angle. The molecules decay along their paths with an e-fold length equal to their mean free path for ionization in the plasma whose temperature and density can vary with distance. Each atom in a molecule yields some fraction of one Franck-Condon atom (FC), where the fraction is calculated as a function of the plasma temperature. The FCs are directed randomly: some strike the plate, while others are directed at some angle into the plasma. They either are ionized or they undergo charge exchange (CX). The fraction-per-atom undergoing CX produces a fraction of a new atom with random direction and with the energy of the ions at the point of reaction. Then, the CX-neutral can either be ionized or undergo another CX. The process continues until the fraction remaining is negligible. The code records the flux of atoms to a plate for each type of process and the fluxes that penetrate beyond a specified distance into the plasma. Fixed reflection coefficients are used for the atoms that strike the plate. For example, by comparing the results when a reflection coefficient is zero with that when it is unity, we can determine the effect of each type of particle on the penetration across the separatrix,. The code accumulates the total reflux of each type of neutral to the wall, as well as the fluxes across a specified boundary such as the separatrix.

We can estimate the reflux to the wall by treating the multiple CXs as a diffusion process. Harrison [15] considers the flux of type "a" particles from the plate at a depth x out from the plate: $\Gamma^a(x) = \Gamma^a(0) \cdot \exp\{-x/\Delta_a\}$, where

$$\frac{1}{\Delta_a} = \frac{3}{\lambda_{cx}^a} + \frac{3}{\lambda_i^a},$$

where λ_{cx}^a and λ_i^a are the mean free paths for type "a" atoms for charge exchange and for ionization, respectively. The factors of 3 on the right-hand side were omitted by Harrison, but are needed since we are considering the penetration in only one direction. At a depth x , CX reactions produce type "b" particles at a rate

$$dk^{a+b}(x) = \Gamma^a(x) \cdot \frac{dx}{(1/3)\lambda_{cx}^a}.$$

Half of the type "b" atoms are directed so that they can strike the wall, so that the flux to the wall from dx is

$$d\Gamma^{b+a}(x) = (1/2)dk_{cx}^a(x) \cdot \exp\{-x/\Delta_b\},$$

where $\Delta_b = \sqrt{(\lambda_i^b \lambda_{cx}^b)/3}$. The total flux of type "b" at the wall is obtained by integrating $d\Gamma^b$ over all x , giving finally:

$$\frac{\Gamma^b}{\Gamma^a(0)} = \left(\frac{1}{2}\right) \cdot \left(1 + \frac{\lambda_{cx}^a}{\lambda_i^a} + \lambda_{cx}^a \sqrt{\frac{3}{\lambda_{cx}^b \lambda_i^b}}\right)^{-1}.$$

From this expression, it can be seen that the reflux fraction approaches 0.5 whenever $\lambda_i \ll \lambda_{cx}$, as it is for energies either below or very far above about 40 eV for D^o or T^o atoms. The reflux fraction has a minimum value of about 0.13 at about 40 eV. This simple expression agrees rather well with the results from the computer codes.

Figure 2.7-11 shows the calculated production rate of ions for three cases. In one case the plasma density is uniform across the SOL at $1 \times 10^{20}/m^3$, while the temperature increases exponentially from 10 eV at the

Figure 2.7-11. Results from the 1-D code GASREFLEX showing the ion production (ions/cm³/s) per unit of ion flux (ions/cm²/s) being neutralized at the plate. The four curves show the ions from molecules, from Franck-Condon neutrals, from the neutral charge-exchange products, and the sum of these three. Multiple charge exchange is included.

plate to 100 eV at the separatrix. In the other two cases shown the temperature and density vary in the manner indicated on the figure. The Franck-Condon atoms penetrate only a few cm, while the multiple charge exchange is clearly the mechanism that allows deeper penetration into the SOL. In all cases it can be seen that not more than 10^{-5} ionizations per initial gas atom leaving the plate occur at 15 cm, which is the closest that the separatrix comes to the divertor plate.

2.7.12 VACUUM PUMPING AT THE DIVERTOR

When the plasma is in its normal position, the divertor surface is at least 150 mm from the plasma. This allows several mean free paths (mfp) for ionization of all expected neutral species, including fast neutrals from exchange and from reflections at the plate when the electron density at the plate is $2 \times 10^{19}/\text{m}^3$ or higher. The mfps for neutrals are short even when the electron temperature at the plate is well below 10 eV because the temperature rises rapidly with distance into the plasma.

It can be seen in Fig. 2.7-12 that the gas that is produced by the neutralization of ions on most of the plate surface must pass into an increasing plasma density. This reduces the probability that a gas atom will penetrate the edge plasma and enter the confined plasma, and therefore promotes recycling. Only at the mouth of the vacuum duct is the gas directed toward decreasing plasma density and into the duct. Recycling on those flux lines that lead to the lower plate near the duct opening is provided partly by gas from nearby regions of the plate, but mostly by gas from the duct itself. Only a very small fraction of the ions entering the duct is removed by the vacuum pumps. The larger fraction enters the plasma and is recycled, causing the opening to the duct to recycle gas nearly as efficiently as the solid plate.

The particle flux at the plate can be calculated from the heat load by requiring power balance across the sheath at the plate:

$$\gamma \cdot J_t \cdot T_t \cdot \sin(\theta) = (dP/dA)_t.$$

Here, $\gamma \approx 8$ is the energy per ion carried across the sheath in units of eT_t , J_t is the ion current density at the sheath, T_t is the plasma temperature at the sheath ($T_e = T_i$ is assumed), θ is the angle between the poloidal component

Figure 2.7-12. Edge plasma ions flowing in the flux tube that passes through area A plasma at the bottom divertor. This gas is therefore recycled almost as efficiently as gas from other regions of the plate.

of the flux lines and the surface normal (see Fig. 2.7-12), and $(dP/dA)_t$ is the local heat load. If $(dP/dA)_t = 2.5 \times 10^6 \text{ W/m}^2$, $\sin(\theta) = 0.5$, and $T_t = 8 \text{ eV}$, then this power balance gives $J_t = 7.8 \times 10^4 \text{ A/m}^2$ of ions, or $J_t/e = 4.9 \times 10^{23} \text{ ions/s/m}^2$ across the sheath. This is the approximate flux of ions normal to the neutralizer plate at the entrance to the vacuum duct. The width of plate that directs neutrals into the duct is about 130 mm and extends all around the machine, giving it an area of $A_{\text{flux}} = 2.2 \text{ m}^2$. Since most of the ions that are neutralized here become gas molecules in the duct, the rate of gas input into the duct is:

$$Q_{\text{in}} = (J_t/e) \cdot A_{\text{flux}} = 1.1 \times 10^{24} \text{ atoms/s}.$$

Gas is removed from the duct through the conductance C_{pump} to the pumps, and through the opening of area A_{plasma} for gas to re-enter the plasma. The pumping speed of the plasma through A_{plasma} is

$$S_{\text{plasma}} = (1/4) \cdot v_{\text{mo}} \cdot (1 - R_{\text{plasma}}) \cdot A_{\text{plasma}}$$

If we assume that the albedo of the plasma is $R_{\text{plasma}} \approx 0.4$ (the reflux fraction is discussed in Section 2.7.10.1.) and take $A_{\text{plasma}} = 3.4 \text{ m}^2$ as shown in the figure, then $S_{\text{plasma}} \approx 640 \text{ m}^3/\text{s}$. The molecular speed is $v_{\text{mo}} \approx 1450 \text{ m/s}$ for DT at 500 K. (Notice that $S_{\text{plasma}} \gg C_{\text{pump}} \approx 25 \text{ m}^3/\text{s}$ for the lower divertor, so that it is essentially just S_{plasma} that determines the molecular density in the duct.) Particle balance requires:

$$Q_{\text{out}} = 2 \cdot n_{\text{mo}} \cdot (S_{\text{plasma}} + S_{\text{pump}}) = Q_{\text{in}} = 1.1 \times 10^{24} \text{ atoms/s}$$

which gives $n_{\text{mo}} = 8.7 \times 10^{20} \text{ mo/m}^3$ for the gas density in the duct, which corresponds to a duct pressure of 42 mTorr when the temperature is 500 K. At this high density, the gas conductance should be treated as a viscous flow rather than molecular flow as assumed here. Since the gas conductance of a given geometry is greater in the viscous regime, we design for the molecular regime to be conservative.

This pressure results from the power balance at the sheath and from the specified 8-eV electron temperature at the plate. We have some control over the plasma temperature near the plate by controlling the recycle rate, which

we can affect through the balance of gas pumping and net fueling rates. The vacuum pumps will remove gas at the rate $Q_{\text{pump}} = 2 \cdot n_{\text{mo}} \cdot S_{\text{pump}}$. With $S_{\text{pump}} = 25 \text{ m}^3/\text{s}$ at each of the two (i.e. top and bottom) divertors, this would give $Q_{\text{pump}} = 2 \times 10^{22}$ atoms/s, which is much larger than the net fueling rate of 2×10^{21} atoms/s, and even larger than the total fueling rate of about 5×10^{21} atoms/s including gas from the pellet injector. When we include the upper pumping ducts, which could remove the same amount as the lower ducts, the possible pumping rate is much greater than the total input rate. This imbalance could bleed the SOL density and reduce the duct pressure until particle balance is established. If the power flow remained constant, T_e at the plate would rise to maintain power balance, and this could cause the erosion problems that we want to avoid. Our control is through controlling S_{pump} , the pumping speed of the vacuum pumps.

We chose the total duct conductance to give $S_{\text{pump}} = 50 \text{ m}^3/\text{s}$ for the two divertors combined, in order to obtain a low base pressure of about 10^{-7} Torr to allow fast recovery after a pulse if TIBER should be operated in a pulsed mode. This vacuum pumping speed is much less than in NET, mainly because of TIBER's much smaller volume and surface area for outgassing. Both surface area and volume are small in TIBER, not only because of its smaller size but also because the region outside of the blanket and shield is in a separately maintained high vacuum. Only the inner surface and volume need be pumped through the divertors.

The vacuum pumping system in TIBER must provide the necessary pumping speed for deuterium and tritium and for helium ash in three different stages of operation: initial pumpdown, steady-state plasma operation, and pulsed-plasma operation.

2.7.12.1 Initial Pumpdown

During initial pumpdown, the gas pressure, p , in the torus decreases with time, t , during initial pumpdown as:

$$p(t) = [p(0) - Q_{\text{og}}/S]e^{-t/\tau} + Q_{\text{og}}/S,$$

where $S \approx 50 \text{ m}^3/\text{s}$, and the pumping time constant $\tau = V/S = 2.5 \text{ s}$. We use $V = 125 \text{ m}^3$, which includes only the inside of the torus. This assumes that the spaces between shielding modules and behind and around divertor modules is

pumped separately by the external system for the superconducting magnets. The gas load on the system during pumpdown, Q_{og} , is due to outgasing of surfaces. Clean stainless steel that has been baked and glow-discharge cleaned outgases at a rate of about 10^{-7} Torr $\cdot$ L/s $\cdot$ m² [3, 19]. Moderately clean surfaces before baking can have a rate that is 100 times higher. Again we assume that only the inner surfaces of the torus plus the inner surfaces of the ducts will contribute to the gas load. The inner surface area of the torus is about 200 m² and the total for the ducts and pipes is about 160 m² for a total of about 360 m² for outgasing. This gives $Q_{og} \approx 3.6 \times 10^{-5}$ to 3.6×10^{-3} Torr-liters/s, and $p(\infty) = Q_{og}/S \approx 10^{-9}$ to 10^{-7} Torr. This base pressure meets the NET goal of 10^{-7} to 10^{-8} Torr after baking and glow-discharge cleaning.

2.7.12.2 Steady State Operation

The gas load under steady-state plasma operation is just equal to the gas-equivalent of the net fueling, plus the gas that is injected into the SOL with the pellets. The net fueling current is about 350 A_a of D^o plus T^o, and the extra pellet gas is estimated to add another 420 A_a or $Q = 68$ Torr-liters/s of molecular gas. This gives a pressure at the entrances to the ducts of $p_{duct} = Q/S = 1.4 \times 10^{-3}$ Torr. In operation one would throttle the pumps to control the recycling rate and, therefore, the density of the edge plasma. TIBER's pumping speed is sufficient to allow this throttling. In steady state, the He pressure in the duct and the relative concentration $n_{He}/n_{D/T}$ of He in the edge plasma near the plates will adjust until He atoms are removed at the rate that α -particles are produced. Our estimate is that the edge plasma will contain about 5% He.

2.7.12.3 Pulsed Mode of Operation

In a pulsed mode of operation, the pressure $p(t_d)$ at the end of a dwell time t_d after a pulse is

$$p(t_d) = [p(0) - p(\infty)]e^{-t/\tau} + p(\infty) \approx p(0)e^{-t/\tau},$$

with $t = t_d$, and $p(0) \approx 5 \times 10^{-4}$ Torr is the pressure in the torus at the end of the previous pulse [3]. This gives $p(t_d) \approx 2 \times 10^{-6}$ Torr if we assume $t_d = 20$ s as for NET. This meets the NET goal of 3×10^{-5} Torr, and the 1×10^{-5} Torr design requirement used in the FEDC tokamak systems code [20].

2.7.12.4 Comparison with NET

Figure 2.7-13 shows TIBER II and NET drawn to the same scale. Although the cross-sectional areas of the plasmas are roughly equal, NET's larger R_0 makes its volume and surface areas larger, while the surface area of protective tile and of the sides and backs of modules greatly increases NET's surface area for outgassing. The main reasons for the smaller ducts in TIBER are the much smaller surface area for outgassing (because the secondary vacuum for the magnets removes the gas from between modules) and the smaller torus volume and shorter ducts. Table 2.7-6 shows the comparison.

2.7.13 IMPURITY CONTROL ISSUES

The major TIBER/ITER impurity control physics issues are:

1. Minimal impurity level ($Z_{eff} < 1.5$). Impurity level restrictions are necessary to avoid excessive radiation levels and excessive depletion due to impurities of the allowed plasma pressure. Successful experimental demonstration is required of techniques for tokamak divertor operation with low impurity content and high heating power.

2. Adequate lifetime for first wall components due to erosion. The TIBER/ITER first wall components must have a reasonable lifetime. The major source of wall erosion is probably due to sputtering by ions and fast neutrals. Wall material redeposition will ameliorate the erosion problem. It will be necessary to study and characterize erosion and redeposition phenomena using tokamak experiments and test facilities.

3. Compatibility between a compact poloidal divertor and enhanced confinement. The outstanding issues are:

- A. The minimum distance allowable from the null point to divertor plate to achieve enhanced confinement.
- B. The need for mechanical and/or plasma baffling of the neutral particle backflow.
- C. The dependence of the divertor operating parameters on the magnetic configuration, plasma density, impurity level, and radiation losses, heating power, etc.
- D. The role of impurity and hydrogen radiation in the divertor power balance.

Fig. 2.7-13. TIBER II and NET drawn to the same scale. The smaller ducts in TIBER are due to: smaller surface area for outgasing, smaller volume, and shorter ducts.

Table 2.7-6. Comparison of the vacuum pumping parameters for TIBER and for NET.

		TIBER	NET
Torus volume	m ³	125	1100
Outgas area	m ²	360	10 ⁴
Outgas load -	Torr-l/s	3.6 x 10 ⁻⁵	10 ⁻³
- after bake & glow-discharge			
Total speed	liters/s	3.5 x 10 ⁴	2.5 x 10 ⁵
Base pressure	Torr	10 ⁻⁹	10 ⁻⁸
Pump time constant	s	3.6	4.4

E. The compatibility of RF heating with divertor operation.

F. Operational limits of divertors in terms of maximum current, maximum beta, and maximum plasma density.

Some of these issues can be addressed through experiment, and some only through modeling.

4. Existence of the high recycling operation of poloidal divertor. High-density, low-temperature plasma at the divertor plate minimizes erosion and impurity production and reduces the pumping requirements for helium. High-density, low-temperature divertor plasmas have been observed on divertor experiments. Recent results on D-III-D and JET, however, show that (at least temporarily) this type of operation may not occur during "H-mode" operation due to decreased divertor particle and heat flows. The existence of high recycling modes in the TIBER/ITER should be examined using existing two-dimensional numerical models.

5. Helium exhaust. Helium pumping must be possible with a modest-sized pumping system. This requires that the neutral gas pressure be in the 10⁻³ to 10⁻² Torr range, and that the helium ion concentration in the divertor be at least 5%. Measurements of the helium concentration in the plasma center, the plasma edge, and in the divertor are needed in operating tokamaks. Part of

the gas in the divertor chamber will be energetic and non-isotropic, and this may allow a reduction of the pumping speed requirements because this hot gas will have a higher throughput than thermalized, cold gas.

6. Impurity transport. Impurity transport in diverted plasmas needs to be such as to avoid central peaking of the impurity concentration. Such central peaking is often observed with divertor experiments, in particular during "H-mode" operations. Further, there needs to be sufficient transport for the central helium ions to ensure an adequate helium removal rate. Experiments are needed which can characterize impurity transport behavior.

7. Acceptable tritium retention in divertor and first wall materials. The quantity of absorbed tritium is limited by the allowed tritium inventory. The tritium absorption properties of the candidate wall materials need to be evaluated, and the likely tritium absorption in ITER needs to be estimated in light of the recycling experience on large, high power tokamaks.

8. Impurity control under the special conditions required for current drive. Current ramp-up and transformer recharge often require special plasma conditions, such as very low densities, to be efficient. At such low densities it is not likely that a high recycling divertor will work.

Numerical modeling of divertor behavior can be applied to many of these issues. The existence of the high recycling mode can be demonstrated with zero and one-dimensional fluid models. The two dimensional PLANET/DEGAS divertor model could also be applied, using realistic geometrical descriptions of proposed TIBER/ITER impurity control designs, to estimate and optimize:

- A. The degree of particle recycling (none, high, or super- high) and consequent plate erosion and tritium inventory.
- B. Fuel-particle containment time in the divertor, which may affect achieving H-mode τ_E scaling.
- C. DT and He pumping performance, which is crucial for steady state operations.
- D. The ability of the divertor to contain and remove heavy impurities.

The last two issues would involve an upgraded version of PLANET, which includes impurity transport.

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2.8 FUELING*

2.8.1 INTRODUCTION

A suitable fueling scheme plays an essential role in sustaining both a steady-state and a feasible long-pulsed fusion reactor. In this work, we assess the requirements for fueling a high density, high temperature fusion-grade plasma for the TIBER Engineering Test Reactor design. Two main approaches of fueling are studied here, namely, pellet injection and coaxial compact-torus plasmoid injection.

We note here that some recent experiments have observed anomalous inward transport of density from pellets deposited in the outer regions of the plasma [1], so it may not be necessary to fuel the plasma at the center. However, we prefer to have the capability of depositing fuel density in a controlled manner at any point in the machine that enables us to control the density profiles. The importance of profiles is stressed in a recent review paper [2]. Deep central fueling is also desirable for reducing the plasma-wall interactions and DT gas recycling near the wall, enhancing particle and energy confinement times [3,4], and the production of appreciable bootstrap current. We are motivated by all these advantages to seek a deep-penetration central fueling scheme.

2.8.2 PELLET FUELING

Pellet injection has received much attention for tokamak fueling partly because of the attractiveness of the possibility of depositing fuel inside the plasma relative to edge gas puffing and some successes in recent experiments [5,6]. However, there are still many constraints and uncertainties in the pellet fueling scheme that have to be overcome or resolved in the future. Some critical issues of pellet fueling are discussed in this section. An assessment of the fueling requirements for the TIBER ETR is studied. A baseline of pellet fueling scheme for the TIBER design is then calculated. As we will show, conventional pellet fueling will probably not be adequate for deep-penetration fueling of reactor-grade plasmas such as the ETR.

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* Additional discussion of engineering aspects of pellet injection including fueling systems for start-up (low density and temperature conditions) can be found in Section 3.

2.8.2.1 Pellet Ablation Modeling

The neutral shielding model [7,8] so far has good agreement with past pellet injection experiments [5,6], which operated at a low plasma temperature and density relative to a fusion reactor. In a reactor-grade plasma, in addition to the plasma electrons, fusion-born alpha particles may also contribute to the pellet ablation. The low alpha density is partially compensated by the high alpha energy, which makes the alpha heat flux an appreciable fraction of that of the electron.

In a separate analysis [9], we have studied the effect of the alphas in a pellet ablation modeling based on the neutral shielding theory. We follow the approach of Ref. [7] in solving the two ablating species eigenvalue problem. The fast alpha is characterized by an averaged energy obtained from a slowing-down velocity distribution weighted favorably with velocity, as in Ref. [10]. The mean energy of the alphas is in the range of 1.6 to 1.9 MeV, which results in an alpha penetration distance comparable with 10-keV electrons. We find the alphas add a fairly constant and small enhancement to the pellet ablation rate in the range of T_e from 10 to 30 keV for a given alpha density. For T_e below 10 keV, the ablation is dominated by the alphas such that the ablation rate is determined mainly from the alpha density. However, the alpha density is insufficiently low in this range of T_e to create any drastic increase in the ablation rate.

The modified pellet ablation scaling law including the fast alphas is given by [9]:

$$\frac{dr_p}{dt} = \begin{cases} 1.72 \times 10^{-8} \left(1 + Ca \frac{n_\alpha}{n_e}\right) r_p^{-2/3} T_e^{1.64} & \text{for } 10 \text{ keV} \leq T_e \leq 30 \text{ keV} \\ 6.24 \times 10^{-2} \left(1 + Ca \frac{n_\alpha}{n_e}\right) r_p^{-2/3} n_e^{1/3} \left(\frac{T_e}{10^4 \text{ eV}}\right)^{X_\alpha} & \text{for } T_e < 10 \text{ keV} \end{cases} \quad (2.8-1)$$

where Ca is 5.0 (2.8) for $\frac{n_\alpha}{n_e}$ less (greater) than 0.03, $X_\alpha = 1.64 \exp(-50 \frac{n_\alpha}{n_e})$, and the units are given by r_p in cm, n_e in cm^{-3} , and T_e in eV. The above expression for the ablation rate is derived for a hydrogen pellet. It can be shown that the ablation rate is proportional to $\frac{m_0^{2/3}}{\rho}$, where m_0 is the atomic mass and ρ is the solid density. Therefore, the ablation rate should be

reduced by factors of 0.68 and 0.57 for deuterium and tritium pellets, respectively. Following the classical slowing-down theory to calculate the fast alpha population, we find the alphas provide about 5% enhancement of the ablation in the range of TIBER plasma parameters.

It should be noted that various suggestions have been made to improve the neutral shielding pellet ablation model, especially for high temperature and density plasma. Extra shielding effects considered include magnetic field distortion [11], ionization and dissociation process [12], and cold plasma shielding [13-15], which reduce the ablation rate. On the other hand, using a multi-group calculation of the electrons [15] instead of the conventionally used monoenergetic representation increases the pellet ablation. To a crude approximation, the two effects of increasing and reducing the pellet ablation rate cancel each other. The accuracy of the various current models of pellet ablation has to be determined from future experiments operating in higher plasma temperature and density regimes.

2.8.2.2 Fueling Requirements and Design Baseline

The required fueling current to sustain a constant fuel in the plasma is calculated from the convective loss, fusion burn, and neutral beam injection currents. Since the current neutral beam program development does not intend to use tritium, we assume that the beam contains only deuterium atoms. Thus, the pellet will have different compositions of deuterium and tritium in order to maintain an equal mixture of DT in the plasma if neutral beam injection is used for current drive and/or heating. The fueling currents for deuterium and tritium are given as

$$I_D = I_{\text{burn}} + 0.5 I_{\text{conv}} - I_{\text{beam}} \quad (2.8-2)$$

$$I_T = I_{\text{burn}} + 0.5 I_{\text{conv}} \quad (2.8-3)$$

where I_{burn} , I_{conv} , and I_{beam} are the fusion burn, convective loss, and neutral beam currents, respectively. Hence the composition mixture of the pellet is given by:

$$F_D = I_D / (I_D + I_T) ,$$

$$F_T = I_T / (I_D + I_T) ,$$
(2.8-4)

Next, we have to decide how much particle content is to be installed in a single pellet. It is clear from Eq. (2.8-1) that a large pellet will have a better penetration, and also it requires a small injection frequency. However, the total number of injected particles in a pellet, N_p , should be kept at a small fraction of the total number of particles N contained in the plasma in order to avoid a large density perturbation, pulsing of the fusion power generated, and the possibility of MHD and thermal instabilities. We adapt the commonly used value of $\frac{N_p}{N} \leq 0.1^6$. Hence we calculate the mass and volume of the pellet from

$$m_p = \frac{N_p}{6.023 \times 10^{23}} (2 F_D + 3 F_T) ,$$
(2.8-5)

$$Vol_p = \frac{m_p}{0.202 F_D + 0.310 F_T} .$$
(2.8-6)

The required repetition rate of pellet injection to maintain a constant fuel level is then given by

$$Rate = \frac{I_D + I_T}{1.6 \times 10^{-19} N_p} ,$$
(2.8-7)

where 1.6×10^{-19} is the charge on the electron e . Alternatively, for no neutral beam injection, the repetition rate can also be found from the particle confinement time τ_p ,

$$Rate = \frac{N}{\tau_p N_p} .$$
(2.8-8)

The most restrictive requirement for pellet fueling in a reactor-grade plasma is the necessity to achieve very high injection velocity to reach the desired penetration. We demonstrate this problem by calculating the pellet

penetration distance with the pellet injection velocity for both the steady-state and pulsed-inductive modes of operation of TIBER (see Sections 2.5 and 2.6 above). The penetration distance is computed by integrating Eq. (2.8-1), in which the local density and temperature are obtained from the radial profiles of the form

$$n_e(r) = n_{eo} \left(1 - \left(\frac{r}{a}\right)^2\right)^{a_n},$$

$$T_e(r) = T_{eo} \left(1 - \left(\frac{r}{a}\right)^2\right)^{a_T},$$

where the exponents a_n and a_T are 1.02 and 0.58, respectively, in the current TIBER baseline; n_{eo} and T_{eo} are the peak axial values; and a is the minor radius.

Figures 2.8-1 and 2.8-2 show the penetration curves for three different sizes of pellets for the two modes. The injection velocity has to be very large (> 10 km/s) for any appreciable penetration. The penetration is better for a larger size pellet (see Eq. (2.8-1)), but we cannot increase the size to an extent of obtaining a satisfactory penetration since the pellet size is limited by the above mentioned criterion of $\frac{N_p}{N} \leq 0.1$. The dotted curves are with pellets of sizes corresponding to this criterion. We see that the pellet penetrates better in the pulsed mode plasma because of a lower electron temperature, despite the increase in electron density. As we may recall from Eq. (2.8-1), the ablation rate has a more sensitive dependence on the electron temperature than the density. It has also been suggested that a sabot (high-Z material coating) for the pellet could be used to overcome this insufficient penetration problem. However, we do not intend to adopt this scheme since it is not likely that this scheme can have significant improvement of the penetration without creating an unacceptable level of impurity from the sabot.

As mentioned in Section 2.8.1, it is not clear whether it is required to deposit the fuel on x axis for central fueling. Some experimental evidence indicates that fuel penetration to $1/3$ to $1/2$ of the plasma radius is perhaps sufficient with the aid of the anomalous inward pinch effect. As a result of the stringent injection velocity requirement, we only benchmark the pellet

Figure 2.8-1. Penetration distance vs pellet injection velocity for various pellet sizes for the steady-state mode with $T_{eo} = 24.7$ keV and $n_{eo} = 2.14 \times 10^{20} \text{ m}^{-3}$.

Figure 2.8-2. Penetration distance vs pellet injection velocity for various pellet sizes for the pulsed-inductive mode with $T_{eo} = 17.6$ keV and $n_{eo} = 2.82 \times 10^{20} \text{ m}^{-3}$.

fueling scheme with penetration of 1/3 and 1/2 the plasma radius. In the steady-state mode, this results in corresponding injection velocities of 22 km/s and 52 km/s, respectively.

We emphasize that "penetration" to a given fraction of the minor radius means, of course, that the pellet has completely ablated away at this point; its fuel mass has been smeared over the ablation path up to this radius. A fuel deposition profile for halfway pellet penetration with the 52 km/s injection is shown in Fig. 2.8-3. The deposition is small in the plasma periphery because the ablation rate is low, due to a smaller electron temperature and density and a larger pellet radius. On the other hand, even though the ablation rate is high deep inside the plasma, the deposition is also small because the pellet radius is very small; accordingly, the particle content in a given radial displacement is less. Therefore, we expect the maximum deposition to be somewhere in between the plasma surface and the deepest penetration point, as indicated in Fig. 2.8-3.

Due to the low inertial strength of DT ice pellet, the acceleration of the pellet must be limited [16] to a pressure of about 0.5 MPa, or about 3 MPa if bounded by a guide tube, otherwise pellet fracture will occur. This necessitates very long acceleration flight tubes for high velocity injection. An approximation expression for the minimum launch length required is given in Ref. [16] as:

$$L = 2.5 \times 10^{-8} \rho r_p^2 v_p^2, \quad (2.8-9)$$

where r_p and v_p are the pellet radius and injection velocity, respectively. For example, even if we could produce the 50 km/s injection velocity for penetration to half the minor radius in the above TIBER design, the minimum length of the pellet injector would be about 250 m!

In summary, we have studied the various pellet fueling requirements for the TIBER design. A pellet fueling design baseline is recommended for each of the steady-state and pulse inductive operation modes as shown in Table 2.8-1. However, as we emphasize in the next section, it is unlikely that conventional pellet acceleration techniques will be able to meet the high injection velocity requirements.

Figure 2.8-3. Fuel deposition profile for the steady-state baseline with halfway penetration from a 50 km/s injection velocity.

Table 2.8-1. Fueling and related plasma parameters for the TIBER design.

	Steady-state	Pulsed-inductive
Major radius, R_0	3 m	3 m
Minor radius, a	0.834 m	0.834 m
T_{eo}	24.7 keV	17.6 keV
T_{io}	25.2 keV	18.0 keV
N_{eo}	$2.14 \times 10^{20} \text{ m}^{-3}$	$2.82 \times 10^{20} \text{ m}^{-3}$
$\langle T_e \rangle$	19.2 keV	13.7 keV
$\langle T_i \rangle$	19.6 keV	14.0 keV
$\langle n_e \rangle$	$1.06 \times 10^{20} \text{ m}^{-3}$	$1.40 \times 10^{20} \text{ m}^{-3}$
Axial magnetic field B_0	5.55 T	5.55 T
Net fueling current required ^a	156.7 A	198.1 A
Fuel mix ratio required ^a	22% D / 78% T	50% D / 50% T
Maximum fuel load per shot ^b	4.8 mg (168 A)	4.7 mg (164.5 A)
Pellet diameter	3.2 mm	3.5 mm
Minimum repetition rate ^c	0.93 Hz	1.12 Hz
Injection velocity required ^d :		
a/3 penetration	22.2 km/s	12.4 km/s
a/2 penetration	52.4 km/s	29.0 km/s
Launch length required:		
a/3 penetration	35 m	10 m
a/2 penetration	250 m	65 m

^a87.3 A of deuterium are separately supplied by the neutral beam current-drive system.

^bMaximum allowed fuel load is set at 10% of the plasma content.

^cAt maximum fuel load per shot.

^dAt maximum pellet acceleration pressure of 3 MPa.

2.8.3 SURVEY OF FUEL INJECTION SYSTEMS

In this section, we present a brief survey of the various current and proposed fueling schemes. Table 2.8-2 lists seven possible candidates for fuel injectors, together with the status, projected ultimate velocity, estimated cost, and limitation and uncertainty of each.

The pneumatic gun, centrifugal injector, and railgun all have been operational and, particularly, the former two injectors have achieved great success in pellet injection experiments. However, these injectors are limited by their inability to reach high enough injection velocity (> 10 km/s) to be considered adequate for the TIBER design. As a result, other conceptual ideas have been suggested to overcome this problem. The electron beam (or laser) rocket [17,18] and the theta-pinch accelerator [19] concepts utilize an ablated portion of the pellet and a theta-pinch configuration plasma, respectively, as propellant to accelerate the target pellet. Both schemes in principle can achieve injection velocity in the range of 10 to 20 km/s, which is only marginal for the TIBER pulsed-inductive mode, but not quite sufficient for the steady-state mode. The laser-vapor-jet [20] concept has been computed to reach an injection velocity of about 50 km/s by a simple assessment of the scheme. The idea is to deliver a laser pulse to the pneumatic gun pre-accelerated pellet at the interface with the tokamak plasma. The resulting plasma plume expands and sweeps through the tokamak plasma until it is slowed down by the magnetic field pressure gradient. Further information on the conceptual laser-vapor-jet fueling scheme can be found in Section 3 of this report.

The laser-vapor-jet concept seems to be the only pellet injection scheme that can achieve deep penetration in TIBER. However, it imposes a severe technological constraint in utilizing a rep-rateable 15-kJ laser. Moreover, all the conceptual advanced pellet fueling schemes have many uncertainties that need further investigations and probably over a relatively long R&D time period. Also, all pellet injection schemes suffer from the weak tensile strength of the pellet that requires a fairly long acceleration path for high velocity pellet. We conclude that pellet injection will probably not be adequate for the TIBER ETR design.

To continue the quest for depositing fuel atoms in a controlled manner at any point inside the plasma despite the deficiency of pellet injection, we pursue an alternative fueling scheme of shooting compact torus (CT) plasma

Table 2.8-2. Brief survey of fuel injection systems.

Type of injector	Status	Achieved velocity (km/s)	Projected ultimate velocity (km/s)	Estimated cost (m\$)	Limitation/uncertainty
Pneumatic gun	Operational	1 to 2	2 to 3 ^a	~5	Temperature/pressure of propellant gas
Centrifugal injector	Operational	<1	1.5 ^b	~5	Mechanical stress of arbor
High-temperature two-stage gas gun	Operational	3.4	5 to 10	~5	Barrel erosion of pellet, pellet tumbling, requires sabot
Railgun	Operational	2	7 to 8 ^c	~15	Tensile strength of pellet, back arcing
E-beam (laser) rocket	Conceptual	---	10 to 15 ^a	---	E-beam (laser) plasma interaction, shock wave, tensile strength of pellet
Theta-Pinch accelerator	Conceptual	---	15 to 20 ^a	---	Tensile strength of pellet, plasma interactions
Laser Vapor-jet injector	Conceptual	---	50 ^d	15 to 20	Large laser power, vapor-plasma interaction
Coaxial plasma Gun	Operational	>1000	800 ^e	15	Plasma interactions (reconnection/tearing/disruptions)

^a"Wish" limit.^bInertial limit of flywheel material.^cPractical limit.^dComputed limit (Logan, 1986).^eRequired for TIBER central fueling.

rings accelerated by a coaxial plasma gun. This possible application of a collective accelerator based on magnetically confined plasma rings was first suggested by Hartman and Hammer in 1981 [21]. An analysis and application of this CT injection scheme is present in the next section. In our opinion, this method has one outstanding advantage relative to the other advanced schemes in Table 2.8-3: namely, a currently operating experimental system with performance parameters in the range required for future reactor fueling.

2.8.4 FUELING WITH ACCELERATED COMPACT TORUS

The compact torus is a self-contained, toroidally-shaped entity comprised of embedded poloidal and toroidal fields forming closed surfaces. In the present RACE (Ring ACcelerator Experiment) configuration of Hartman and Hammer at Lawrence Livermore National Laboratory [22], the toroids are formed by a magnetized coaxial Marshall gun. The gun is fired following the establishment of an initial magnetic field. The resulting plasma, with embedded toroidal accelerating field, emerges from the gun and stretches the initial field within a flux conserving region. Reconnection of this initial field produces a dipole-like, poloidally-directed field and results in an isolated, trapped plasma ring. In typical rings, the internal poloidal and toroidal fields are of comparable magnitude.

The CT ring so formed has three interesting features: (a) the lifetime of the magnetic topology of the ring is sufficiently long that we can apply an acceleration over distances much greater than the characteristic ring dimensions, (b) the plasma confinement time generally exceeds the magnetic lifetime, and (c) the CT is a very robust structure and is able to sustain large accelerating forces, comparable to those required to maintain radial equilibrium, while retaining its structural integrity. Accordingly, the relatively large mass per ring ($\sim$ milligrams) coupled with a high acceleration results in a structure with directed kinetic energy appreciably greater than its stored magnetic or thermal energy.

In principle, there are several ways to provide an accelerating magnetic force to a compact toroid. Currently, Hartman and Hammer [22] in the RACE system at LLNL are employing a coaxial railgun where the ring acts as a sliding short and the magnetic pressure behind the ring provides the $\underline{J} \times \underline{B}$ accelerating force. This coaxial mode of acceleration is simply an extension of the formation technique within the magnetized, coaxial Marshall gun.

During its passage along the coaxial barrel, radial expansion of the toroid is prevented by a radial equilibrium force generated by image currents in the conducting walls.

A review of the experimental data for the production and acceleration of compact toroids up to 1984, is given in Ref. [23]. In particular, force free, low-pressure compact toroids were reported at a number of establishments, including Los Alamos National Laboratory where peak electron temperatures of 150 keV and lifetimes greater than 1 ms (much larger than that required for a fueling system--see below) were obtained. In addition, translation experiments at LANL and at Kurchatov in the Soviet Union have demonstrated velocities of ~ 200 km/s. In the current RACE experiments of Hartman and Hammer at Lawrence Livermore National Laboratory [22], compact toroids are produced with 50-cm outer diameter. Masses of 0.5 mg have been accelerated to velocities of ~ 400 km/s, while velocities of over 2500 km/s have been achieved. At a velocity of 1400 km/s and a ring mass of 0.018 mg, approximately 18 kJ of the capacitor-bank energy of 50 kJ were converted into kinetic energy, resulting in a wall-plug to injected efficiency of $\sim 30\%$. The total capacitor bank energy available for these experiments is 260 kJ, although full use has not been made of this to date.

Clearly, from the current experimental database, we are confident that the high force per unit mass obtainable with the coaxial accelerator, its potential high efficiency ($< 50\%$), and the large number of particles in the CT provide for a controllable high velocity fueling system with potential outputs in the multi-megajoule range.

2.8.4.1 Toroid-Plasma Dynamics

Generally, there are six interrelated constraints we must consider to determine the requirements for transporting the fuel in the compact toroid from the exit of the coaxial plasma gun to the desired deposition point in the plasma:

1. The lifetime of the CT must be long relative to the time for its production, acceleration, and penetration in the plasma.
2. Once the ring has left the coaxial guide tubes of the railgun it will tend to tilt through 90° in the external toroidal field of the tokamak (analogous to a galvanometer coil turning in an applied

field). The poloidal field of the CT then reconnects with the toroidal field of the tokamak, and the CT configuration rapidly diffuses. This is shown schematically in Fig. 2.8-4. The tilting time and the time to reconnect a large fraction of the ring flux may differ in general, due to possible constraints on reconnection rates imposed by plasma conductivity. This will be considered later. For the purpose of this work, however, we assume plasma conductivity does not impede reconnection. Note also that reconnection does not necessarily mean that the fuel stops at this point because the forward momentum pressure may be sufficient to overcome the tension of the external toroidal field (see below). However, although in this case the fuel plasmoid will inwardly distort the embedded reconnected field, unless tearing occurs the fuel will be deposited only on the field lines where reconnection occurred. Accordingly, the injection velocity must be such that the reconnection following tilting occurs near the desired fuel deposition point.

3. Notwithstanding the tilt/reconnection requirement above, the initial velocity must be sufficiently high to overcome the retardation due to the toroidal magnetic field gradient in the tokamak. In this case, the injection velocity would be determined by the requirement that the final CT velocity be zero at the desired deposition point.
4. If the CT were to be injected into free space, it would expand approximately spherically at a rate determined by its internal Alfvén speed. This is because the radial equilibrium force supplied by the image currents in the coaxial guide has been removed. Within the tokamak, however, the ring will expand or contract depending on the balance between the internal ring pressure and the external pressure (magnetic and thermal). Typically, as we shall see, the ring will contract in the direction perpendicular to the external toroidal field while expanding in the direction parallel to this field. Note that these changes in dimensions (and, therefore, by flux conservation, changes in the internal CT field) must be considered consistently in the calculation of tilting and penetration above.

Figure 2.8-4. Schematic of compact-toroid (CT) acceleration and interaction with the external tokamak toroidal field. As it traverses the external field, the CT rotates through a 90° angle due to the torque on its internal magnetic moment. Alignment of the CT poloidal field with the external field results in field-line reconnection and fuel deposition.

5. To avoid the possibility of plasma disruptions or large, shockwave-induced hydrodynamic drag, we would expect that the CT injection velocity would have to be less than the characteristic Alfvén velocity of the plasma we are fueling.
6. Although unity beta compact toroids have been formed in field-reversed theta-pinch experiments at Los Alamos National Laboratory [24], ballooning calculations for spheromak type, small Larmor radius compact toroids, such as those produced in the RACE experiment at LLNL, have determined that rings are internally unstable for betas much above 20%. Consequently, this may limit the amount of mass which can be carried in a CT of given internal magnetic field and temperature.

Accordingly, armed with these constraints and conditions above, we can proceed to determine typical ring parameters. In the following formalism, base SI units are used throughout, except where noted.

The ultimate lifetime of the CT, tilting and reconnection notwithstanding, will be determined by the resistive decay of its internal fields. Typically, the time for a magnetic field to soak through a characteristic distance ℓ is

$$\tau = \mu_0 \ell^2 / \eta \quad (2.8-10)$$

where η is the classical Spitzer resistivity multiplied by a neoclassical enhancement factor, γ_{nc} , as

$$\eta = 1.65 \times 10^{-9} \frac{\gamma_{nc} Z_{eff} \ell n \Lambda}{T_e^{3/2} (\text{keV})} \quad \Omega \cdot \text{m} . \quad (2.8-11)$$

Substituting into Eq. (2.8-10) and replacing ℓ by the radius r of the ring, we obtain a CT lifetime of

$$\tau = \frac{7.62 \times 10^2 r^2 T_e^{3/2} (\text{keV})}{\gamma_c Z_{eff} \ell n \Lambda} . \quad (2.8-12)$$

From observations of the RACE coaxial gun [22], T_e is in the range 10 to 100 eV; we will use 25 eV in subsequent calculations. In addition, we require clean rings for fueling which have Z_{eff} close to unity (see later).

Accordingly, for a typical ring of 0.1-m radius carrying 5 mg of DT fuel, the decay lifetime is in the millisecond range. Because this is much larger than characteristic times for acceleration and subsequent slowing down in the plasma (approximately several microseconds--see below), we can ignore ring decay in the following calculations.

Next we consider the tilt constraint. Because the radius of curvature of the toroidal field of the tokamak is large relative to the characteristic size of the ring, we can model the former as an infinite cylinder but retain the field gradient resulting from the finite aspect ratio. The CT leaves the coaxial gun with its poloidal field perpendicular to the toroidal field of the tokamak--see Fig. 2.8-4. As it slows down, the CT rotates under the torque as a result of the interaction of its magnetic moment with the external toroidal field, B_ϕ , i.e.,

$$\Gamma(s) = \underline{\mu}(s) \times \underline{B}_\phi(s) = \mu(s) B_\phi(s) \sin\theta(s) , \quad (2.8-13)$$

where θ is the angle between the poloidal field of the ring and the toroidal field of the tokamak (see Fig. 2.8-4) and where the torque is evaluated at a given coordinate s within the plasma measured from the exit of the coaxial gun. The corresponding angular acceleration of the ring at this point is

$$\alpha(s) = \frac{\Gamma(s)}{I_m(s)} , \quad (2.8-14)$$

where I_m is the moment of inertia of the ring about the rotation axis. We can write the following approximation for μ as

$$\mu = IA \approx \frac{1.36}{\mu_0} R \langle B \rangle A , \quad (2.8-15)$$

where I is the current in the CT, A is the enclosed area, and $\langle B \rangle$ is the mean field of an equivalent circular coil radius R . Replacing $\langle B \rangle$ and R by the internal poloidal field B_{ring} and radius r of the ring, respectively, and substituting into Eq. (2.8-14) gives the angular acceleration as

$$\alpha(s) \approx \frac{1.36 r(s) B_{\text{ring}}(s) A(s) B_\phi(s) \sin\theta(s)}{\mu_0 I_m(s)} . \quad (2.8-16)$$

Note that because the dimensions of the ring vary as it traverses the plasma, all variables in Eq. (2.8-16) must be referenced to the local plasma coordinates. Therefore, in a time δt , at s , the ring rotates through an angle $\delta\theta$ about θ where

$$\delta\theta(s) = \omega_0(s) \delta t + \frac{1}{2} a(s) (\delta t)^2, \quad (2.8-17)$$

and where $\omega_0(s)$ is the initial angular velocity at θ . Accordingly, for a given injection velocity and a concurrent slowing down model (see below), we can determine the position s in the plasma where the CT has rotated through a full 90° so that the CT poloidal field aligns and reconnects with the toroidal field of the tokamak.

Now should the energy density of the ring, $1/2 \rho(s) v_{\text{ring}}^2(s)$, at reconnection be greater than that of the external field, $B_\phi^2(s)/2\mu_0$, then the forward momentum of the plasmoid will bend the reconnected field lines inward. Thus after reconnection we can set up an instantaneous equation of motion for a small displacement δs , which can be used to determine the final stopping point of the ring, i.e.,

$$v^2(s) = U^2(s) + 2a(s) \delta s. \quad (2.8-18)$$

Here, $U(s)$ is the initial velocity of the ring at s , and $a(s)$ represents the deceleration due to the fact we are losing the kinetic energy of the ring and storing it the distorted field, i.e.,

$$a(s) \sim - \frac{B_\phi'^2(s) A(s)}{2\mu_0 m}, \quad (2.8-19)$$

where m is the mass of fuel in the ring and where $B_\phi'(s)$ is the local value of the distorted toroidal field at the reconnected plasmoid; note that we use a prime to distinguish this compressed field from the local value of the unperturbed toroidal field $B_\phi(s)$. Modeling of this effect will be discussed in the next section. Notice however, that unless resistive tearing and further reconnection occurs, then, although the toroidal field may be distorted inward, the fuel remains tied to these fieldlines. This means that the fuel deposition point averaged around the circumference of the machine is

merely the radial position at which 90° flipping and initial reconnection occurred. If indeed tearing does occur, this may have important implications for disruption initiation. However, it would mean that fuel would be deposited at radii inside the 90° reconnection surface.

To compute the ring dynamics before reconnection occurs, we note that the instantaneous dimensions of the ring depend on its expansion or contraction both parallel and perpendicular to the toroidal field. In the perpendicular direction, we can determine this by the difference between the external and internal pressures. Accordingly, in the paraxial limit:

$$\Delta P(s) \sim \frac{B_\phi^2(s)}{2\mu_0} + P(s) - \frac{B_{ring}^2(s)}{2\mu_0} - P_{ring}(s), \quad (2.8-20)$$

where P and P_{ring} are the external plasma and internal ring pressures, respectively. The corresponding velocity of contraction (or expansion) at this point becomes

$$\nu(s) = 2\Delta P(s)/\rho(s), \quad (2.8-21)$$

where $\rho(s)$ is the mass density of the ring; this velocity should not be confused with the forward velocity $V(s)$ of the ring at this point.

Parallel to the toroidal field, we have

$$\Delta P(s) \sim \frac{B_{ring}^2(s)}{2\mu_0} + P_{ring}(s) - P(s), \quad (2.8-22)$$

with a corresponding velocity of the same form as Eq. (2.8-19). Note that in the limit the plasma pressures are small relative to the ring field (a reasonably good assumption in most of our cases), then this expansion velocity parallel to the field would reduce to the local Alfvén velocity, i.e.,

$$\nu(s) \sim 1.38 \times 10^{16} B_{ring}(s)/\sqrt{n_{ring}(s)}, \quad (2.8-23)$$

Equations (2.8-22) and (2.8-23) possibly overestimate the expansion parallel to the field since they neglect the retarding effect of magnetic tension of the ring lines of force.

As it travels through the plasma, the ring will undergo negligible heating by the external electrons and alphas because the thermalization time of these particles is much greater than the penetration time of the ring. Consequently, we neglect external heating of the ring plasma and model the ring as a cylinder with axis perpendicular to its direction of motion and local dimensions determined by the above pressure balance relations.

Next, we represent the toroidal field of the tokamak in terms of the magnetic field on axis, B_0 , as

$$B_\phi(R) = B_0 R_0/R ,$$

or, in terms of s

$$B_\phi(s) = B_0 R_0 / (R_0 + a_0 + \Delta - s) , \quad (2.8-24)$$

where R_0 and a_0 are the plasma major and minor radii, respectively, Δ is the distance from the exit of the coaxial gun to the edge of the plasma, and R and s represent the local plasma coordinate measured from the center of the tokamak and from the exit of the gun, respectively. We note that the forward motion of the CT is retarded by the toroidal field gradient with a deceleration term of the form

$$a(s) = - \frac{P \cdot \frac{dA}{ds}}{m} , \quad (2.8-25)$$

where m is the mass of fuel and where the integral is evaluated over the curved surface of the cylindrical ring. In terms of the radius $r(s)$ and length $\ell(s)$ of the ring at this point, this becomes

$$a(s) = - \frac{r(s) B_0^2 R_0^2 \ell(s)}{2\mu_0 m} \int_0^{\pi/2} \frac{\cos\xi}{[R(s)-r(s)\cos\xi]^2} - \frac{\cos\xi}{[R(s)+r(s)\cos\xi]^2} d\xi, \quad (2.8-26)$$

where $R(s)$ is the radial position coordinate measured from the center of the machine which, from Eq. (2.8-24) above, can be expressed in terms of s as $R(s) = R_0 + a_0 + \Delta - s$, and where ξ is the angle variable for surface integration.

Accordingly, the ring velocity $V(s)$ at any coordinate s before the 90° reconnection point becomes

$$V^2(s) = V_{inj}^2 + 2 \int a(s) ds, \quad (2.8-27)$$

where V_{inj} is the initial injection velocity.

Reference to Eqs. (2.8-17) and (2.8-27) shows that if the torque on the CT is small it will slow down and stop in the plasma at a point determined by the initial injection velocity V_{inj} as expressed by Eq. (2.8-28). By contrast, if the tilting moment is high, the CT will rotate through 90° and reconnect before its forward velocity has been reduced to zero by the field gradient. At this point, as discussed above, bending of the reconnected toroidal field will occur if the forward momentum pressure is sufficient; however, fuel will be deposited at the flux surface where reconnection occurs. Clearly, in the latter case, to ensure that the ring reconnects at the desired fueling point, the initial velocity requirement V_{inj} will be greater than that obtained from the solution of Eq. (2.8-27) because the ring still has some forward velocity at the reconnection point. Therefore, the injection velocity from Eq. (2.8-27) is the minimum required, the actual value being determined by concurrent solution of Eq. (2.8-17) for a rotation through $\theta = 90^\circ$.

Finally, we have two further constraints from above to consider. First, to guard against the possibility of induced disruptions, we limit the ring velocity to less than the local plasma Alfvén velocity, i.e.,

$$V_{ring}(s) < 1.38 \times 10^{16} B_\phi(s) / \sqrt{n_{plasma}(s)}. \quad (2.8-28)$$

Second, the ring beta at the exit of the coaxial gun (i.e., where $s = 0$) must be limited to around 20% or less, i.e.,

$$2 n_{ring}(0) kT < 0.2 B_{ring}^2(0) / 2\mu_0. \quad (2.8-29)$$

2.8.4.2 Modeling and Results

In this section, we apply the formalism above to fueling the TIBER Engineering Test Reactor. The fueling related plasma parameters for TIBER were shown above in Table 2.8-1. Here, we switch from the pellet injection scheme to the CT injection scheme to solve the same fueling requirements, with

on-axis deposition of the CT. Note that we only do the calculations for the steady-state case because both the steady-state and pulsed cases have the same geometry and magnetic field, but with different plasma density and temperature, which contributes to a difference only of order β .

In order to simplify the solutions of the above coupled non-linear equations describing the motion of the CT, we have developed a multi-zone model to compute the CT injection into a tokamak plasma. The plasma is modeled as an infinitely long cylinder, but with the retention of the magnetic gradient due to the finite aspect ratio. The trajectory of the CT is computed from the lower bound of a zone to the corresponding upper bound using averaged quantities within that zone. As we follow the motion of the CT in each zone, we calculate in order:

- (1) The expansion/contraction of the CT due to the difference between its internal pressure and the total external pressure (Eqs. (2.8-20) through (2.8-22 above)),
- (2) The deceleration of the translational motion across the tokamak toroidal field lines as a result of the field gradient (Eqs. (2.8-23) through (2.8-27 above)), and
- (3) The rotation of the tilt angle of the CT from the interaction of its magnetic moment with the tokamak magnetic field (Eqs. (2.8-13) through (2.8-17 above)).

The contraction (expansion) of the CT influences the motion of the CT in several ways. First, it affects the deceleration in the translational forward motion, because the decelerating force is proportional to the projected cross-sectional area of the CT (Eq. (2.8-25)). Second, it affects the tilting rate of the CT. The increase (decrease) of the toroidal and poloidal magnetic field of the CT is determined by the contraction (expansion) of the CT since the internal toroidal and poloidal magnetic fluxes are conserved. The change in the dimensions of the CT also varies its moment of inertia. Consequently, by Eq. (2.8-16) the angular acceleration for tilting is similarly affected, being proportional to the internal magnetic field strength and inversely proportional to the moment of inertia.

In the following calculations, we use the maximum allowable ring beta of 20% (Eq. (2.8-29)) in order to minimize the CT injection velocity requirement, as discussed later. The mass and dimensions of the CT are varied parametrically with a fixed representative temperature of 35 eV (see above).

We take the initial length of the CT as 2/3 of its initial diameter. The plasma conditions shown in Table 2.8-1 are used for all subsequent calculations.

First, for illustrative purposes, we calculate the injection velocity requirement for depositing the CT at the center of the plasma, ignoring, for the moment, the effect of fieldline reconnection. Figure 2.8-5 shows the variation of the parameters of interest for a 5-mg, 30-cm initial diameter CT as it penetrates the plasma. The initial velocity is chosen such that the CT stops at the center of the plasma. We stress that the effect of fieldline reconnection is neglected in Fig. 2.8-5 merely to demonstrate the retarding effect of the toroidal field gradient. In Fig. 2.8-6 we include fieldline reconnection so that the initial velocity must now be such that the CT has rotated through 90° and reconnects at the center of the plasma.

From Figs. 2.8-5 and 2.8-6, we see that the CT expands in the direction parallel to the tokamak toroidal field (ℓ increases) while being compressed in the perpendicular direction (d decreases) as expressed in Eqs. (2.8-20) through (2.8-22). The maximum contraction is limited by the condition that the pressure equilibrium is achieved (i.e., $\Delta P = 0$), so that the minimum contracted diameter of the CT at location s , $d(s)$, is given by:

$$\frac{d(s)}{d(0)} = \left\{ P_p + (P_p^2 + P_B^2)^{1/2} \right\}^{1/2}, \quad (2.8-30)$$

where

$$P_p = \frac{\ell(0)}{\ell(s)} \frac{\mu_0 P_{\text{ring}}(0)}{B_\phi^2(s) + 2\mu_0 P(s)} \quad \text{and}$$

$$P_B = \frac{B_{\text{ring}}^2(0)}{B_\phi^2(s) + 2\mu_0 P(s)}.$$

We notice that the injection velocity for the CT to achieve a full 90° rotation at the center is much larger (Fig. 2.8-6) than that for the CT to stop at the center when we relaxed the tilt condition (Fig. 2.8-5). Evidently, from Fig. 2.8-5, with an initial injection velocity of only 20 km/s, the CT flips in about 1.4 μ s after injection, while it would take

Figure 2.8-5. CT dynamics as it traverses the plasma assuming, for illustrative purposes only, that no field line reconnection occurs. The CT exits the coaxial gun at A and enters the edge of the plasma at B. Its initial velocity is selected so that its forward velocity is reduced to zero at the center of the plasma at C.

Figure 2.8-6. CT dynamics as it traverses the plasma for the real situation where field line reconnection occurs following a 90° rotation. The CT exits the coaxial gun at A and enters the plasma edge at B. Its initial velocity is selected so that its magnetic moment becomes fully aligned with the external field at the fuel deposition point C at the center of the plasma; notice that the CT still has appreciable forward momentum at this point.

6.3 μ s to slow down and stop at the center. Consequently, we have to inject the CT with a larger velocity so that the penetration time can be reduced to be about the same as the 90° tilt time. As we can see from Fig. 2.8-6, the injection velocity must be increased to about 780 km/s to meet this condition.

Once the CT has reconnected, we noted above that it may still contain sufficient energy density to distort inward the embedded, reconnected toroidal field. In our radial zoning scheme, we apply a simple model to approximate the residual forward motion of the plasmoid/field combination after reconnection by taking the kinetic energy lost in each zone and adding it to the energy stored in the compressed field $B_z(s)$ in the subsequent zone (see Eqs. (2.8-18) and (2.8-19)). In the case of the parameters shown in Fig. 2.8-6, we find that the forward momentum of the CT after reconnection at the center is sufficient to inwardly distort the field an additional 20 cm before being brought to rest, tearing and further reconnection notwithstanding. This phenomenon will be considered further in the next section.

We observe that the injection velocities calculated from the constraints of slowing down and tilting are much less than the local plasma Alfvén velocity, which is about 7700 km/s. Thus, the condition is satisfied that the CT injection velocity be less than the local plasma Alfvén velocity for avoiding possible disruption initiation. This criterion is also obeyed in our subsequent calculations with various CT mass and sizes.

In Fig. 2.8-7, we show the penetration distance as a function of the injection velocity for a CT carrying 5 mg of DT fuel with the CT diameter as a parameter. The penetration distance is computed with the two different criteria: i.e., the more probable situation that the tilt angle goes to zero (solid line) and, for illustrative purposes only, that the forward velocity vanishes with no tilting/reconnection (dashed line). As we noted above, the injection velocity requirement is dominated by the tilting requirement for the parameter range here. If we neglect tilting (dashed line in Fig. 2.8-7), then a small diameter CT can penetrate further than a larger diameter counterpart of the same mass and injection velocity. This is due to the fact that a smaller size CT has a smaller cross-sectional area projection perpendicular to the tokamak toroidal field and a smaller pressure gradient across it (Eqs. (2.8-25) and (2.8-26)). This results in a weaker retarding force. On the other hand, as we see from Fig. 2.8-7, the penetration distance and CT

Figure 2.8-7. Penetration distance to the fuel deposition point for a 5 mg CT as a function of the injection velocity; the initial diameter of the ring is shown as a parameter. The dashed lines show, for illustrative purposes only, the effect of requiring zero forward velocity at the deposition point assuming no field line reconnection. The solid lines show the real situation where 90° tilting/reconnection occurs at the deposition location.

diameter relation is reversed when the tilting condition is applied because the critical characteristic determining the tilting rate is the angular acceleration resulting from the interaction of the CT magnetic moment with the external field (Eq. (2.8-16)). For a given mass, a larger diameter CT has a larger moment of inertia and a smaller internal field (the latter arising from the lower density for a given CT beta limit). Accordingly, a larger CT of the same mass penetrates further when the tilting/reconnection condition is applied. There is thus a rationale to use the maximum allowable CT beta to reduce the CT internal magnetic field in the generation phase.

In Fig. 2.8-8, the variation of the injection velocity required for the CT to reach the center of the plasma is shown as a function of the CT mass. Only the tilting criterion is considered since it is the more restrictive requirement. The required injection velocity is shown for CT diameters of 5, 25, and 50 cm. The reduced velocity requirement for larger mass CT is simply the effect that the deceleration (angular acceleration) is inversely proportional to the mass for a given force (torque).

Figure 2.8-8 also serves as a guide for the selection of the velocity for central CT fueling for various choices of CT mass and diameters for the set of TIBER ETR baseline parameters in Table 2.8-1. Evidently, the velocity is in the range of 500 km/s to 1500 km/s. The injection velocity is lowered by utilizing a CT of large mass and large size. This is considered further in the next section.

Figure 2.8-9 shows the injected energy of the CT as a function of the mass and diameter, with velocities corresponding to those in Fig. 2.8-8. This energy requirement increases with mass since the reduction in injection velocity with increasing mass is at a rate less than $1/\sqrt{m}$.

Referring to the fueling current requirement of 156.7 A for the TIBER ETR design as indicated in Table 2.8-1, we can evaluate the average power required for CT injection to achieve central fueling under steady-state operation. This average line power is shown in Fig. 2.8-10 with an assumed efficiency of 50%. Clearly, the average power decreases for larger CTs injected at smaller repetition rates, which is in contrast to the energy for a single shot, as shown above. We note that the average power deposited in the plasma under

Figure 2.8-8. Injection velocity requirements as a function of CT mass for central penetration with tilting/reconnection at the fuel deposition point; the initial CT diameter is shown as a parameter.

Figure 2.8-9. Injected energy per shot as a function of CT mass for central penetration with tilting/reconnection at the fuel deposition point; the initial CT diameter is shown as a parameter.

Figure 2.8-10. Average line power to achieve steady-state central fueling as a function of the injected CT mass. The initial CT diameter is shown as a parameter.

steady-state fueling conditions is just over 1 MW, a value which is small relative to both the internal alpha power (63 MW) and externally-supplied current-drive power (62 MW).

2.8.4.3 Conclusions: The Advantages, Uncertainties, and Applications of CT Fueling

We have seen above that, in principle, accelerated compact toroids provide the only near-term method of achieving deep-penetration-fueling in reactor grade plasmas, a fact which is given additional credence by the existence of a currently operating accelerated CT experiment with commendable performance parameters [22]. In summary here, we assess the strengths and weaknesses of the method and address some of the uncertainties in our current modeling.

The advantages over other advanced fueling schemes can be summarized as follows:

1. It would appear to be the only near-term scheme which would have direct applicability to central fueling of reactor-grade plasmas such as the prospective Engineering Test Reactor.
2. In principle, the fuel deposition point can be precisely determined via control of the injection velocity by simple adjustment of the $\underline{J} \times \underline{B}$ accelerating force in the coaxial accelerator. This is in distinct contrast to conventional pellet fueling in the outer regions of the plasma where fuel deposition is smeared over the ablation path of the pellet and where further inward-pinching of the fuel density must be invoked for deeper penetration.
3. Because, to a first order, the dynamics of the CT following its exit from the coaxial gun are determined by the magnetic field geometry of the tokamak rather than the plasma conditions, initial verification tests of CT fueling could be performed in a simple experimental assembly without the need for a full reactor-grade plasma. Again, this is in contrast to conventional pellets where pellet ablation dynamics are a direct consequence of the high plasma temperature and density.
4. Unlike the mechanical nature of conventional pellet production and acceleration, the formation and acceleration of compact-toroids are performed under electrical control. This makes precise fueling control much easier (see above regarding control of deposition point

via velocity) and facilitates regulation of fueling parameters under high repetition rates if required.

5. Also, unlike conventional pellet manufacturing and acceleration techniques, there is no tritium debris associated with the compact toroid as it emerges from the gun. Excess tritium utilized in the CT formation process can be electronically gated out of the coaxial accelerator and left behind to be recycled in the source.
6. Compact toroids manufactured and accelerated in a system of given dimensions and magnetic field can be compressed to small radii before injection into the tokamak by the simple addition of a conical converging section at the end of the coaxial gun. This alleviates potential problems of restricted port size and access to the tokamak chamber. Note, however, that the increased stored energy in the compressed magnetic field of the CT is supplied by a corresponding decrease in its forward kinetic energy. Hence, in this case, we must compensate the original velocity so that we retain the desired injection velocity into the plasma.
7. The CT carries magnetic helicity as well as mass and, in principle, provides a high-Q method of non-inductive drive for the plasma current. This is considered in a separate paper [25]. The current perturbations due to the CT must not be so large as to stimulate a major disruption. This criterion should be met since the CT helicity is typically much less than the tokamak helicity inventory.
8. We mention here in passing that CT fueling could be a valuable contributor to alpha-particle diagnostics. A currently proposed [26] diagnostic technique in prospective DT machines is to seed the plasma with boron or lithium impurities. Subsequent measurement of emerging neutrons or gamma rays from (α, n) or (α, γ) nuclear reactions with these impurities yield information on alpha particle spectra in the machine. Clearly, a fueling technique such as the CT, which could provide localized deposition of these target materials at any desired point, would be of great value in precise determination of the alpha-particle distribution function.

The largest uncertainty in our modeling of the plasma interaction concerns the behavior of the CT at the reconnection point. We have shown that the injection velocity condition is determined by the requirement that the CT has

tilted through 90° and undergoes reconnection at the desired fuel deposition point. The consequence is that the reconnected plasmoid still retains an appreciable forward momentum which, in our simple model above, is sufficient to distort the reconnection field a further 20 cm inwards. Accordingly, a deeper insight is required of CT dynamics at this point. For example, does resistive tearing and further reconnection occur? Do the outer poloidal field lines of the CT connect with the tokamak field leaving the inner "core" to continue further inward? Are reconnection rates constrained by plasma conductivity? If significant tearing occurs, is there a possibility of disruption initiation or will it behave like mild sawteething? Is it possible that Rayleigh-Taylor short-wavelength modes, driven by the rapid deceleration of the CT after reconnection, will result in facilitated tearing and dissipation? Does the bulk of the fuel remain tied to the fieldlines where initial reconnection occurred or will it be deposited at radii inside the 90° reconnection flux surface? Clearly, formal assessment of these questions will require application of 3-dimensional MHD theory and a definitive experiment.

The other uncertainty in this technique concerns the ability to make clean, low- $\bar{z}$ compact toroids in the Marshall gun source. Marshall guns have the reputation of being rather "dirty" in that ablated material from insulators and wall structures can accumulate in the plasma discharge. A set of experiments are currently being proposed for the RACE experiment at LLNL to address this question. Could, for example, the high mass debris be electronically gated from the main DT plasma ring by virtue of the lower velocity of the former? The prospects for this look encouraging.

Finally, we return to the fueling requirements for the TIBER Engineering Test Reactor. In Table 2.8-3, we present a suggested list of typical specifications based on the results of Section 2.8-4 above. As we saw, a larger mass per fueling shot resulted in a larger injected energy per shot but a smaller repetition rate and lower average power requirements (Figs. 2.8-9 and 2.8-10). In terms of dimensions, we saw it was advantageous to inject a larger diameter CT, providing access-constraints on port size permit this, and providing that the expanding CT can be accommodated as it traverses the plasma. We noted above that it is a simple matter to compress a compact toroid after acceleration and before injection by means of a conically converging coaxial section at the expense of a slight loss in velocity. By this means, restrictions on port size and machine access can be relatively

Table 2.8-3. Suggested parameters for CT injection fueling of the engineering test reactor.

Injection velocity (km/sec)	750
DT mass/shot (mg)	4.8
Rep-rate (Hz)	0.93
Energy/shot (MJ)	1.35
Average wall-plug power (MW)	~ 2.5
Initial CT diameter at entrance port (m)	0.5
Capital cost (M\$)	~ 15 ^a

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^a The total direct capital cost of the TIBER ETR is estimated at ~ 1.2 B\$ without contingency.

easily accommodated. Clearly, the mass per shot, repetition rate, and diameter could be chosen to optimize the overall cost of the system in terms of injection velocity, energy, and wall-plug power requirements. We have not attempted this in Table 2.8-3, but rather have simply listed a consistent representative set of parameters.

Note in Table 2.8-3 that we have selected a low repetition rate of around one per second and a corresponding high energy per shot (~ 1.4 MJ). Assuming an overall efficiency of 50%, our capacitor bank storage requirement of ~ 2.7 MJ will cost ~ 9.5 M\$ at a unit price of 3.5 \$/J [27]. We estimate the total installed cost for a tritium-tight fueling system as approximately 15 M\$, the bulk being in the energy storage. Clearly, we can reduce the energy storage requirement by injecting smaller mass toroids at higher repetition rates. However, the difficulty of rapid-recycle operation of large capacitor banks would make this an untenable scenario at repetition rates of a few Hz or more. In this context, it is interesting to note that a new concept is currently under development for generating pulsed power with high energy per shot and high repetition rates [28]. This is based on a compensated pulsed alternator ("Compulsator") which embodies a novel operating principle of compensated inductance and overcomes the inherent limitation of AC machines as pulse generators. As an example, a Compulsator under test at the University of Texas [29] has performance specifications of 1 MJ per pulse at a repetition rate of 240 Hz; the pulse width is ~ 1 to 2 ms, but shorter pulse lengths should be easily achievable. With this type of device we would have

the potential for rapid energy transfer, as in conventional capacitor banks, but with the significant advantage of repetition rates greater than 10 Hz. Cost savings should also be appreciable.

2.8.5 SUMMARY AND CONCLUSIONS

An assessment is studied of the fueling requirements for sustaining a constant fuel level for both a steady-state and long-pulsed modes of operation for the TIBER Engineering Test Reactor. We emphasize the desirability of depositing fuel in a controlled manner at any point in the plasma because of the advantages of central fueling and density profile control.

For the pellet fueling scheme, we extend the neutral shielding pellet ablation model to include the effect of the fusion-born alpha particles, which only accounts for a mild 5% enhancement in the ablation rate. The high electron temperature and density of a fusion-grade plasma makes pellet penetration very difficult. We find injection velocities of about 50 km/s and 30 km/s are required for 1/2 plasma penetration for the steady-state and pulsed-inductive modes, respectively. The pellet launch length has to be large (many tens of meters) in order to avoid pellet breakage due to the low inertial strength for DT pellets. We find these two restrictive requirements impose a severe limitation on the application of the pellet injection scheme for adequate fueling of the TIBER design.

We also examine the application of a novel but practical technique for central fueling of reactor-grade tokamak plasmas by compact torus (CT) driven by a coaxial accelerator. We employ a self-consistent, radial-zoning scheme to model the transport of the CT; this scheme requires quantification of several interrelated constraints, including ring decay, tilting, field-line reconnection, deceleration in the external field gradient, and ring expansion/contraction. Subsequent application of this formalism to the TIBER design enables the parameterization of fueling requirements in terms of injection velocity, fuel mass, penetration distance, CT dimensions, and repetition rate. Assessment is also given to hardware implications for near-term applications.

A brief survey is done on the various operational and conceptual pellet injection schemes and the CT injection scheme for plasma fueling. We conclude that the CT injection scheme is the only advanced method capable of controlled and localized deep-penetration fueling which is developable in the immediate

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near-term and for which an experimental system is currently operating. As a result, the TIBER fueling system is baselined on the application of the CT injection system as indicated in Table 2.8-3.

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2.9 ALPHA PARTICLE EFFECTS

2.9.1 INTRODUCTION

The TIBER reactor will require long-pulse or steady-state operation. While existing devices such as TFTR and JET will address some single-particle alpha effects and CIT will examine some collective behavior, TIBER will be the first device where the full range of collective behavior and equilibrium effects will be faced. Indeed, a key physics objective of TIBER is to provide experience with such effects in a situation similar to that expected for later power reactors. Consequently, a thorough study of the influence of alpha effects on the design of TIBER has been initiated. The issues involved include:

Alpha slowing and heating effects

- Self-consistent thermalization and heating profiles.
- Losses, especially due to ripple trapping, including influence on first wall and coil designs.
- Instabilities--effects and control.

Alpha "ash" transport and control

- Ash transport and possible accumulation.
- Effect on operation window and power density.
- Development of control methods if required.
- Relation.

Burn control

- Techniques for both driven and ignited operation.
- Controlled and emergency shut-down techniques.
- Interrelation with other subsystems.
- Optimal control methods for testing operation.

Since this work is only beginning, the effort to date has not covered these items in equal depth. Some preliminary results on instability losses due to ripple trapping, ash effects on the operating window, and control requirements for driven operation are presented in following sections.

2.9.2 INSTABILITIES

Whether or not any alpha-driven instabilities will occur and significantly affect the TIBER design and operation remains an open question. Various possible effects were discussed at a workshop held in June 1987 in Germantown,

MD [1]. Comments from that review are very briefly outlined here. During the next study period a more specific study of candidate effects relative to TIBER is planned.

2.9.2.1 Comments from Workshop Review

Alpha fishbone oscillations could arise from an internal, ideal $m = 1$ interchange mode centered at $q = 1$, which is FLR stabilized in the absence of dissipative terms [2]. However, the trapped alpha precessional resonance provides a destabilizing effective viscosity, allowing alphas to slow to about 400 keV before they are resonantly lost. This theory contrasts with the Princeton fishbone theory which predicts fast alpha losses near 3.5 MeV. Both mechanisms may occur.

Estimates [3] show that the global Alfvén eigenmode, the shear Alfvén gap mode, and the ballooning mode simply are all destabilized by alphas. Scaling of these modes favors a small tokamak with a large aspect ratio.

Studies [4,5] indicate that the loss-cone type of instability is below threshold even when all the trapped alphas are ripple lost from a plasma with no magnetic field gradient. Previously, the magnetosonic wave had been found essentially stable. Also, the shear Alfvén wave is difficult to destabilize by alpha-particle velocity space effects. Thus, velocity-space inversion conditions appear unlikely in TIBER-like ETRs, although possible inversion associated with sawtooth action still needs study.

Numerical results [5] indicate that moderate-to-high- n ballooning modes are destabilized by the trapped-alpha precessional resonance in ETR. A tentative conclusion is that plasma stability improves dramatically as elongation increases from 2 in CIT (with a critical beta of 1%) to 2.4 in ETR (with a critical beta of 3%). Further checks are necessary, however, to confirm this result. The overall trend, however, is toward lower critical beta due to the alpha precessional resonance. This mode is the high- n version of the fishbone instability.

Overall conclusions from the workshop discussion can be summarized as:

- Alpha-driven fishbones appear to be the most threatening instability for fast alpha losses.
- Nonlinear evolution of alpha-driven Alfvén modes needs careful study.
- The various alpha-driven modes may appear in near-term devices like CIT, providing important data for TIBER.

- Increased plasma elongation may stabilize some of the alpha-driven MHD modes in regions of nonzero shear.

These problems clearly need much further investigation because of the high uncertainty in their evaluation and their potential importance to reactor operation.

Some additional effects that were not covered in the workshop discussion will also be considered. One important problem is the electric field associated with alpha transport and losses. Preliminary estimates [5] indicate a significant plasma rotation (due to $E \times B$ forces) can result. While this does not appear dangerous relative to plasma breakup, the effect on background plasma and impurity transport needs evaluation.

2.9.3 TOROIDAL FIELD RIPPLE-LOSS OF FAST ALPHAS

2.9.3.1 Introduction.

Fast alpha particles can be trapped in the ripple of magnetic field between toroidal field (TF) coils and drift out of the plasma. This loss of fast alphas can potentially have two undesirable effects: reduced alpha heating and localized alpha bombardment of the first wall. For the TIBER design, the loss will probably be sufficiently low that the reduction in heating will not cause a significant change in the plasma energy balance. However, the local wall bombardment remains a potential concern. The extent of this problem remains uncertain at this point, though, because of major discrepancies in calculational models that have not been resolved. This problem is discussed in some detail here, and a resolution of the discrepancies is planned during next year's study period.

The outline of this section is as follows: Section 2.9.3.2 briefly reviews existing work on the loss of fast ions due to toroidal field ripple, then discusses the present work in more detail. Section 2.9.3.3 enumerates various modeling issues and ways to improve the models. Section 2.9.3.4 points out several new aspects of the TF ripple-loss problem which can have an important impact on the results. Consequences for TIBER are explained in Section 2.9.3.5.

2.9.3.2 Existing Work.

The effect of TF ripple in a tokamak was first identified as an important problem by Anderson and Furth [6]. Subsequent research has been extensive (see Ref. [7]). This section summarizes the basic physical processes and recent study based on these effects.

2.9.3.2.1 Basic Ripple Processes

With discrete toroidal field coils, the magnetic field strength is lower between the coils than in the plane of the coils. Secondary magnetic wells form when the magnitude of the ripple, $\sigma(R, Z)$, is sufficiently large [3]:

$$a^* = RB(\bar{B}_p \cdot \nabla B)/N\delta B_t^3 < 1, \quad (2.9-1)$$

where R is the major plasma radius, $\bar{B}_p$ is the poloidal component of the magnetic field vector, B_t is the toroidal component of the field, B is the magnitude of the net field, and N is the number of toroidal field coils. This generalizes the formula for a circular tokamak [6] to a noncircular equilibrium where all magnetic field terms in Eq. (2.9-1) are for unrippled values. Thus, for example, for the Fusion Engineering Device [9], Fig. 2.9-1 shows an example of contours of constant ripple and the outer flux surface of the equilibrium; ripple wells form between the mid-plane ($Z = 0$) and the contour of $a^* = 1$. Ions moving in this rippled field are affected several ways. They can be collisionally ripple-trapped in a well by pitch-angle scattering near the banana tip of a barely unripple-trapped orbit as they traverse that ripple well. They can also be captured in a ripple well (near the banana tip) as the large-banana-width orbit collisionlessly carries the particle into a region of higher ripple. The inverse of these two processes, detrapping [10], also can occur by scattering at the banana tip of its orbit or as the large-banana-width orbit carries a fast ion out of a ripple well in less than one bounce period. Ripple-trapped particles no longer follow banana orbits, but rather oscillate toroidally between two maxima in the B -field while moving vertically due to curvature and gradient- B drift. Such ions are either lost to the wall or move into a region where ripple wells do not exist. The former case is a direct wall-loss mechanism. In the latter case, the particles move radially outward where diffusive processes occur. Initial work on this problem is reported in Refs. [11-14]. A more recent analysis by

Figure 2.9-1 For FED, an example of contours (solid line) of constant ripple (%), the edge of the noncircular plasma (dashed line), and the curve (chain dashed line) of $\alpha^* = 1$.

Goloborod'ko et al. [15] uses a kinetic treatment for collisionless losses. However, Yushmanov [16] points out that the drift kinetic equation used by Ref. [15] is inapplicable for fast alphas where the toroidal drift precession distance during one bounce period is greater than the distance between TF coils.

Since collisionless effects are much faster than collisional processes for MeV alphas, some early work [6,11-13] assumed that ripple-trapped ions are instantaneously lost to the wall. However, using a model based on Ref. [17], Hively [18] found that such losses are dramatically reduced by collisionless ripple detrapping in an INTOR-type device. Other calculations based on Ref. [19] and independent studies by Ref. [20] agree with this conclusion.

Without ripple, banana-trapped particles precess toroidally by $\Delta\phi$ during one bounce period. With ripple, neoclassical orbits do not close exactly due to a variable lingering period as the banana tip crosses a ripple well or ripple inflection region leading to a radial displacement. For high energy alphas (> 1 MeV), the toroidal precession is large compared to the period of the toroidal field coils, $\Delta\phi > 2\pi/N$. Stochastic diffusion of these fast alphas occurs when the radial displacement, d , is sufficiently big compared to the minor radius, (a) , of the plasma, $d > a/N$. Stochastic ripple diffusion leads to rapid loss of ions to the wall in 10 to 100 bounce periods. The collisionless form of this process is reported in Refs. [20-22]. Yushmanov [16] has studied the collisional stochastic diffusion problem. A generalization of stochastic diffusion occurs for lower energy alphas. In this case, the toroidal precession distance during one bounce is $\Delta\phi > 1/N$, and most particle orbits undergo a systematic outward displacement, d , since the phase of each successive displacement changes by $N\Delta\phi > 1$. However, banana-trapped particle orbits undergo large radial jumps when $\Delta\phi = n\pi/N$, due to resonance between the toroidal drift precession and the ripple periodicity. Yushmanov [23] refers to the resulting transport as resonant diffusion, and he describes a smooth transition among the banana-drift, ripple-plateau, and resonant diffusion regimes, assuming small banana-width orbits. Goloborod'ko et al. [24] use a more general theory for large banana-width orbits and find resonances when the ratio of the toroidal drift-precessional frequencies to the bounce-frequency is a rational number (m/n) for m and n integers (localized in both velocity and coordinate space). Calculations by Tani et al. [25,26] find that most of the fast, trapped alphas are lost from an INTOR

plasma via stochastic ripple diffusion. However, Hively points out [7] that the validity of this calculation is questionable because the mapping procedure used by Tani et al. is not area preserving. Recently, Zajtsev et al. [27] have used a flux-surfaced-averaged Fokker-Planck equation to model fast alpha confinement based on a small-banana-width form for the stochastic diffusion; their results agree with Tani. However, the validity of their model is questionable because a small banana-width model is employed for stochastic diffusion.

Banana-plateau and banana-drift diffusion are related processes in which neoclassical orbits fail to close exactly due to the same ripple-induced, variable lingering period as the banana tip crosses a ripple well or a ripple inflection region. Collisions decorrelate the ripple phase causing diffusion, rather than oscillatory motion. If the collisionality is low (collisionless regime), collisions decorrelate the ripple phase after many bounce periods, leading to banana-drift diffusion [28]. At moderate collisionality, collisions decorrelate the ripple phase between successive bounces, producing banana plateau diffusion [29]. When the collisionality is high, collisions decorrelate the phase after a small fraction of a bounce period, but this case does not apply to fast alphas. Goldston and Towner [14] discuss a generalization of these processes for large banana-width alpha orbits during slowing down. Mynick [30] has constructed a generalized theory of banana-drift transport which also includes low- n NHD modes and high- n perturbations arising from microtearing. Very recently, Zweben et al. [19] have calculated alpha losses based on a guiding-center orbit simulation [31]. They find significant alpha losses in TFTR due to stochastic ripple diffusion, due to sawteeth (modeled as a stationary $m/n = 1/1$ island in the plasma core) and due to higher order resistive tearing modes (also modeled as stationary islands).

While the above diffusive processes are due to abrupt changes in the toroidal adiabatic invariant as an alpha orbit moves through magnetic ripple, diffusion also occurs as a result of jumps in the magnetic moment [25]. Putvinskii and Shurygin [26] point out that Ref. 32 did not take account of the local character of the gyroresonance interaction; they then go on to obtain a corresponding diffusion and critical value of ripple. There is no analytical theory which accounts for simultaneous changes in the toroidal adiabatic invariant and the magnetic moment, although Monte Carlo, guiding-center simulations (e.g., Refs. [18,19,20]) include both jumps.

2.9.3.2.2 Simulations of Combined Processes

In his 1980 review, Kolesnichenko [34] discussed only the theory of Belikov et al. [12] corresponding to collisionless ripple trapping without detrapping. Improved ripple modeling was done by Tani [26,35,36,37] for the slowing down and scattering of neutral injected ions and alphas, including collisionless ripple trapping/detrapping, banana drift, and ripple diffusion losses. Fowler and Rome [17] developed a more complete model for neutral beam injection.

Recent simulations by Tani et al. [26] and Hively [18] have studied alpha particle ripple losses for INTOR in detail. Table 2.9-1 compares the physics models of Refs. 18 and 26. Reference 26 finds moderate alpha losses (10 to 20%) for a circular INTOR plasma with an edge ripple of 0.75%. Reference 18 finds such losses to be small (2.5 to 3%) for both circular and noncircular INTOR plasmas. Table 2.9-2 compares key results of the two studies in more detail. The remainder of this subsection elucidates this comparison in detail.

Regarding the distribution of alpha flux at the wall (see Fig. 2.9-2), there are several important differences between Refs. 18 and 26. The peak fluxes differ by a factor of three: Tani's maximum flux is 0.6 MW/m^2 versus Hively's value of 0.2 MW/m^2 . The toroidal distribution found by Tani is high and broad for poloidal angles less than 40° and is abruptly lower (but still broad) for larger poloidal angles. Hively's results are strongly peaked between the TF coils as well as strongly peaked at low poloidal angle, i.e., at the outboard midplane. (The sign reversal in poloidal angles between Fig. 2.9-2a and Fig. 2.9-2b is because the toroidal fields are reversed for the two cases. The direction of the toroidal current was also opposite, reversing the labeling of the toroidal axis.) Figure 2.9-2a shows large statistical variations in Tani's flux due in part to his small bin size; no such fluctuations are evident in Fig. 2.9-2b.

Figure 2.9-3 compares the energy loss spectrum obtained by Tani et al. [26] and Hively [18]. The overall features are similar, including large losses near the birth energy and abruptly lower losses during thermalization. However, a chi-squared statistical comparison of the two spectra shows significant differences. Tani's spectrum decreases monotonically from the birth energy to 2.6 MeV, is flat from 2.6 to 1.6 MeV, then rises slowly from 1.6 to 0.2 MeV, and finally decreases somewhat in the lowest (thermal) bin.

Table 2.9-1. Comparison of physics models.

Model	Tani et al. [26]	Hively [18]
Equilibrium ^a	Circular, low-beta	Noncircular, high-beta with vertical offset and poloidal null
Ripple profile ^a	Sheet current model scaled by ripple magnitude at edge	Finite cross-section, D-shaped TF coil
Plasma profiles ^a	Parabolic in r/a	Transport consistent
Wall ^b	At plasma edge	D-shaped, displaced from plasma edge
Collisions ^c	Slowing down Pitch angle scatter	Slowing down Pitch angle scatter Speed diffusion Toroidal electric field

^aThese aspects of the model affect the realism but not the basic physics.

^bLosses decrease rapidly with plasma-wall separation [40], so this aspect of the model affects both the realism and the physics.

^cCollisional effects due to speed diffusion and toroidal electric field are unimportant [18].

Hively's spectrum has more losses near 3.5 MeV and has large fluctuations in the loss during thermalization. In particular, no losses occur near 0.5, 1.9, 2.7, and 3.1 MeV. This observation is thought to correspond to energies where the toroidal drift precession during one bounce period is a multiple of the TF coil separation distance. Then the orbit is locked in resonance with the TF ripple and is well-confined (not stochastically diffusing). Thus, as fast alphas thermalize, they alternately undergo ripple loss, followed by being stably confined. Such stable confinement "islands" have been seen previously for the fast alphas in a tokamak with ripple due to TF coils [38] and due to a bundle divertor [39]. The loss spectrum of Tani et al. [26] is shown in Fig. 2.9-2a for an edge ripple of 1.5%, versus Hively's result [18] in Fig. 2.9-2b for an edge ripple of 0.75%. Since the ripple magnitude does not

Figure 2.9-2. Wall flux of lost alphas from a circular INTOR plasma with an edge ripple of 0.75% as found by (a) Tani et al [26] and (b) Hively [18].

Table 2.9-2. Comparison of results for circular plasmas.

Issue	Claims by:		
	Tani [26]	Hively [18]	Comparison
Importance of collisionless detrapping	Small	Moderate	
Importance of stochastic diffusion	Large	Moderate	
Toroidal wall loading distribution	Broad	Peaked	Fig. 2.9-2
Poloidal wall loading distribution	Broad	Peaked	Fig. 2.9-2
Fluctuation in wall loading distribution	Large	Small	Fig. 2.9-2
Energy/particle loss fraction (%)	10/20	2/3	
Fluctuation in energy loss spectrum	Moderate	Large	Fig. 2.9-3
Reason for this spectral fluctuation	Statistics	Reasonance	Fig. 2.9-3

change the resonance energies, the key feature of this comparison is the presence of low loss regions. No such stable confinement islands are seen in Tani's spectrum.

2.9.3.2.3 Ripple Loss Experiments

Tokama experiments have found fast ion losses [41] and thermal diffusion [42] in reasonable accord with theory. A more detailed comparison by Scott [43] found the loss (ripple trapping and stochastic diffusion) of neutral-beam-injected ions as expected theoretically (see Ref. 17). However, very well diagnosed fusion-product-confinement experiments are needed to resolve the discrepancies in modeling of these high energy particles notes here.

2.9.3.3 Modeling Issues in Existing Simulations

Issues which impact the realism and/or accuracy of the alpha trapping model include:

- Taking the first wall at its real position vs the plasma edge. The resulting loss fraction is a strong function of the plasma-wall separation [44].
- Using realistic profiles for the temperatures and densities rather than theoretical transport-consistent values. The consequent alpha

Figure 2.9-3. Alpha energy loss spectrum in a circular INTOR plasma as found by (a) Tanai et al. [26] and (b) Hively [18]. Figure 2.9.3.2(a) was adapted from Ref. [35] by converting the spectrum to a linear scale and reappportioning the data to the same energy bins as used by Ref. [18] to allow a chi-squared comparison (see text).

particle source function, as well as the influence of the plasma on slowing down and scattering, depends on these profiles.

- Determining the source function from the ion densities and temperatures rather than assuming an ad hoc alpha particle source profile. Losses are strongly dependent on the actual alpha source distribution.
- Integrating the alpha particle guiding center motion with sufficient accuracy that numerical diffusion is not introduced. Particle following in field-line coordinates is presently the most computationally efficient technique [31], but inversion to real space must be done carefully.
- Using sufficient particles in a Monte Carlo simulation to obtain adequate loss statistics. Source biasing improves these statistics; a statistical error analysis is needed.
- Using realistic ripple fields for TF coils with a finite cross section is essential. Simple analytical fields are easy to model but are not realistic for modern D-shaped coil shapes. Also, the ripple profile changes as the edge ripple varies, due to the shift in the TF configuration [18].

There are a number of important physics issues that must be included in future calculations; for example:

- An adequate collision model for alpha thermalization and pitch angle scattering is essential. Collisional effects are small due to speed diffusion and toroidal electric field acceleration [18]. Pitch angle scattering without slowing down is unphysical [19] because the physics of stably confined, locked orbits is omitted.
- The rippled tokamak-equilibrium should be divergence free ($\nabla \cdot \bar{B} = 0$). A nonzero divergence causes unphysical orbital motion since the time rate of change of the magnetic moment is proportional to the divergence of $\bar{B}$ to lowest order [43-45]. The rippled part of the magnetic field must also be curl free if a superposition of axisymmetric and rippled parts is used. The nonzero curl of the rippled field is equal to a corresponding ripple current density ($\bar{J}_{rip}$) which will modify the particle motion due to a $\bar{J}_{rip} \times \bar{B}$ force. Various models, e.g. Ref. [36], violate these criteria. When a superposition of axisymmetric and ripple parts is used, the

axisymmetric toroidal component must be subtracted to avoid double counting of the toroidal field. A superposition is, in fact, erroneous since the rippled equilibrium is fully 3-dimensional. Future calculations should include such 3-D equilibria [46,47] to represent "self-healing" properties of the plasma to ripple perturbations. Moreover, the alpha particle pressure will modify the 3-D tokamak equilibrium and should be included.

- Modeling the alpha particle interaction with the ripple field requires care. For example, a time-scale enhancement for simulating collisional processes for trapped orbits relative to collisionless ones produces large errors.

2.9.3.4 New Physics to Be Included in Future Models

There are several additional physics aspects that need to be included in the next step calculations. These effects include the large gyroradius of the fast alphas, electric fields in the plasma (particularly in the radial direction), and sawteeth. Some comments about each follow.

2.9.3.4.1 Finite Gyroradius Effects

The fast alpha gyroradius in a reactor grade plasma can have a significant effect on confinement, thermalization, and losses. These effects include:

- Distributing the alpha particle source distribution over the gyroradius. The source distribution will then be less centrally peaked, leading to more losses and less central alpha heating.
- Changing the topology of the collisionless alpha orbits due to second-order corrections in the canonical toroidal angular momentum.
- Distributing the alpha particle orbit interaction with the ripple over the gyroradius. This effect will be particularly important at the banana tip of a trapped orbit as it traverses a ripple well.
- *Distributing the wall loss position of the fast alphas over the gyroradius.* Localized wall heating is then less peaked.

2.9.3.4.2 Electric Fields Effects

Electric field effects caused by the alphas will change the dynamics of both the background plasma and alphas as follows:

- Large-banana-width orbits will separate the doubly charged alpha orbits from their birth flux surfaces, creating significant poloidal and radial components of electric field.
- A radial electric field modifies the collisionless trapping process by adding additional drifts and by changing the component of alpha velocity parallel to the net magnetic field. Similar effects will occur for the other processes involved, especially at the banana tip of a trapped orbit.
- The background plasma dynamics in the presence of such an electric field should be considered (e.g., plasma rotation), particularly profile changes that modify alpha particle dynamics.
- These effects need to be considered self-consistently with the alpha distribution function, which is modified by the resulting electric field.

2.9.3.4.3 Sawteeth

Since sawteeth have been observed to expel high-energy fusion products from the plasma [50], sawtooth control or elimination may become crucial for TIBER. Modeling of alpha dynamics in the presence of sawteeth has not yet been attempted, and some aspects can be identified here:

- Sawtooth-flattened background plasma profiles will affect the alpha confinement and thermalization.
- A high-beta plasma has the plasma center shifted outward, causing the alpha orbits to sample a larger ripple field and leading to larger losses.
- Aside from ripple losses, the degree of sawtoothing will have a strong effect on thermal alphas (ash) accumulation in the plasma core. This problem is discussed further in Section 2.9.4.

2.9.3.5 Consequences for TIBER

Calculations by Hively [18] find a 2.5 to 3% energy loss of fast alphas from INTOR due to TF ripple. The corresponding reduction in plasma heating by the alphas is insignificant. The flux of alphas lost to the wall is peaked poloidally at the outboard wall, between the TF coils toroidally, and above (or below) the plasma midplane. The local alpha heat flux is double that without ripple losses but can be accommodated by appropriate wall design. The

structural wall damage is probably unimportant, but the synergism between alpha sputtering and plasma sputtering could enhance the impurity generation, limiting the burn time [44]. These results need to be verified, but if true, they suggest that TF ripple losses are manageable. This still is a key design issue, however. The point is that to minimize costs the maximum allowable ripple is desired. Thus the present uncertainties in the loss rates and their effects must be reduced by improved modeling, experiments, and optimization studies.

Calculations by Tani et al. [26] find 10 to 20% energy loss of the fast alphas from INTOR; Refs. 25 and 27 agree with their result. Such high losses would have a large impact on the first-wall design, which would need to accommodate about 1MW/m^2 thermal loading along with the associated sputtering and blistering. Ignition would be more difficult since most of the trapped, fast alphas are lost; compensation via a larger machine-size would be expensive.

2.9.4 ALPHA ASH EFFECT ON PLASMA PERFORMANCE

2.9.4.1 Introduction

It is well known that the buildup of alpha-particle impurities in a plasma enhances the radiative power loss and the fuel ion depletion. Hence, in a steady-state operating reactor like TIBER II, the ash accumulation could lead to a serious deterioration of plasma performance. In this study we use a steady-state, global power balance equation, which is constrained by the density limit, the beta limit, and the energy multiplication requirements to investigate how the plasma operation regimes for TIBER II change as a result of the presence of the alpha ash.

2.9.4.2 Global Power Balance and Constraints

To find the regimes of plasma operation, we use a global, steady-state power balance equation,

$$F = -P_{\text{con}} - P_{\text{rad}} + P_{\alpha} + P_{\text{OH}} + P_{\text{aux}} = 0, \quad (2.9.-2)$$

where the P terms represent the power loss due to thermal transport processes (P_{con}), the power loss due to bremsstrahlung radiation (P_{rad}), the alpha

heating power (P_α), the ohmic heating power (P_{OH}), and the auxiliary power (P_{aux}).

The maximum electron density ($n_{e_{max}}$) of the operating regimes is given by either the Murakami limit [51,52] as

$$n_{e_{max}} = \langle n_{mu} \rangle = \nu_{mu} B_0 / (R_0 q_*) \quad (m^{-3}) \quad (2.9-3)$$

with $\nu_{mu} \approx 1.5$ or the Greenwald density limit $\langle n_{GR} \rangle$ [53]. The maximum volume-averaged beta is constrained by the Troyon limit [52] as

$$\beta_{tot} < \beta_{crit} = 4 \times 10^{-2} I / (a B_0) , \quad (2.9-4)$$

where I is the current in MA, B_0 is the magnetic field in tesla, a is the minor radius (in m), R is the major radius (in m), and q_* is the cylindrical equivalent safety factor [52]. The total plasma β is taken as the sum of the electron β (β_e), the fuel ion β (β_{DT}), the thermal alpha β ($\beta_{t\alpha}$), and the fast alpha β ($\beta_{f\alpha}$):

$$\beta_{tot} = \beta_e + \beta_{DT} + \beta_{t\alpha} + \beta_{f\alpha} = (1 + \gamma_{f\alpha}) \beta_{th} , \quad (2.9-5)$$

with $\beta_{th} = \beta_e + \beta_{DT} + \beta_{t\alpha}$ and $\gamma_{f\alpha} \approx 15\%$ for simplicity.

By assuming that profiles of the local density, temperature and current density are in the form $x = x_0 (1 - r^2/a^2)^{a_x}$ for $x = n, T$, or J with $a_J = 3a_T/2$ and that the electron temperature $\approx$ the ion temperature T (in KeV), we obtain an analytic expression of Eq. (2.9-2) in terms of the average density and temperature as [52,54]

$$-a_{con} m T_{10} (1 + f_n) / \tau_E - a_{rad} m^2 T_{10}^{0.5} + a_\alpha m^2 T_{10}^S + a_{OH} T_{10}^{-1.5} + a_{aux} = 0 , \quad (2.9-6)$$

where f_n is ratio of the ion density to the electron density (n_i/n_e), $s = 2$ or 3 for $T \geq 7.5$ KeV or ≤ 7.5 KeV, τ_E is the energy confinement time (in seconds),

$$T_{10} = T/10,$$

$$m = n e R_0 q_* / (1.5 \times 10^{-20} B_0),$$

$$a_{\text{con}} = 0.54 [Rq^*/(1.5B_o)] g_{\text{con}},$$

$$a_{\text{rad}} = 1.5^2 \times 1.68 \times 10^{-2} Z_{\text{eff}} g_{\text{rad}}$$

$$a_{\alpha} = 0.349 f_{DT}^2 g_{\alpha},$$

$$a_{OH} = 2.1 \times 10^{-3} (q_*^2/q_o^2) Z_{\text{eff}} (1 + \kappa_*^2/2\kappa)^2 g_{OH},$$

$$f_{DT} = n_{DT}/n_e,$$

n_{DT} = the fuel ion density, and

κ_* is the effective plasma elongation.

The g terms are factors resulting from averaging over the plasma volume and are functions of α_n and α_T . Assuming that the alpha ash is the only impurity in the plasma, then

$$Z_{\text{eff}} = \frac{n_{DT} + 4n_{\alpha}}{n_e}, \quad (2.9-7)$$

with n_{α} = the average density of the ash.

The main insight here is that the ash effects come into the calculation through the factors Z_{eff} , f_{DT} and f_n in Eq. (2.9.-6) and through the beta limit in Eq. (2.9-3). For further details about the above equations one can refer N. Uckan's earlier papers [52,54].

The energy multiplication requirements can be added into the constraints of the operating regimes by defining an energy multiplication factor as

$$Q = \frac{P_f}{P_{OH} + P_{\text{aux}}},$$

where P_f is the fusion power.

2.9.4.3 Determination of the Ash Level

The amount of the alpha ash in a plasma is determined by the equation for quasi-neutrality

$$n_i + 2n_\alpha = n_e \quad (2.9-8)$$

and the steady-state alpha-particle density balance equation

$$\frac{n_\alpha}{\tau_\alpha} = g_\alpha \frac{n_i^2}{4} \langle \sigma v \rangle, \quad (2.9-9)$$

where g_α describes profile effects and τ_α , the alpha confinement time, is assumed to be proportional to the plasma particle confinement time. That is,

$$\tau_\alpha = \text{constant} \times \tau_p = f_\tau \times \tau_p. \quad (2.9-10)$$

Here we have introduced the important parameter f_τ as the ratio of alpha to fuel particle confinement times. We also assume that $\tau_p = 3\tau_E$. The form of τ_E will be determined by the scaling relation used later.

By solving Eqs. (2.9-8) and (2.9-9), we obtain the ash density as

$$n_\alpha = \frac{n_e}{2} - \frac{-1 + \sqrt{1 + 2C\langle\sigma v\rangle f_\tau \tau_p n_e}}{2C\langle\sigma v\rangle f_\tau \tau_p}. \quad (2.9-11)$$

Hence, the alpha density level increases as the confinement time ratio, f_τ , increases.

Substituting Eqs. (2.9-5), (2.9-7), and (2.9-11) into Eqs. (2.9-3), (2.9-4), and (2.9-5), we obtain contours that define the operating regions (or operating windows) in n_e - T space. For various values of f_τ , we obtain different contours. Such a contour presentation is similar to the ignition plots in Refs. 52 and 54, but the latter did not include the ability to vary the ash density ratio. Further, in order to study the effect on operation other than ignition alone, we have added energy multiplication requirements.

2.9.4.4 Application to TIBER

There are several scalings for the plasma confinement time. Because our interest is to show the trends, not to calculate the absolute numbers, we will use Kaye-Goldston H mode scaling [52] as the example in the following discussions. Then τ_E takes the form

$$\tau_E = (\tau_{EOH}^{-2} + \tau_{Eaux}^{-2})^{-1/2} \quad (2.9-12)$$

where

$$\tau_{EOH} \approx 0.105 m a B_0 R_0 \quad (2.9-13)$$

and

$$\tau_{Eaux} = \frac{R_0 q_\star}{1.5 B_0} \frac{f_{LH}^{2.4}}{m^{0.75} T^{1.38}} \left(\frac{a B_0^2}{q_\star} \right)^{1.5} \left(\frac{1 + \kappa_\star^2}{2} \right)^3 \frac{(a/R_0)^{0.65}}{\kappa^{0.7} a^{0.2} q_\star^{1.7}} \quad (2.9-14)$$

In TIBER, there is an auxiliary power input. So the driven windows (i.e., $P_{aux} > 0$), together with the ignition windows (i.e., $P_{aux} = 0$), are considered. The parameters used for TIBER are given in Table 2.9-3.

2.9.4.5 Results

The change of the operating windows due to the ash effect is examined under two different cases of the constraints. Case I is the upper limit case with both higher density limit and higher total beta limit. Case II is the lower limit case with lower density limit and lower total beta limit.

Case I:

$$\begin{aligned} n_e &\leq \langle n_{GR} \rangle \\ \beta_{tot} &\leq \beta_{crit} \end{aligned}$$

Case II:

$$\begin{aligned} n_e &\leq \langle n_{mu} \rangle \\ \beta_{tot} &\leq 0.75 \beta_{crit} \end{aligned}$$

In general the ignition window gets smaller as f_τ increases. Finally, the window closes when f_τ reaches a certain value, depending on the limits (i.e., the case) assumed. Although the driven windows can be expanded from the ignition windows by adding in auxiliary power, as shown by comparing Figs. 2.9-4 and 2.9-5, they still close as f_τ increases. The values of f_τ and the ash density ratio for closing the ignition window or the driven window are given in Table 2.9-4 for the two cases considered.

Figure 2.9-4. The crossed area is the driven window with $Q \geq 7$ at $f_\tau = 4.5$ where the ignition window is closed for case I.

Figure 2.9-5. The crossed area is the driven window with $Q \geq 6$ at $f_r = 1.0$ where the ignition window is closed for case I.

Table 2.9-3. Plasma parameters.

Parameter	TIBER II
Major radius, R_0 (m)	3.0
Minor radius, a (m)	0.834
Elongation, κ	2.4
Triangularity, δ	0.4
Toroidal field, B_0 (T)	5.5
Plasma current, I (MA)	8
Axis safety factor, q_0	1.24
Cylindrical equivalent safety factor, q_*	2.576
Murakami density limit, (10^{20} m^{-3})	1.07
Greenwald density limit, (10^{20} m^{-3})	1.831
α_n	1.02
α_T	0.58

Another problem here is that if the ash obeys (even partially) classical transport laws, it will tend to diffuse inward until a large concentration accumulates on the center line. Thus initially, no ash might diffuse outward, so that initially $\tau_a \sim \infty$! Then, the time needed to reach a density ratio of 14% is about 8 seconds. Also, the value of f_r should slowly decrease from a high initial value, approaching an asymptotic limit for steady-state operation. The key question, however, is whether or not one can get over the initial stage of high f_r values.

2.9.4.6 Possible Control Schemes

There are several possible ash control methods listed in Table 2.9-5. Preliminary studies to date concern beam-driven flow reversal [55] and electric field pumping [56]. However, more study is needed for the power requirement of the beam-driven flow reversal method and an effective way to generate the radial electrical field needed for electrical field pumping. Moreover, while the central ash density can be changed by applying the above methods, extra work is needed to find how to effectively sweep the ash out of

Table 2.9-4. Results for Cases I and II.

	Case I (upper limit case)	Case II (lower limit case)
$f_{\tau i}$	~ 4.5	~ 1.0
$f_{\tau D}$	~ 14.0	~ 7.0
$f_{\alpha i}$	$\sim 20\%$	$\sim 5\%$
$f_{\alpha D}$	$\sim 30\%$	$\sim 14\%$

$f_{\tau i}$ = the minimum value of f_{τ} for closing the ignition window.

$f_{\tau D}$ = the minimum value of f_{τ} for closing the driven window with $Q \leq 6.0$.

$f_{\alpha i}$ = the minimum value of the density ratio (n_{α}/n_e) at $T \approx 16$ KeV for closing the ignition window.

$f_{\alpha D}$ = the minimum value of the density ratio (n_{α}/n_e) at $T \approx 16$ KeV for closing the driven window with $Q \geq 6.0$.

the edge of the plasma and into the divertor scope-off region. Little information is yet available for the other techniques listed in Table 2.9-5. Further study and a comparison of these methods for application to TIBER are planned for the next study period.

2.9.4.7 Conclusions

The operation window is very sensitive to the alpha ash density. For example, we find that for the TIBER with Kaye-Goldston H mode scaling, the ignition window constrained by the Murakami density limit and 0.75 Troyon beta limit is closed for $f_{\tau} = \tau_{\alpha}/\tau_p = 1.0$ (i.e., $\sim 5\%$ ash at $T \approx 16$ KeV). If auxiliary power is used to allow sub-ignition operation, the window is expanded somewhat but still closes at $f_{\tau} \approx 7.0$ (i.e., $\sim 14\%$ ash at $T \approx 16$ KeV) with $Q \leq 6$. Even for the case with the Greenwald density limit and Troyon beta limit, the ignition window still closes at $f_{\tau} \approx 4.5$ (i.e., $\sim 20\%$ ash at $T \approx 16$ KeV) and the driven window with $Q \geq 6$ closes at $f_{\tau} \approx 14.0$ (i.e., $\sim 30\%$ ash at $T \approx 16$ KeV). The details of ash transport (i.e., the value of f_{τ}) are very complicated, so the expected value remains uncertain. However, if good energy confinement were obtained, approaching neoclassical condition, a very

Table 2.9-5. Possible ash control methods.

Beam-driven flow reversal pumping
Electric field pumping
Controlled sawtooth pumping
Controlled instability pumping
Controlled wave pumping

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large f_T would be predicated. Consequently, it appears that an ash control technique will be essential for operation of a TIBER-type device, and a study of such techniques will be the main focus of the next work period.

2.9.5 TIBER BURN CONTROL

2.9.5.1 Introduction

The region of plasma parameter space of interest in engineering test reactor conceptual designs is often unstable against disturbances in plasma temperature. Present understanding of plasma transport indicates that in such unstable regions, plasma self-heating due to alpha particle thermalization can exceed power losses due to transport, radiation, etc. Thus, the plasma undergoes a thermal excursion until another higher temperature equilibrium is achieved or the discharge is terminated by a disruption. This type of plasma instability is commonly referred to as a "thermal instability" and is important because either consequence of an excursion yields substantially reduced, and probably unacceptable, overall reactor performance in terms of neutron fluence, availability, etc.

A number of burn-control issues need to be addressed within the context of the TIBER design study. First, a control method appropriate to TIBER should be selected to stabilize the plasma against possible thermal disturbances. Toward this end, a systematic, qualitative survey of previously proposed methods was conducted and candidate methods were identified. Second, the general applicability and performance of the candidate methods must be determined for application to TIBER. Toward this end, a specialized, time-dependent plasma model was developed and control scenarios were simulated.

Third, required refinements in analysis should be identified in order to assure that confidence in the control performance be commensurate with the mid-term development path of an engineering test reactor.

2.9.5.2 Survey, Comparison, and Selection of Methods

Many workers have investigated methods of controlling against thermal excursions (see, e.g., [57]). Control methods suggested include: enhanced transport via toroidal field ripple, plasma current control, refueling control, fuel polarization, additional plasma heating, adiabatic compression/decompression, impurity radiation, soft beta-limit saturation, and alpha "ash" density control. The physical basis and control strategies of these methods are briefly summarized in Table 2.9-6 (one may consult citations in Ref. 57 for additional information). In order to select an appropriate control method, the proposed methods were qualitatively reviewed.

Criteria were identified for the comparison and selection which were felt to reflect the significant requirements for application on TIBER. The criteria fall roughly into three categories--physics, technology, and operational aspects--and are detailed in Table 2.9-7. Each control method was evaluated with regard to the criteria as either favorable, neutral, or unfavorable with respect to the physics and design implications to TIBER. The results of this evaluation are given in Table 2.9-8.

Two control methods were selected for further consideration. A principal method was selected which was intended to represent a low risk and conservative option. An alternative method was also selected which would embody significantly enhanced desirability while perhaps being less conservative.

In selecting the principal method, the physics criteria were given significantly more importance than the other categories in order to reflect the conservative selection philosophy. Based on this approach, control via combined heating and refueling was identified as the principal method. The physical basis for this method exists on two observations. First, a plasma in a sub-ignited equilibrium requires a finite amount of auxiliary heating. In such an equilibrium, plasma self-heating may be suppressed by a reduction of the auxiliary heating. Second, injection of cold fuel causes the plasma to redistribute its energy with the cold fuel, resulting in a lower plasma temperature at thermal equilibrium, thus suppressing the self-heating.

Table 2.9-6. Summary of burn-control methods.

Control method	Physical mechanism	Control strategy
Toroidal field ripple	Ion energy transport	Magnetic field well enhances ion energy transport. If not strong enough, well depth may be actively controlled.
Plasma current control	Gross energy confinement	Vary I_p to affect energy confinement losses.
Refueling control	Fusion rate and injected cold fuel thermalization	Injected $n(\text{tritium})/n(\text{deut})$ influences fusion rate. Cold fuel thermalization reduces plasma temperature.
Fuel polarization	Fusion rate	Fusion reactivity enhanced by polarization may be varied by controlling the density of polarized fuel ions.
Additional heating	Energy source	Decrease/increase level of auxiliary heating from nonzero steady-state input.
Adiabatic compression/decompression	Energy source/sink	Vary plasma volume to effect heating/cooling due to adiabatic compression/decompression.
Impurity radiation	Radiative energy loss	Inject impurities to control line, recombination, bremsstrahlung, and cyclotron radiation.
Soft beta-limit limit.	Gross energy confinement	Degrade energy confinement by operating at soft beta-limit.
"Ash" density control	Fusion rate	Control fusion rate via dilution of fuel by helium ions.

Table 2.9-7. Evaluation of burn-control methods.

Criteria	Methods (see key below)											
	1	2	3	4	5	6	7	8	9	10	11	12
Exp. verf.	+	0	0	(0,0)	+	0	+	0	-	0	-	+
Param. flex.	0	0	0	(-,0)	-	-	-	-	-	-	0	0
Theor. uncer.	+	0	0	(-,-)	+	0	0	-	0	0	0	+
Tech. uncer.	0	0	0	(-,+)	+	0	0	-	+	0	0	0
R & D req'd.	0	+	+	(-,+)	+	0	+	-	+	+	-	0
Power req'd.	0	+	+	(-,+)	0	-	+	0	+	0	0	+
Comp. cost	+	+	+	(-,+)	-	0	+	-	+	0	0	+
Resp. tim	+	0	0	(-,0)	+	0	0	0	0	-	-	+
Coupled eff.	0	0	0	(-,0)	-	-	-	0	-	0	0	-
Complexity	-	0	0	(-,+)	0	-	+	0	+	+	0	-
Failure cons.	0	0	0	(0,0)	0	0	-	0	-	-	0	0

1 = Additional heating

7 = Impurity concentration

2 = Density

8 = Fuel polarization

3 = Fuel composition

9 = Soft-beta limit

4 = Ripple (active, passive)

10 = Controlled sawteeth

5 = Compression/decompression

11 = Ash control

6 = Current control

12 = Heating + refueling

+ = Favorable

0 = Neutral

- = Unfavorable

Passive control via toroidal field ripple was identified as the alternative method. In this scenario, thermal ions may be trapped in magnetic field wells created by discrete toroidal field coils. Here, in accordance with neoclassical transport theory, they complete numerous banana orbits and may rapidly drift out of the plasma, enhancing the thermal diffusivity. Estimates [58] have shown that this enhanced conductivity scales approximately according to $\chi_I \sim T^{3/2}$, which may be sufficient to suppress thermal instabilities in higher temperature plasmas ($T > 20$ keV). This method was

selected because, although the physics uncertainties were significant, its desirability was substantially enhanced by the passive nature of the control and because the required hardware would already be present in the form of discrete toroidal field coils.

2.9.5.3 TIBER Burn Control Simulations

2.9.5.3.1 Description of Plasma Modeling and Simulation Code

In order to determine control requirements and performance, a computer code was developed to simulate the plasma/controller interaction. The resulting code incorporated the significant physics embodied in the MUMAK [59] code, including: profile-averaging; lower hybrid, neutral beam, and electron cyclotron current-drive and efficiencies; ash and arbitrary impurity dilution; alpha thermalization; electron-ion drag; detailed plasma geometry; calculation of all beta components; numerous confinement scalings (Kaye-Goldston with alpha power contribution being the baseline); and time dependence. The code then used MUMAK output to determine various TIBER plasma equilibria; disturbances were initiated and control responses followed.

2.9.5.3.2 Transient Response Characteristics

Two configurations will be discussed: (1) the TIBER baseline case and (2) TIBER with inductive/lower hybrid current-drive cycles. A significant parameter describing the driven nature of these configurations is $Q = P_{fus}/P_{aux}$. Case 1 is strongly driven, with $Q = 5$, while Case 2 is somewhat less strongly driven, $Q = 12$. In order to simulate the effects of thermal disturbances, the electron and/or ion temperatures were offset by 10% of their equilibrium values, and the electron and ion temperatures were followed as they evolved. It should be noted that a great deal of uncertainty exists with regard to the magnitude and mode of such disturbances. For the purposes of this analysis, offset of both electron and ion temperatures yielded the most pessimistic results; therefore, design recommendations were based on this disturbance mode.

The ion temperature evolution for the baseline case with no burn control is shown in Fig. 2.9-6. It should be noted that, in the absence of control, the ion temperature settled to within 1% of equilibrium in 1.5 seconds. The plasma transport power losses were sufficient, in this case, to stabilize the

Figure 2.9-6. Plasma ion temperature evolution for 10% offset of ion and electron temperatures at time = 0.7 seconds. TIBER baseline case. No feedback control.

operation point because the alpha-particle-induced heating was not a dominant source. Thus, the plasma was considered stable against thermal disturbances.

The ion temperature evolution for Case 2 without burn control is shown in Fig. 2.9-7. In this case, it should be noted that the ion temperature settled to within 1% of equilibrium in 6 seconds. Although this case is also stable against thermal disturbances, it should be noted that this settling time has quadrupled from that of the baseline case. This trend continues for increasing values of Q until an instability threshold is reached at approximately $Q = 30$.

2.9.5.3 Conclusions

On the basis of the analysis presented, both TIBER cases considered are stable against thermal disturbances. Within the validity of the zero-dimensional modeling and the assumptions implicit in this modeling, no control system is required to stabilize the plasma against thermal disturbances, unless plasma response characteristics conflict with other TIBER operational requirements.

2.9.5.4 Future Requirements

Although general conclusions can be drawn from the present analysis, a number of issues should be resolved, among these: control of high- Q and ignited plasmas, the effects of plasma profile evolution, and related control for safety purposes. With regard to the issue of control for high- Q and ignited plasmas, analysis is underway for the ignition case of TIBER. In this configuration, steady-state heating is low or nonexistent, thus control authority with additional heating is degraded and control via refueling alone or an alternative method will be required. In either case, the plasma is unstable against thermal disturbances and requires control to allow stable operation. Two issues regarding plasma profile effects must be addressed. First, a more detailed, one-dimensional transport analysis is required to self-consistently account for profile effects. Second, methods of controlling plasma profile characteristics and their effectiveness in burn-control must be ascertained. Realistic plasma profiles (considering MHD equilibrium as well as transport) for TIBER should be determined, and their influence, as well as various disturbance modes on thermal stability and control, must be determined to insure sufficient confidence in thermal control aspects of TIBER design.

Figure 2.9-7. Plasma ion temperature evolution for 10% offset of ion and electron temperatures at time = 1.3 seconds. TIBER contingency case. No feedback control.

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2.10 DISRUPTION MODELING

2.10.1 INTRODUCTION

Plasma disruptions in tokamaks represent a particularly challenging design and safety problem. This off-normal event can occur with such frequency in experimental tokamaks that withstanding the effects of a plasma disruption is an important design consideration not only for the effect of a single event but for the cumulative effect over the anticipated lifetime of the device.

During a plasma disruption, plasma confinement is lost, resulting in the rapid release of the plasma thermal energy, which for the baseline TIBER II plasma (i.e., 6% beta, 10-MA current) represents approximately 109 MJ of energy. This phase of the plasma disruption is called the thermal quench. Subsequent to the thermal quench, the plasma magnetic energy is lost by inductively displacing the plasma current into adjacent conductive structures. The magnitude of the magnetic energy for the baseline TIBER II plasma is about 250 MJ. This phase of the disruption is called the current decay. The consequences of the release of these energies are surface erosion by vaporization of plasma interface components and excessive forces generated by eddy currents flowing in these structures. Since highly elongated plasmas are intrinsically susceptible to vertical instabilities, the severity of these consequences can be compounded for the TIBER II. A worst case scenario would occur when the plasma moves off midplane and is not restored before the disruption begins. The result is a disrupting plasma that simultaneously undergoes a vertical instability.

Since the understanding of the physics of plasma disruptions is still evolving, predictions of the time frame for the thermal quench and current decay cannot be predicted from present theory. However, based on experience gained from the operation of existing tokamaks, time scales for these phases were proposed for the purpose of this design study, as presented in Table 2.10-1. Disruption calculations based on the parameters of Table 2.10-1 were performed with the DSTAR code [1] to determine the magnitude of eddy currents, the resulting structural forces, the extent of surface erosion of plasma interface components, and the effect that the adopted disruption parameters and a vertical instability have on these predictions. Section 2.10.2 contains a brief description of the capabilities of the DSTAR code and the approach taken in simulating plasma disruptions for the TIBER II

Table 2.10-1. Adopted TIBER II plasma disruption parameters.

<u>Time scales</u>	
Thermal quench	- 0.1 ms
Current decay	- 10 ms (1.0 MA/ms decay rate)
<u>Energy partitioning</u>	
Thermal	
70%	(76.3 MJ) particle heating of divertor
30%	(32.7 MJ) uniformly radiated to first wall
Magnetic	
33%	(82.5 MJ) Joule heating of the plasma lost to the first wall as 70% (57.8 MJ) radiant heating and 30% (24.7 MJ) particle heating
67%	(167.5 MJ) displaced currents.

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design. Section 2.10.3 describes the DSTAR model developed for these analyses. Section 2.10.4 contains a discussion of the results. Conclusions drawn from this study are contained in the final section, Section 2.10.5.

2.10.2 DSTAR CODE DESCRIPTION/APPLICATION

The DSTAR code represents a comprehensive model for predicting the consequences of plasma disruptions. This code has evolved by combining existing capabilities for addressing the following issues:

- Evolution of unstable regions of the plasma.
- Transport and deposition of the plasma thermal energy.
- Motion of the plasma as a result of background magnetic fields.
- Interaction of the plasma with impurities resulting from structural erosion.
- Vacuum vessel eddy currents and resulting forces.
- Armor material erosion by vaporization.

The major components of the DSTAR code are:

- The Tokamak Simulation Code (TSC) [2] developed by the Princeton Plasma Physics Laboratory (PPPL) for the electromagnetic evolution and transport of the plasma.
- An "average" ion atomic physics model [3] developed by PPPL for predicting impurity ionization and emitted radiation rates.
- A two-dimensional diffusion model for impurity transport in plasmas developed at the Idaho National Engineering Laboratory (INEL) for predicting the evolution of the impurity/plasma radiant interaction zone.
- An armor tile thermal model also developed at the INEL for determining surface temperatures and vaporization.

The DSTAR code was used in two modes in analyzing the disruptive event for the TIBER II design. The code was initially run with the wall impurity interaction portion of the code bypassed. The pressure history of the plasma was specified through user input to decrease linearly over the adopted thermal quench time to a value that was found to give the adopted current decay. Once determined, this pressure history was also applied to a plasma that was initially displaced off the midplane in order to simulate a disrupting plasma that undergoes a vertical instability. The intent of these analyses was to determine the eddy currents and forces that develop in the poloidal field/equilibrium field (PF/EF) coil cases and first wall/blanket/shield (FW/B/S) structures and to determine the impact that a vertical instability has on these currents and forces.

A second set of calculations was performed where the wall-impurity interaction portion of the code was not bypassed. In this portion of the code, the diffusion of energy from the plasma due to the stochasticity of the magnetic field is based on the following estimate for the plasma diffusion coefficient: $D = \lambda^2 / \tau$, where $\tau = L / c_a$, c_a is the Alfvén speed, and L is the length of the magnetic field-lines wandering in the radial direction within the characteristic scale length. A value for this characteristic scale length was determined that resulted in the adopted thermal quench time. Once determined, the partitioning of the plasma thermal energy between particle heating and radiant heating is known, since the wall impurity interaction model of the DSTAR code predicts this partitioning. In addition, predictions

of plasma temperature, Z_{eff} , and the resulting plasma current decay are solved self-consistently based on the energy balance between resistive heating, impurity radiation, and plasma particle loss. These predictions of plasma conditions were compared with those obtained from the first set of calculations to determine the overall consistency of the adopted disruption parameters of Table 2.10-1.

2.10.3 DSTAR TIBER MODEL DESCRIPTION

Figure 2.10-1 contains a poloidal cross section of the model developed for these disruption analyses. Shown here are the poloidal magnetic flux contours in the plasma and the vacuum regions, the PF/EF coils and coil cases, and two concentric toroidal shells representing the anticipated toroidal eddy current conduction paths within the FW/B/S structures of this tokamak, as noted. The first wall structure (the inner toroidal shell) was modeled as single-turn stainless steel toroidal coils assumed to be 1 cm in thickness. The blanket/shield structure (the outer toroidal shell) was modeled in a similar manner. The coils that compose these shells are defined to be toroidal discontinuous at the same toroidal location. This discontinuity is achieved by applying an infinite gap resistance to these coils at this same location. By allowing the coils of the inner and outer toroidal shells to be electrically connected at this gap, DSTAR simulates the anticipated eddy current conduction paths for these structures, which are toroidally discontinuous sectors. Currents flowing toroidally along the first wall of a given sector, for example, would have to turn at the gap between sectors, flow radially outward, and find a return path back through the blanket/shield structure.

The PF and EF coil cases were modeled as single-turn continuous coils. The coil cases of the center PF coil stack are 7-cm thick on the outer radius, 5-cm thick on the inner radius, and 1.25-cm thick top and bottom of the coil. The EF coil cases are 10-cm thick all around the coil. The single-turn coils used to model these structures were given resistances based on stainless steel at these dimensions. The PF and EF coils were assumed to have resistances that were several orders of magnitude less than the values assigned to the coil cases.

The armor tiles of the inboard wall and divertor areas were modeled by 105 thermal segments. These segments are toroidally axisymmetric, taken to be

Figure 2.10-1. Poloidal cross section of DSTAR TIBER II model showing plasma, PF/EF coils, and blanket structure.

0.5-cm thick, and composed of graphite. Since armor tiles are not attached to the outboard wall, this wall was not modeled. The boundary conditions for heat transfer were radiative heat transfer for the inboard wall tiles to a constant sink temperature and convective cooling for the divertor tiles. The DSTAR code calculates the imposed transient surface heat flux for these thermal segments by utilizing e-folding scrape-off layer formulations. The energy e-folding distances for normal operating conditions was taken as 1.5 cm, and during the disruption this length was doubled (an adopted disruption parameter assumption).

2.10.4 DISRUPTION ANALYSIS RESULTS

The predicted equilibrium conditions for a 10 MA, 6% beta plasma are a central plasma temperature of 34.0 KeV at a density of 2.0×10^{20} per m^3 , a safety factor on-axis of 1.4 and at the edge of 3.1, a plasma thermal energy of about 109 MJ, and the magnetic field energy of about 250 MJ. The plasma axis was at 3.05 m, and the x-point was at a radius of 2.51 m and an axial position of ± 2.35 m. The poloidal cross section of this plasma as it forms in the plasma chamber is illustrated in Fig. 2.10-1. The resulting maximum steady-state heat loading on the divertor was predicted to be 3.2 MW/m^2 , resulting in a surface temperature of about 800 K. The initial temperature of the inboard tiles was approximately 1900 K. As discussed in Section 2.10.3, the disruption calculation were conducted in two parts. The first part used the DSTAR code with the wall impurity portion of the code bypassed. The second part included wall impurity effects. The results obtained are described in the subsections which follow.

2.10.4.1 Current Decay Results

Figure 2.10-2(a) contains the predicted plasma boundary at various times during the disruption, illustrating the motion and shrinkage of the plasma. Figure 2.10-2(b) contains the predicted plasma current for this disruption. As was mentioned in Section 2.10.2, this current decay was obtained by selecting an appropriate plasma pressure history. In this case, the plasma pressure was linearly decreased in 0.1 ms to result in electron temperatures of approximately 15.0 eV and was maintained at this value for the remainder of

Figure 2.10-2. (a) Poloidal cross-section showing plasma outline and (b) plasma current versus time during an on-midplane disruption.

the event. The plasma Z-eff was taken as 3 for this calculation. This 15-eV temperature is in agreement with values required for the DSTAR code to simulate measured maximum current decay rates in TFTR disruptions [4].

Figures 2.10-3(a)-6(a) contain the predicted toroidal eddy current pattern that develops in the PF/EF coil cases and FW/B/S structure, and Figs. 2.10-3(b)-6(b) contain the force distribution in the FW/B/S structure produced by these currents flowing through the poloidal field of this device. Part (a) of these figures depicts the poloidal cross section of the tokamak now rotated 90° with respect to the toroidal axis of the tokamak. For clarity, the currents of Part (a) have been separated according to the direction in which they are predicted to flow, with negative currents in the upper part of the figure and positive currents in the lower part. These eddy current patterns correspond to transient times of 0.35, 1.85, 6.85, and 9.94 ms, respectively. The currents of Fig. 2.10-3(a) are primarily the consequence of a radially inward redistribution of the plasma current density produced by the thermal quench. These patterns change as a result of inductive coupling from the decaying plasma current and because the plasma moves toward the inboard side of the plasma chamber. By the end of the disruption, the eddy currents are flowing in the positive direction in the first wall and finding return paths through the inboard FW/B/S region. The maximum eddy current in this structure is approximately 2 MA/m at 9.94 ms. The maximum structural force of 16 MN/m (a resulting pressure of about 6 atmospheres) occurs at 6.85 ms. Toroidal eddy currents in the PF/EF coil cases develop more or less linearly over the course of the disruption to a maximum of about 2.1 MA/m for the EF4 coils. The resulting PF/EF coil-case-forces distribution at 9.94 ms is given in Fig. 2.10-7. A maximum force of 240 MN/m occurs in the EF coil cases on the outboard side of the reactor. Figures 2.10-8 and -9 contain the net radial and vertical forces on the entire PF/EF coil case and FW/B/S structures. Since this disruption is symmetric with respect to the the midplane of this tokamak, the net vertical force is zero. The combined consequence of these currents is a net radial compressive load on the PF/EF coil set of 55 MN and on the FW/B/S structure of 46 MN.

Figure 2.10-10(a) contains the predicted plasma boundary for a disrupting plasma that is simultaneously experiencing a vertical instability. The associated plasma current decay is given in Figure 2.10-10(b). The influence of the vertical instability is readily apparent. The plasma moves vertically

Figure 2.10-3. (a) Toroidal PF/EF coil and blanket structure eddy current pattern and (b) blanket structure force distribution at 0.35 ms for an on-midplane disruption.

Figure 2.10-4. (a) Toroidal PF/EF coil and blanket structure eddy current pattern and (b) blanket structure force distribution at 1.85 ms for an on-midplane disruption.

Figure 2.10-5. (a) Toroidal PF/EF coil and blanket structure eddy current pattern and (b) blanket structure force distribution at 6.65 ms for an on-midplane disruption.

Figure 2.10-6. (a) Toroidal Pf/EF coil and blanket structure eddy current pattern and (b) blanket structure force distribution at 9.94 ms for an on-midplane disruption.

MAXIMUM HORIZONTAL FORCE (N) = -2.0782×10^7

MAXIMUM VERTICAL FORCE (N) = 2.4355×10^7

Figure 2.10-7. PF/EF coil case force distribution due to eddy currents resulting from a disruption on the midplane.

(a)

(b)

Figure 2.10-8. (a) Net radial force and (b) net vertical force on the PF/EF coil case structure versus time for a disruption occurring on the midplane.

(a)

(b)

Figure 2.10-9. (a) Net radial force and (b) net vertical force on the blanket structure versus time for a disruption occurring on the midplane.

Figure 2.10-10. (a) Poloidal cross-section showing plasma outline and (b) plasma current versus time during an off-midplane disruption.

downward into the lower divertor area. The rate of plasma current decay is increased by this vertical motion, resulting in a 6.75-ms current decay, which is 3.25 ms shorter than that of the on-midplane disruption.

Figures 2.10.11(a)-14(a) contain the predicted toroidal eddy current patterns for the PF/EF coil cases and the FW/B/S structure. Part (b) of these same figures contains the resulting force distribution for the FW/B/S. It is interesting to note that the current and force patterns of the FW/B/S structure are virtually identical to those of the on-midplane disruption at 0.35 ms; as is expected, since these currents arise from a redistribution of the plasma current density as a result of the thermal quench. The similarity rapidly vanishes, however, as the plasma begins to move downward into the divertor. The maximum eddy current in the FW/B/S structure is about 2 MA/m, and the maximum force is about 23 MN/m. By 6.35 ms, the eddy currents in the upper EF4 coil case reached a value of 4.5 MA/m, producing an instantaneous power from resistive dissipation of about 110 MW. This current decays with an e-folding time of 40 ms, resulting in 3.8 MJ of energy being deposited in the coil case. Given the massive size of these coil cases, the anticipated thermal rise is less than 1 K. The predicted PF/EF coil case force distribution at 6.35 ms appears in Fig. 2.10-15. A maximum force of 900 MN/m was predicted for the EF4 coil case. Figures 2.10-16 and -17 contain the net radial and vertical forces on these structures. In addition to the net radial compression experienced during an on-midplane disruption, a net vertical compressive load is predicted to occur as well. The magnitudes of these loads are 26 MN radial and 11 MN vertical for the PF/EF coil set and 50 MN radial and 10.5 MN vertical for the FW/B/S structure.

2.10.4.2 Thermal Quench Results

Figure 2.10-18(a) contains predicted plasma thermal and magnetic energy for a disruption calculation where DSTAR now self-consistently predicts plasma conditions. For this set of calculations, the characteristic scale length of the magnetic field-line fluctuations was adjusted to obtain a 0.1 ms thermal quench. This was achieved with a value of approximately 17 cm. The predicted partitioning of the thermal energy, Fig. 2.10-18(b), was 27% into particle heating of the divertor and 73% into impurity radiation heating of the interior of the plasma chamber. The fraction of plasma magnetic energy

Figure 2.10-11. (a) Toroidal PF/EF coil and blanket structure eddy current pattern and (b) blanket structure force distribution at 0.35 ms for an off-midplane disruption.

HORIZONTAL POSITION (M)

MAXIMUM HORIZONTAL FORCE (N) = $-1.1992E+06$
 MAXIMUM VERTICAL FORCE (N) = $-1.6512E+05$

(b)

Figure 2.10-12. (a) Toroidal PF/EF coil and blanket structure eddy current pattern and (b) blanket structure force distribution at 1.85 ms for an off-midplane disruption.

Figure 2.10-13. (a) Toroidal PF/EF coil and blanket structure eddy current pattern and (b) blanket structure force distribution at 5.35 ms for an off-midplane disruption.

MAXIMUM HORIZONTAL FORCE (N) = 7.4059E+07
MAXIMUM VERTICAL FORCE (N) = 8.9987E+07

Figure 2.10-15. PF/EF coil case force distribution due to eddy currents resulting from a disruption off the midplane.

(a)

(b)

Figure 2.10-16. (a) Net radial force and (b) net vertical force on the PF/EF coil case structure versus time for a disruption occurring off the midplane.

Figure 2.10-17. (a) Net radial force and (b) net vertical force on the blanket structure versus time for a disruption occurring off the midplane.

(a)

(b)

Figure 2.10-18. (a) Plasma thermal K and magnetic M energies and (b) Integrated radiant R and particle P plasma power loss terms versus time for an on-midplane disruption

predicted to result in plasma resistive heating, and subsequently in armor tile heating during the current decay, is predicted to be 90%, with 52% lost through particle heating and 48% through radiant heating.

Figures 2.10-19 and -20 contain emitted impurity radiation contours at 0.005, 0.06, 0.25, and 1.34 ms, respectively. These figures are poloidal views of the plasma and the divertors of the TIBER II as modeled by the DSTAR code. As of 0.005 ms, Fig. 20.10-19(a), the divertor tile surface temperatures had not reached a value where significant vaporization occurs. As a result, the predicted radiation contours are still at their initial values. By 0.06 ms, Fig. 20.10-19(b), the situation has changed dramatically as vaporized impurities from the divertor area have diffused through the scrape-off layer into the plasma. In addition to the impurities originating from the upper and lower divertors, impurities are being emitted from the inboard wall tiles as well. This results in the maximum radiation contours residing near the inboard wall and midway between the plasma axis and top/bottom boundary of the plasma. Although the maximum surface heat flux was predicted to be $4.5 \times 10^6 \text{ MW/m}^2$, which is an incredibly high value, this heat flux drops almost exponentially to a value of $1.0 \times 10^2 \text{ MW/m}^2$ by 1.0 ms due to impurity cooling of the plasma edge. Figure 2.10-20 illustrates the further evolution of these impurity radiation contours. The position of maximum radiant intensity shifts from the axis of the plasma to the top and bottom of the plasma. This behavior is a result of a combination of factors. Subsequent to the thermal quench, the heat load to the divertors is not sufficient to maintain the armor tile surfaces at temperatures that would produce significant vaporization. As a consequence, the rate of impurity influx to the plasma becomes less than that diffusing out of the plasma. The impurity concentration in the plasma drops, impurity radiant losses are lowered, and the plasma pressure rises due to resistive heating. The plasma once again reaches temperatures that are high enough to cause divertor vaporization, introducing more impurities into the plasma. The period of this cyclic influx is initially 2.5 ms, but the period decreases to 0.5 ms by the end of the current decay. The maximum divertor tile surface erosion was predicted to be 10 microns for this disruptive event, with approximately 250 g of armor tile material being vaporized.

Figure 2.10-19. Impurity radiation intensity contours in the plasma during an on-midplane disruption at times of (a) 0.005 ms and (b) 0.06 ms.

Figure 2.10-20. Impurity radiation intensity contours in the plasma during an on-midplane disruption at times of (a) 0.25 ms and (b) 1.34 ms.

A 10-ms current decay was not predicted for this case. Figure 2.10-21 contains the predicted plasma current, which indicates the current decay time to be 20 ms. Although the plasma temperatures correspond well with those determined for the 10 ms current decay prediction presented in Section 2.10.4.1, the Z_{eff} s differed by about a factor of two between these cases.

Two additional scoping calculations were performed with different magnetic turbulence scale lengths to determine the effect this parameter has on the predicted thermal quench time, the current decay time, and the plasma energy partitioning. The scale lengths examined were 24 cm and 10 cm, respectively. A thermal quench of 0.07 ms and current decay of 16 ms resulted for a scale length of 24 cm. The more rapid decay for the 24 cm case is due to lower predicted plasma temperatures. The partitioning of thermal energy of the plasma was 60% into particle heating of the divertor and 40% in radiant heating of the interior of the plasma chamber. The partitioning of the plasma magnetic energy is that almost 100% of the magnetic energy goes into heating the plasma, of which 48% is lost by particle heating and 52% by radiant heating. A thermal quench of 0.3 ms and current decay time of 23 ms resulted for a scale length of 10 cm. The partitioning of the plasma thermal energy is 11% to particle heating and 89% to radiant heating. The reason the 24-cm scale length resulted in a higher fraction of thermal energy going into particle losses is that more of the plasma thermal energy diffuses into the divertor before the impurities can diffuse into the plasma. The partitioning of the magnetic energy is that 65% of this energy goes into heating the plasma, of which 31% is lost by particle heating and 69% by radiant heating.

2.10.5 DISRUPTION MODELING CONCLUSIONS

Disruption calculations have been performed for the TIBER II design with the DSTAR code. Induced eddy currents and resulting structural forces have been determined for a plasma disruption that has a 0.1-ms thermal quench and a 10-ms current decay. The desired current decay was obtained by selecting an appropriate plasma pressure history and Z_{eff} (a plasma temperature of 15 eV and a Z_{eff} of 3, after the thermal quench). Table 2.10-2 summarizes these results by presenting the calculated maximums. Since the TIBER II plasma will be susceptible to vertical instabilities, a calculation was included for a disrupting plasma that simultaneously undergoes a vertical instability. These

Figure 2.10-21. Plasma current decay predicted by a self-consistent treatment of wall impurity radiant losses.

Table 2.10-2. Summary of TIBER II plasma disruption results.

	On-midplane ^a	Off-midplane ^b
Current decay time (ms)	10	6.75
Maximum eddy currents (MA/m)		
PF/EF coil case	2.1	4.5
FW/B/S	2.0	2.0
Maximum eddy current force (MN/m)		
PF/EF coil case	240	900
FW/B/S	16	23

^aPlasma pressure history specified to give linear decay in 0.1 ms to plasma $T_e = 15$ eV and constant thereafter, $Z_{\text{eff}} = 3$.

^bOn-midplane plasma pressure history plus a vertical instability.

results are also summarized in Table 2.10-2. Of particular interest is that the current decay is accelerated when the vertical instability is included.

Since the DSTAR code has the capability of self-consistently determining plasma pressure based on plasma transport, impurity transport, impurity radiation, and resistive heating, predictions of plasma current decay time, temperatures, Z_{eff} , and energy partitioning are simultaneously made. These parameters are summarized in Table 2.10-3 where a comparison has been presented with the adopted disruption parameters of Table 2.10-1. Of interest here is that the characteristic scale length of magnetic turbulence required for a 0.1 ms thermal quench ($\lambda = 17$ cm) results in plasma conditions which produced a 20-ms current decay, a thermal plasma energy partitioning of 70% radiant and 30% particle, and 90% of the magnetic energy going into plasma heating with the subsequent loss of this heating being about equally split between particle and radiant heating. This result is in contrast to the adopted values of a 10-ms current decay, a thermal energy partitioning of 70% particle and 30% radiant, and 33% of the magnetic energy going into heating, with 30% being particle and 70% radiant heating.

These differences are significant, but they are not as far apart as they would at first appear to be for three reasons. First, since only a factor of

Table 2.10-3 Comparison of adopted parameters and DSTAR results.

Adopted value	DSTAR prediction (Pressure history calculated)		
	Characteristic scale length (cm)		
	17	24	10
<u>Time Scales</u>			
Thermal quench (ms)			
[0.1]	0.12	0.07	0.3
Current decay (ms)			
[10]	20	16	23
<u>Partitioning of Kinetic Energy</u>			
Particle heating			
[70%]	27%	60%	11%
Radiant heating			
[30%]	73%	40%	89%
<u>Partitioning of Magnetic Energy</u>			
Portion into Heating			
[33%]	90%	100%	65%
Particle heating			
[30%]	52%	48%	31%
Radiant heating			
[70%]	48%	52%	69%

two in plasma resistivity is required for reconciling the current decay, this is well within modeling uncertainty. Second, had the impurity concentration in the plasma been higher during the current decay, plasma temperatures would have been lower and Z-eff higher. The combined affect would be an increase in plasma resistivity and a faster current decay. A higher plasma impurity concentration and a more rapid current decay would also change the fraction and partitioning of magnetic energy. Finally, the partitioning of the thermal energy is very sensitive to the assumed thermal quench time. This effect can

be seen in the variation of thermal energy partitioning with characteristic scale length as presented in Table 2.10-3.

It should also be mentioned that the adopted disruption parameters of Table 2.10-1 are only best estimates, and that while a complete validation of the DSTAR/TSC codes has not as of yet been accomplished, these codes have been very successful in predicting plasma current decay and eddy currents from existing experiments [4].

The conclusions to be drawn are that the adopted parameters represent estimates appropriate for initial design analysis, and that agreement between predictions and the adopted parameters is within present modeling uncertainty. Efforts to improve and validate these models are presently ongoing, and are expected to be factored into future disruption analyses.

References

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2.11 DISRUPTION CONTROL

2.11.1 INTRODUCTION

Major disruptions in tokamaks are widely believed to be caused by resistive kink (tearing) instabilities with helicity numbers $m = 1, 2, 3$, and $n = 1$. These modes are triggered by free-energy sources associated with unfavorable plasma-current profiles at magnetic surfaces with $q = m/n$, where q is the safety factor. In particular, growth of the $m = 2$ MHD oscillations has frequently been identified as a precursor signal to the onset of a major disruption. Thus, suppression of the magnetic island growth caused by this tearing mode is the most direct method of disruption control. Another plausible method is to operate the tokamak with $q(r = 0) < 1$ by externally stabilizing the sawtooth oscillations ($m = 1$, $n = 1$) such that $q(r = a) < 2$ [1].

An important task during our disruption-control studies is to identify the potential for improved plasma performance, once a viable disruption-control method has been determined. The dominant effect is a reduction of thermal shocks on the first wall and electromagnetic stress on the in-vessel conducting components, resulting in longer reactor lifetimes and less stringent engineering design criteria. The reactor may operate with a lower safety factor and a higher current, giving rise to the possibility of a higher critical beta for ballooning instabilities and a modest improvement in confinement. Operation with higher densities may also be possible.

Glasser, Furth, and Rutherford [2] first suggested stabilizing the MHD modes by controlling the current profile locally at the $q = m/n$ surfaces with a conducting shell. Chan and Guest [3] proposed the use of localized EC (electron-cyclotron) heating for this purpose and demonstrated its effectiveness for reasonable microwave power densities and pulse durations. Recently, Yoshioka et al. [4], Rutherford [1], and Ignat et al. [5], separately explored the possibility of the use of localized current drive and heating with LH and EC waves. Concurrently, White et al. [6] examined various feedback schemes for controlling resistive kink modes, and the effects of local heating and current drive feedback were derived. The most convincing demonstration of $m = 2$ mode suppression was performed on the T-10 tokamak [7], where pulsed EC heating at the $q = 2$ surface was applied. In this experiment, a set of magnetic probes was used to track the $m = 2$ mode intensity, and the

probe output is proportional to $\omega_2 A_2$ with A_2 and ω_2 being the mode amplitude and frequency, respectively. Shown in Fig. 2.11-1(a), when the EC resonance is inside the $q = 2$ surface, the probe signal increases due mainly to ω_2 growth. However, as depicted in Fig. 2.11-1(b), as the EC resonance is moved to the outside of the $q = 2$ surface, the probe signal drops by a factor of 10 as a result of a decrease in A_2 . The suppression of the $m = 2$ mode is clearly demonstrated, while the location of the microwave absorption is also crucial in making this scheme effective. Essentially, flattening the current profile locally stabilizes the $m = 2$ tearing mode.

2.11.2 EFFECT OF LOCALIZED RF HEATING AND CURRENT DRIVE ON $m \geq 2$ TEARING MODES

Most of the work on disruption control with RF waves assumes that the bulk plasma current is driven by the OH primary, for which the thermal electrons are carrying the current. In this section, the effect of RF absorption and current drive in the vicinity of the $q = m/n$ surface on the tearing modes is examined in a plasma in which the bulk current is driven by a combination of OH and RF power. This is of particular interest to TIBER, where a variety of operating modes are being considered involving current ramp-up and flat-top with OH and/or RF.

Following the analysis of White et al. [6], which is strictly valid only for a tokamak equilibrium with circular flux surfaces, we write:

$$\vec{B} = \nabla \psi \times \hat{z} - \left(\frac{n r}{m R}\right) B_z \hat{\theta} + B_z \hat{z} \quad , \quad (2.11-1)$$

where ψ is the helical flux and m, n are the poloidal and toroidal mode numbers, respectively. The evolution of ψ is governed by Faraday and Ohm's law:

$$\frac{\partial \psi}{\partial t} + \vec{v} \cdot \nabla \psi = - \eta (j - j_{RF}) \quad , \quad (2.11-2)$$

where j is the total current, $\vec{v}$ is the fluid velocity, j_{RF} is the RF-driven current, and η is the plasma resistivity. Through Ampere's law, j is expressed as:

$$\mu_0 j = - \nabla^2 \psi - \left(\frac{2n}{m R}\right) B_z \quad . \quad (2.11-3)$$

Figure 2.11-1. Effect of electron-cyclotron heating on $m = 2$ oscillations in T-10 when the resonance zone was located inside (a) or outside (b) the $q = 2$ surface.

On the right-hand side of Eq. (2.11-2), the term $(j-j_{RF})$ represents the thermal (OH) portion of the current, with j_{RF} being treated as an equivalent current source. In the absence of OH, this term vanishes at steady state while, during ramp-up or ramp-down, it is nonzero and equal to the back emf current.

Assuming a perturbation of single helicity m/n , we can write ψ as

$$\psi = \psi_0 + \delta\psi(r,t) \cos(m\theta) .$$

For modulation of RF power in phase with the rotating island, we write

$$j_{RF} = j_{RF0} + \delta j_{RF}(r,t) \cos(m\theta)$$

$$j = j_0 + \delta j(r,t) \cos(m\theta)$$

$$\eta = \eta_0 + \delta\eta(r,t) \cos(m\theta) .$$

Flux surface-averaging of Eq. (2.11-1) and retention of the first-order terms yields:

$$\frac{\partial \psi}{\partial t} = -\eta_0(\delta j - \delta j_{RF}) - \delta\eta(j_0 - j_{RF}) . \quad (2.11-4)$$

A quasilinear treatment [6] of Eq. (2.11-4) around the magnetic island for $m \geq 2$ modes results in an evolution equation for $\delta\psi$:

$$\frac{\partial}{\partial t} \delta\psi = \frac{\eta_0 \Delta'(W)}{W} \delta\psi + \eta_0 \delta j_{RF} - \delta\eta(j_0 - j_{RF}) , \quad (2.11-5)$$

where W is the island width and $\Delta'(W)$ is the fractional poloidal field jump across the island. To stabilize the tearing mode, the condition $\partial \delta\psi / \partial t \leq 0$ must be satisfied, which implies:

$$\frac{\delta j_{RF}}{j_0} - \frac{\delta\eta}{\eta_0} \left(1 - \frac{j_{RF0}}{j_0}\right) \approx \frac{\Delta'(W)W}{16} \left(\frac{2q}{r_q} - 1\right)^{-1} . \quad (2.11-6)$$

For a Spitzer-like η ($\propto T^{-3/2}$), Eq. (2.11-6) is simply given as:

$$\frac{\delta j_{RF}}{j_0} + \frac{3}{2} \frac{\delta T}{T_0} \left(1 - \frac{j_{RF0}}{j_0}\right) \approx \frac{\Delta'(W)W}{16} \left(\frac{2q}{r_q} - 1\right)^{-1}, \quad (2.11-7)$$

where $q' = dq/dr$. Note that the contribution from RF heating depends on the sign of the term $(1 - j_{RF0}/j_0)$. For an OH-driven current, corresponding to $j_{RF0} = 0$, both current drive and heating are effective in suppressing the instabilities. In the case of a pure RF-driven current, the situation is somewhat different. At steady state, $j_{RF0} = j_0$, and heating will have no effect. For current ramp-up, $j_{RF0} > j_0$, which implies that heating and current drive have opposing effects on the island growth, whereas for ramp-down with $j_{RF0} < j_0$, they have the same effect, with the heating effect depending on the magnitude of the back emf current.

To quantify the requirements in terms of $\delta j_{RF}/j_0$ and $\delta \eta/\eta_0$, the right-hand side of Eq. (2.11-6) is estimated using typical profiles [1]:

$$j_0(r) = j_0(0)/[1 + (r/r_0)^{2p}]^{1 + 1/p}$$

$$q(r) = q(0)[1 + (r/r_0)^{2p}]^{1/p}$$

For rounded profiles where $p = 2$ and $r_0/r_w = 2$, with r_w being the wall radius, and for an $m = 2$, $n = 1$ mode, $q = 2$ at $r = r_s = 1.3r_0$. When $q(0) = 1$, $r_s \Delta'(W) = 2.5$ and, typically, $W = 0.1r_w$, leading to $dq/dr = 4.4/r_w$ at $r = r_s$. Thus, the condition in Eq. (2.11-6) becomes:

$$\frac{\delta j_{RF}}{j_0} - \frac{\delta \eta}{\eta_0} \left(1 - \frac{j_{RF0}}{j_0}\right) \approx 0.06. \quad (2.11-8)$$

To estimate the RF power requirement, a more thorough calculation involving coupled ray tracing, flux diffusion, and transport will be needed.

2.11.3 CHOICE OF RF WAVES

In the previous subsection, we have given a preliminary qualitative and quantitative analysis of our primary method of disruption control. We found that, depending on the source of the plasma current (OH or RF) and the phase of the plasma evolution (ramp-up, flat-top, or ramp-down), suppression of the $m = 2$ tearing mode via localized current drive and/or heating around the $q = 2$ magnetic surface is likely for perturbations $\delta j_{RF}/j_0 \sim \delta T/T \sim$ a few %. For

However, when the plasma current itself is driven by RF, the local heating effect depends on the direction and magnitude of the back emf current. At steady state and during current ramp-up, heating is ineffective, while current drive is effective for all phases of the discharge current in tokamak operation.

Various means of disruption control by stabilizing $m = 2$ modes have been suggested, including:

- LH current drive and heating.
- EC heating and current drive.
- LH current drive with EC heating [8,9].

It has also been pointed out that current drive is generally more efficient than heating [1], and pulse-modulation of RF power is more efficient than CW operation [4]. The choice of the RF candidates should be explored for efficiency and applicability to TIBER. Detailed assessment of the RF power and modulation requirements should also be carried out. This task will involve:

- Power balance within $q = 2$ island $\Rightarrow$ transport analysis.
- RF power-deposition and current-generation profiles $\Rightarrow$ ray tracing and Fokker-Planck calculations.
- Local plasma transient response $\Rightarrow$ helical flux evolution.

2.11.4 ELECTRON-CYCLOTRON CURRENT DRIVE NEAR THE $q = 2$ SURFACE

In this subsection, we show an explicit launch scenario for electron-cyclotron current drive, localized near the $q = 2$ surface in the TIBER II equilibrium magnetic field. We model wave propagation within the cold-plasma approximation by means of the TORAY ray-tracing code [10,11], and we estimate the electron-cyclotron power absorption by using Shkarofsky's weakly relativistic approximation [12]. Current-drive efficiency is calculated as in Ref. [13] and includes relativistic and trapped-electron effects. For both power absorption and current drive, we use the assumption that the electron distribution function is close to a Maxwellian.

High-efficiency current drive by electron-cyclotron waves is achieved by selectively heating high-energy electrons that have parallel velocity of one

sign. The relativistic, wave-particle resonance condition is satisfied for such electrons if the parallel index of refraction $|N_{\parallel}| = \mathcal{O}(1)$ and if $\ell\Omega/\omega < 1$, where ℓ is the harmonic number, Ω is the cyclotron frequency, and ω is the wave frequency. Typical values for these parameters are $|N_{\parallel}| = 0.6$ and $\ell\Omega/\omega = 0.9$.

Localization of power deposition and current drive to a small range of flux surfaces is accomplished by choice of the launch geometry and of the polarization. Wave absorption must be strong and occur where propagation is at a relatively small angle with respect to the flux surfaces. For any harmonic ℓ , the extraordinary mode has the stronger absorption.

An attractive launch geometry for the TIBER II study is illustrated in Fig. 2.11-2. The launch location is just outside the separatrix of the diverted equilibrium. Extraordinary-mode power at 450 GHz is launched at polar angle $\theta = 50^\circ$ and at azimuthal angle $\phi = 285^\circ$ in a cylindrical coordinate system coaxial with the tokamak axis of symmetry. Second-harmonic cyclotron resonance absorbs the waves in a region where magnetic-field and plasma parameters have the values $B \approx 7.2$ T, $n \lesssim 0.5 \times 10^{20} \text{ m}^{-3}$, and $T_e \lesssim 10$ keV. Second-harmonic power absorption is limited to the region near the $q = 2$ flux surface and amounts to 87% of the incident power. Third-harmonic absorption is not shown here but becomes significant farther along the ray, just after the second-harmonic absorption decreases; no power reaches the center of the plasma. As a result of relativistic effects the third-harmonic absorption produces negligible driven current.

Second-harmonic extraordinary-mode absorption is used here, but an alternative scenario for localized power deposition and current drive utilizes the ordinary-mode polarization and the fundamental cyclotron resonance with a wave frequency of about 225 GHz.

Figure 2.11-3 shows the effect of 10 MW of electron-cyclotron power on the current flowing in the TIBER II plasma. The plots show the toroidal current enclosed within each flux surface, where the surfaces are parametrized by the normalized poloidal flux $\tilde{\psi} = (\psi - \psi_{\text{axis}})/(\psi_{\text{edge}} - \psi_{\text{axis}})$. The integrated equilibrium current I_p (dashed curve) is characterized for $0 < \tilde{\psi} < 0.6$ by nearly constant slope, which implies that $dI_p/d\tilde{\psi}$ is almost flat. For $\tilde{\psi} > 0.6$ and, in particular, near the $q = 2$ surface at $\tilde{\psi} = 0.73$, $dI_p/d\tilde{\psi}$ decreases with

Figure 2.11-2. Poloidal projection of a ray trajectory that deposits power and drives current outside the $q = 2$ surface in a TIBER II equilibrium. The dashed curve is the $\tilde{\psi} = 0.96$ flux contour. The chain-dot curve is the second-harmonic cyclotron resonance. The asterisks indicate the location of wave absorption.

Figure 2.11-3. Enclosed current as a function of normalized poloidal flux $\tilde{\psi}$. The solid curve differs from the dashed curve due to 10 MW of second-harmonic X-mode power.

increasing $\tilde{\psi}$, indicating a gradient in the current profile that may be associated with major disruptions. The solid curve, which is calculated from the sum of the equilibrium and the driven currents, maintains the nearly constant slope to larger $\tilde{\psi}$. With electron-cyclotron current drive, therefore, the current gradient is reduced near the $q \approx 2$ surface, which may help to control disruptions and minimize transport.

Figure 2.11-3 shows that a microwave power level of 10 MW drives a localized current near $q = 2$ that is comparable to the equilibrium current near $q = 2$. Ample current can thus be driven, but tailoring a steady electron-cyclotron current to suppress $m = 2$ island growth may be difficult. A more practical scheme may be to sense the islands and modulate the microwave power in synchronism with the passage of the island 0-point [6]. Study of such a feedback scheme for disruption control is a future task.

The launch scenario described here attempts to accommodate shifts in the radial location of the $q = 2$ surface by launching from the top or bottom near the inside edge of the plasma; then a modest ($\sim 10\%$) variation of the wave frequency ω can tune the current over the largest possible fraction of the plasma cross section. Future work will quantify this tunability.

A possible secondary application of the microwave launch scenario presented here is localized heating of the plasma core during startup to help initiate the fusion burn.

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2.12 TIME-DEPENDENT STUDIES

2.12.1 INTRODUCTION--THE TIME-DEPENDENT TOKAMAK SYSTEMS CODE (TTSC)

The TTSC solves radially averaged time-dependent particle and energy balance equations together with various engineering constraints. The physics and engineering modules are time-dependent, versions of the MUMAK [1] and TETRA [2] codes as they existed in early 1987. These two codes were merged, made time-dependent, and combined with an MHD equilibrium scaling routine [3]. The MHD scaling routine determines the required poloidal field (PF) and transformer coil currents (OH) to provide equilibrium and the necessary volts-seconds.

TTSC uses the steady-state plasma data from a TETRA output file and finds a reasonable startup scenario. TTSC determines auxiliary heating and fueling requirements, volts-seconds capability, and startup time, all features which MUMAK and TETRA, by their time-independent nature, cannot readily address. In addition, TTSC provides an initial assessment of the required OH and PF coil waveforms and checks to see if $J_{B_{max}}$, B_{max} , and stress limits are exceeded.

TTSC has three basic modes of operations. In the fully time-dependent mode, time-varying particle sources for each ion species and heating sources for all species (including electrons) are known, and the density and energy equations are solved. On the other hand, the desired density and temperature behavior may be specified so the required time-varying particle and energy sources are determined. Finally, the density behavior may be prescribed and the required particle source is calculated while a realistic power source is specified and the time-varying species temperatures determined.

2.12.2 PHYSICS

The basic form of the time-dependent equations is

$$\frac{dF^\beta}{dt} = \sum_a S_a^\beta - \sum_b L_b^\beta \quad \beta = e, D, T, \alpha, H, O, \quad (2.12-1)$$

where F is either the density n or the energy $\frac{3}{2} n_p k T_p$, S_a are sources, and L_b sinks. The S_a and L_b terms are modules from the MUMAK physics code, which is described briefly in the appendixes at the end of this chapter.

TTSC can solve particle and energy equations for each ion species. However, to save computer time since the ion-ion collision time is short, we

assume that all ion species temperatures are the same and solve a global ion energy equation. In addition, since TTSC does not contain a wall model or an impurity transport model, we assume that the impurities (including thermalized alphas) are a percentage of the total ion density. The percentage for each species may be time-varying if specified. Hence the density equations for the impurities are not solved. In addition, we also assume that the ratio of deuterium to tritium is fixed (normally this ratio is 1). Therefore, only a global ion density equation and ion and electron energy balance equations are solved. The impurity densities are determined from the density of the deuterium-tritium mixture while all ion species temperatures are the same.

The deuterium-tritium density equation is

$$\frac{dn}{dt} = R_{\text{fueling}} - n_D n_T \langle \sigma v \rangle_{DT} - \frac{n}{\tau_p^i}, \quad (2.12-2)$$

where R is a particle source term, and

$$\tau_p^i = C \times \tau_E^i, \quad (2.12-3)$$

where C is a constant and τ_E^i is the global energy confinement time calculated from the MUMAK transport routines.

The electron energy balance equation is

$$\frac{3}{2} \frac{dn_e T_e}{dt} = P_{\text{ohmic}} + P_{\text{aux}} + p_\alpha^e - \frac{n_e T_e}{\tau_E^e} - \frac{n_e (T_e - T_i)}{\tau_{EQ}} - P_{\text{rad}} + P_{\text{source}}, \quad (2.12-4)$$

where P_{ohmic} is the ohmic heating source, P_{aux} is an auxiliary heating term (usually RF), p_α^e is the fraction of fast alpha power deposited in the electrons as they slow down, τ_E^e is a global electron energy confinement time, τ_{EQ} is the electron-ion equilibration time, and P_{rad} is a radiation loss term (bremsstrahlung and synchrotron). The modules which calculate these terms come from the MUMAK code and are documented in Ref. [1]. The electron and ion densities are related through charge neutrality

$$n_e = \sum_{\beta} Z_{\beta} n_{\beta} \quad \text{where } \beta = D, T, \alpha, H, O. \quad (2.12-5)$$

If the electron density, $n_e(t)$, is prescribed as either an analytic function or a set of data points, then the ion species densities are

$$n_\beta = \frac{f_\beta n_e}{Z_\beta} \quad \text{where } \beta = D, T, \alpha, H, O, \quad (2.12-6)$$

where f_β are the concentrations of the ion species such that

$$\sum_\beta f_\beta = 1. \quad (2.12-7)$$

In this first version of TTSC, these concentrations are constant in time. The f_β can be adjusted so that the steady-state value of Z_{eff} used by MUMAK and TETRA is obtained:

$$Z_{eff} = \frac{\sum_\beta Z_\beta^2 n_\beta}{n_e} = \sum_\beta Z_\beta f_\beta. \quad (2.12-8)$$

Knowing the electron density and ion species mix, the required ion particle source terms can be obtained from Eq. (2):

$$R_{fuel}^\beta = \frac{dn_\beta}{dt} + \frac{n_\beta}{\tau_p^\beta} + R_\alpha \quad \text{for } D, T, \alpha, O, H, \quad (2.12-9)$$

where

$$R_\alpha = n_D n_T \langle \sigma v \rangle_{DT} \quad \text{for } D \text{ and } T \quad (2.12-10)$$

$$= - n_D n_T \langle \sigma v \rangle_{DT} \quad \text{for alphas}$$

$$= 0 \quad \text{for impurities,}$$

and $\frac{dn}{dt}$ is determined from the specified analytic expression for $n_e(t)$ or by fitting $n_e(t)$ data points to cubic splines.

Assuming that all the ion species have the same temperature, a global ion energy balance equation can be solved to obtain the ion temperature

$$\sum_{\beta} \frac{3}{2} \frac{dn_{\beta} T_i}{dt} = \sum_{\beta} p_{aux}^{\beta} T_{\beta} \sum_{\beta} \frac{n_{\beta} (T_e - T_i)}{\tau_{ei}^{\beta}} - \sum_{\beta} \frac{n_{\beta} T_{\beta}}{\tau_E^{\beta}} + P_{\alpha} + \sum p_s^{\beta}, \quad (2.12-11)$$

where P_{aux} is an auxiliary heating source (usually RF), τ_{EI}^{β} is the ion-electron equilibration term for ion species β , τ_E^{β} is the global energy confinement time for species β , P_{α} is the fraction of alpha power deposited in the ions, and p_s^{β} is an additional energy (source or sink), which is used when the ion temperature is prescribed.

If $T_e(t)$ and the ratio of T_e to T_i prescribed, then Eqs. (2.12-4, -11) may be solved to obtain the required energy sources

$$P_{source}^e = \frac{3}{2} \frac{dn_e T_e}{dt} - P_{ohmic} - P_{aux} - p_{\alpha}^e + \frac{n_e T_e}{\tau_E^e} + \frac{n_e (T_e - T_i)}{\tau_{ei}} + P_{rad}$$

$$p_{source}^{\beta} = \frac{3}{2} \frac{dn_{\beta} T_i}{dt} - p_{aux}^{\beta} - p_{\alpha}^{\beta} + \frac{n_{\beta} (T_e - T_i)}{\tau_{ei}^{\beta}} + \frac{n_{\beta} T_{\beta}}{\tau_E^{\beta}}, \quad (2.12-12)$$

where $\frac{dn}{dt}$ is determined from the analytic expressions for $n(t)$ and $T(t)$ or by using cubic splines.

2.12. INITIAL CONDITIONS

The initial plasma current is specified (typically 200 kA). The initial density is taken to be a fraction of the Greenwald density limit given as

$$n_e = \frac{1 \times 10^{14} f_e I_p}{\pi a^2}, \quad (2.12-13)$$

where $f_e < 1$ (typically 0.2). Using Eq. (2.2-13) and the user-supplied f_{β} from Eq. (2.2-6) the initial ion densities are determined. To find the ion and electron densities, the ion and electron energy balance equations are solved assuming

$$\frac{dn_{\beta} T_{\beta}}{dt} = 0 \quad \beta = D, T, \alpha, H, O, e. \quad (2.12-14)$$

We also force the ratio $\frac{T_i}{T_e}$ to be a user-supplied initial ratio

$$RATTITE = \frac{T_i}{T_e} . \quad (2.12-15)$$

We also allow ne to adjust somewhat and solve the charge balance equation

$$n_e = \sum_{\beta} Z_{\beta} n_{\beta} \quad \beta = D, T, \alpha, H, O . \quad (2.12-16)$$

For these four equations, the corresponding four variables are T_e , T_i , n_e , and τ_{MIN} , where τ_{MIN} is the minimum energy confinement time used in the energy balance equations. The radial transport routines come from the MUMAK code, which assumes a steady-state full-sized plasma. The τ_E^{β} calculated from these routines for the initial plasma is usually quite small and greatly increases the amount of computer time needed to reach steady-state. Therefore, TTSC assumes that

$$\tau_E^{\beta} \geq \tau_{MIN} , \quad (2.12-17)$$

where τ_{MIN} is adjusted to solve the initial energy balance equations. A nonlinear equation-solver HYBRID [4] is used to solve the four equations (2.12-12, 2.12-15, 2.12-16). The initial value for n_e comes from Eq. (2.12-13) while those for T_e and T_i come from Eq. (2.2-15) and the following equation,

$$\beta = f_{\beta} \beta_{LIM} , \quad (2.12-18)$$

where f is a user-supplied number ≤ 1 (initially 0.005), and β_{LIM} is the Yamazaki beta limit with a coefficient of 0.047. When external energy and/or particle sources are prescribed, we assume that the global particle and energy confinement times obtained from the MUMAK model degrade exponentially when either the beta or density limits are exceeded. This prevents thermal runaway when the fusion power becomes significant. When either the density or temperature time-profiles are specified and the required sources calculated at steady-state, the sources may be negative. This can be interpreted as an enhancement in radial transport required to prevent thermal runaway.

2.12.4 MHD EQUILIBRIUM SCALING

Having obtained the initial T_i , T_e , n_i , n_e , and plasma current, the PF and OH coil currents need to be adjusted to provide MHD equilibrium and the proper amount of flux linkage. A unique feature of TTSC is the capability of linking the PF and OH coil waveforms with plasma performance. At each point in time, the OH and PF coil currents are adjusted to provide MHD equilibrium and correct amount of volts-seconds. The number, sizes, and locations of the PF and OH coils come from the output of the TETRA steady-state nonlinear optimization systems code discussed in a companion report

Before the plasma is initiated for a given OH current, the remaining PF coils currents are adjusted to produce (in a least-squared sense) a zero vertical magnetic field at a number of equally spaced grid points along the midplane (from 3 to 32). To determine the OH current, a guess for the initial OH current is supplied and the resulting PF coil currents determined. $B_{\max}$ and $JB_{\max}$ are calculated at the top and bottom of each PF coil and at the top, bottom, and center of the transformer stack. The initial guess is then scaled so that the currents in the OH and PF coils do not exceed the $B_{\max}$ and $JB_{\max}$ limits (typically 14 T and 500 MAT), while being as large as possible. This will allow the OH coil swing to produce the largest amount of volts-seconds. Starting from these initiation coil currents, when the plasma is present, the coil currents are then adjusted to provide MHD equilibrium and the required volts-seconds. To obtain these currents, first the vertical field at the midplane grid points (from 3 to 32) is calculated using the plasma conditions (size, position, poloidal beta, current, elongation, and inductivity) and the PF and OH coil currents and locations from a separate 2-D MHD equilibrium calculation are calculated. These reference set values are obtained from TETRA or directly from the 2-D MHD code NEQ [5]. This reference set provides the vertical field that the TTSC MHD routine then uses to scale the coil currents. For a given OH current, the PF coil currents are adjusted so that the vertical field at the grid points on the midplane match in a least-squared sense those values obtained from scaling the reference set vertical field for the new set of plasma conditions generated by TTSC. The OH current is adjusted so that the complete PF and OH coil sets provide the amount of volts-seconds required by the plasma at that point in time. At each time step, the new plasma conditions and volts-seconds requirement are given to the MHD scaling routine, which then returns the required OH and PF coil currents.

In addition, at each time step, $B_{\max}$ and $JB_{\max}$ are determined at the top and bottom of each PF coil and at the top, bottom, and center of the transformer stack. To pinpoint a possible stress problem, the maximum $B_{\max}$ and $JB_{\max}$ for the entire coil set over the entire simulation are recorded as well as the coil location where this occurs. In addition, at the end of the simulation, usually 2 to 3 seconds into the current flat-top, and presumably, but not necessarily, when the plasma is at steady-state operating conditions, the PF and OH coil currents are linearly extrapolated to the time where either the $B_{\max}$ or $JB_{\max}$ limit (whichever comes first) is met. This then determines the maximum burn time and volts-seconds capability of the given magnet set for the assumed startup scenario. If either the $B_{\max}$ or $JB_{\max}$ limit is exceeded at some point during the startup, this is noted along with the coil location and current value. The burn time and volts-seconds capability is then given as the values obtained by TTSC at the end of the simulation.

2.12.5 ENGINEERING

The radial build and PF and OH coil locations and sizes come from the TETRA code. Given the PF and OH coil currents, an estimate of the actual coil stresses can be obtained. This will be done when the new version of TETRA is incorporated into TTSC.

The auxiliary power inputs can be adjusted to minimize the amount of volts-seconds used in startup without exceeding the beta limit or the capability of the auxiliary heating power supplies.

2.12.6 POSSIBLE STARTUP SCENARIOS

For this first version of TTSC we have examined various startup scenarios for pulsed operation where RF is used as a heating source but not as a current drive source. We have also assumed that the plasma density ramp is prescribed while the auxiliary power (presumably RF) source to the electrons is specified. The resulting particle sources (or sinks) and electron and ion temperatures are determined. The plasma current is assumed to be a piecewise continuous linear ramp based on DIII-D data [5] and shown in Fig. 2.12-1. The electron density is prescribed to follow the plasma current as displayed in Fig. 2.12-2 where we have fixed Z_{eff} at 1.32. The general TIBER pulsed-mode parameters are given in Table 2.12-1, while Table 2.12-2 shows the assumed initial plasma conditions.

Figure 2.12-1. Assumed plasma current waveform. Initial value of 200 kA. Initial slope is 1 MA/s changing to $\frac{4}{9}$ MA/s after reaching 1 MA.

Fig. 2.12-2. Assumed radially-averaged electron density ramp.

Table 2.12-1. Initial plasma conditions.

TIBER parameters.	
R_0 (m)	3.0
a (m)	0.833
κ	2.4
B_T (T)	5.55
I_p (MA)	8.0
n_e ($\times 10^{20} \text{ m}^{-3}$)	1.78
Z_{eff}	1.32
β_{lim} (75% of β_{yam})	0.050
$\frac{d}{dt}$ (V-S)	0.072
t_{startup} (s)	~ 17

Table 2.12-2. Conditions at plasma initiation ($t = 3$ s).

I_p (MA)	0.200
n_e ($\times 10^{20} \text{ m}^{-3}$)	0.018
T_e (keV)	0.13
T_i (keV)	0.08

Figure 2.12-3 shows the prescribed RF auxiliary heating power waveform, in this case a maximum of 20 MW to the electrons. Table 2.12-3 shows the effects of varying the maximum amount of RF power. For RF powers less than 8 MW, the plasma does not ignite. For RF powers from 8 to 10 MW, the plasma ignites but well after the current flattops. To ignite the plasma and reach steady-state

Figure 2.12-3. Assumed auxiliary power source waveform (RF with all the power deposited in the electrons)

Table 2.12-3. TIBER with CW RF heating at $t \sim 27$ s ($t_{\text{flat top}} + 10$ s).

Case	RF power (MW)	T_e (keV)	T_i (keV)	β	P_{FUS} (MW)	Γ_{WALL} (MW/m ²)	V-S (V-S)	V-S _{RES} (V-S)	τ_E^e (s)
30.00	~ 0	1.49	1.49	0.0083	1.27	0.005	56.8	12.78	13.9
30.08	8	6.14	5.81	0.0343	150	0.621	47.9	3.98	3.7
30.10	10	8.78	7.83	0.0493	297	1.23	47.0	3.00	2.94
30.15	15	9.20	8.05	0.0515	314	1.30	46.1	2.05	1.97
30.20	20	9.82	8.04	0.0519	313	1.30	45.7	1.71	1.07

when the current flattops requires at least 15 MW of RF power. Increasing the RF power from 10 MW to 20 MW saves over 1 volts-seconds, which translates to 14 s of increased burn time.

To get a feeling for the capabilities of TTSC we examine in detail a possible startup scenario assuming a maximum of 20 MW of RF power deposited in the electrons with the heating profile shown in Fig. 2.12-3. Figures 2.12-4 and 2.12-5 show the resulting electron and ion temperatures. The initial change in slope in the temperature curves near 4 seconds corresponds to a change in confinement due to a change in the assumed slope in the current rampup. The second change in the temperature slopes near 9 seconds follows the exponential enhancement in radial losses when the beta limit is exceeded. The plasma beta and beta limit (dashed curve) are shown in Fig. 2.12-6. The final change in slope occurring near 17 seconds results from the assumed change in slope in the plasma current due to flattopping. The energy confinement time is a function of density, temperature, plasma current and auxiliary power (including alpha power) and is plotted for the electrons in Fig. 2.12-7. There are several competing effects. As the density, temperature, and plasma current increase, τ_E^e increases. However, as these

Figure 2.12-4. Radially-averaged electron temperature assuming the auxiliary power source in Fig. 2.12-3.

Figure 2.12-5. Radially-averaged ion temperature assuming the auxiliary power source in Fig. 2.12-3 with all ion species having the same value.

Figure 2.12-6. Plasma beta including fast alphas (solid line) and 75% of the Yamazaki beta limit (dashed line). Energy confinement time is exponentially degraded when the solid curve is above the dashed line.

Figure 2.12-7. Global electron energy confinement time.

increase, P_a increases, which decreases τ_E^e . Also, P_{ohmic} decreases in time, which tends to decrease the rate at which the electron energy increases. These effects, together with the exponential enhancement in losses when the beta limit is exceeded, produce the temperature curves shown in Figs. 2.12-4 and 2.12-5.

The reference magnet set which provides a 2-D MHD equilibrium during steady-state is shown in Fig. 2.12-8. TTSC represents this case as 7 sets of PF coils and an OH stack made up of 4 coils, all programmed with the same current. The numbering system for the coils can also be seen in Fig. 2.12-8. Point 16 corresponds to the top of the OH stack, while 16 is at the center. Figure 2.12-9 shows the volts-seconds required by the plasma. This volts-seconds requirement can be divided into an inductive part, primarily a function of plasma size and current (Fig. 2.12-10), and a resistive part (Fig. 2.12-11), a function of plasma density and temperature. Assuming a full-sized plasma startup, the inductive part of the total volts-seconds requirement is independent of the startup scenario. The startup scenario can greatly affect the resistive part. The less auxiliary heating, the longer the plasma takes to reach steady-state and the greater the resistive volts-seconds consumed, as shown in Table 2.12-3. However, in a reasonable startup scenario, the resistive part may amount to less than 10% of the total volts-seconds consumed, so there reaches a point of diminishing return as far as increasing the amount of auxiliary heating. To provide this volts-seconds, the PF and OH coils are charged so as to swing from a -27.66 V-s to 22.79 V-s, as shown in Fig. 2.12-12. The OH and PF coils are swung so as to prevent B_{max} and JB_{max} limits from being exceeded. The OH coil stack current waveform is shown in Fig. 2.12-12, while the B_{max} and JB_{max} at the end of the OH stack are shown in Figs. 2.12-13 and 2.12-14. Figures 2.12-15 through 2.12-21 show the current waveforms for the other PF coil sets identified in Fig. 2.12-8. The rather strange shape in many of the current waveforms between 4 and 8 seconds is a result of changes in beta poloidal, shown in Fig. 2.12-22. For fixed plasma size, the scaling is a sensitive function of poloidal beta.

In addition to examining B_{max} and JB_{max} for the OH stack, we also examined them for each PF coil set. In this particular case, the B_{max} at point 16 limited the burn time. When the slope of B_{max} for the steady-state conditions in Fig. 2.12-23 was extrapolated, then the B_{max} limit of 14 T was reached for a burn time of 321 seconds, with a total volts-seconds consumed of 74.8 V-s.

TIBER-II Model k69c

Figure 2.12-8. Reference PF and OH coil set used by TTSC and TETRA. This plot is generated by the FEDC/ORNL 2-D MHD equilibrium code NEQ. The numbers by the coils corresponds to the TTSC numbering system.

Figure 2.12-9. Volts-seconds consumed by the plasma.

Figure 2.12-10. Inductive volts-seconds consumed by the plasma, a function of plasma current.

Figure 2.12-11. Resistive volts-seconds consumed. This may vary greatly with differing startup scenarios and is a function of changing plasma conditions.

Figure 2.12-12. Volts-seconds provided by the PF and OH coils. At each time-step, the volts-seconds provided by the coils equals that required by the plasma.

Figure 2.12-13. OH transformer stack current waveform. Initial peak limited by $B_{\max}$ limit.

Figure 2.12-14. $B_{\max}$ at point 16, the center of the OH transformer stack at an r location of 0.6 m.

PF COIL-16 3.974×10^{-1} 3.309×10^2
RAOH 0.600 ZHOH 0.000

Figure 2.12-15. JB_{max} at point 16, the center of the OH transformer stack at an r location of 0.6 m.

WAVE=1 0.000 9.072×10^0

RPF= 0.800 ZPF= 1.196 RICO=-6.310

Figure 2.12-16. Current waveform in PF coil 1 located at $r = 0.8$ m and $z = \pm 1.2$ m.

WAVE=3 0.000 1.254×10^1

RPF= 0.800 ZPF= 1.996 RICO=-11.15

Figure 2.12-17. Current waveform in PF coil 3 located at $r = 0.8$ m and $z = \pm 2.0$ m.

WAVE=5 0.000 6.709×10^0

RPF= 0.800 ZPF= 2.796 RICO=-6.131

Figure 2.12-18. Current waveform in PF coil 5 located at $r = 0.8$ m and $z = \pm 2.8$ m.

WAVE=7 0.000 6.116×10^0

RPF= 1.800 ZPF= 4.246 RICO=-4.727

Figure 2.12-19. Current waveform in PF coil 7 located at $r = 1.8$ m and $z = \pm 4.25$ m.

WAVE=9 0.000 4.967×10^0

RPF= 4.095 ZPF= 4.246 RICO=-1.172

Figure 2.12-20. Current waveform in PF coil 9 located at $r = 4.1$ m and $z = \pm 4.25$ m.

WAVE=11 -1.443×10^0 0.000

RPF= 6.195 ZPF= 3.196 RICO= 0.490

Figure 2.12-21. Current waveform in PF coil 11 located at $r = 6.2$ m and $z = \pm 3.2$ m.

WAVE=13 -5.802×10^0 6.777×10^{-1}

RPF= 6.195 ZPF= 1.496 RICO=-0.312

Figure 2.12-22. Current waveform in PF coil 13 located at $r = 6.2$ m and $z = \pm 1.5$ m.

Figure 2.12-23. Poloidal beta.

These numbers are somewhat larger than the baseline parameters because we do not look at the stresses in the OH stack that may be exceeded before the $B_{\max}$ limit is met. In the next version of TTSC we hope to incorporate some analytic stress models. During steady-state, the plasma has a loop voltage of 0.072 V. Therefore, for each volts-seconds saved during startup, the burn time can be extended approximately 14 seconds. In addition, if $JB_{\max}$ was the limit, then this limit would be reached for a burn time of 30,617 seconds. In charging the coils to zero out the vertical field at plasma initiation, the largest $B_{\max}$ is at the top of coil 4. However, as the coil currents change to follow the changing plasma conditions, the highest $B_{\max}$ point quickly switches to the center of the transformer stack.

2.12.7 CONCLUSIONS

We have developed a zero-dimensional time-dependent physics-engineering code which can quickly examine various startup scenarios and, when coupled to the steady-state nonlinear optimization tokamak systems code TETRA, ascertain if the "optimal" steady-state operating conditions are "achievable." TTSC can estimate volts-seconds requirements and capabilities and auxiliary heating and fueling requirements. In addition, TTSC can address potential magnet stress problems by examining $B_{\max}$ and $JB_{\max}$ at various coil locations. TTSC has shown the effectiveness of auxiliary power in reducing the volts-seconds requirement during startup and has found a reasonable startup scenario for the pulsed-operation mode of TIBER. In addition, TTSC can estimate the burn time and volts-seconds capability of the magnet set.

The next version of TTSC will include the most up-to-date physics and engineering routines for MUMAK and TETRA. We also hope to address the effect of RF current drive (in particular, lower-hybrid) on the volts-seconds requirement and on the MHD equilibrium scaling. We also intend calculating stresses in the OH and PF coils to determine the burn time and volts-seconds capability. We will also use TTSC to examine expanding radius startup scenarios, since, presumably, at very low densities the plasma will be smaller and more circular. Finally we will examine the effect of changing Z_{eff} from approximately 5 to the steady-state value of 1.32 on volts-seconds consumed during startup.

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2.13 TIBER II: VERTICAL STABILITY

The TIBER II design has been analyzed for its vertical stability, and a system concept has been developed for stabilizing the $n = 0$ axisymmetric MHD mode. Vertical stabilization of highly elongated plasmas such as TIBER II is possible with the inclusion of passive and active current-carrying elements which are located close to the plasma. Additional work is required to address some of the secondary issues, such as the higher order MHD modes and non-linear and non-ideal plasma motion.

The following subsections address the stability issues studied. In Section 2.13.1 a plasma equilibrium representative of the baseline TIBER II parameters is developed based on the free boundary equilibrium solver, EFIT, which has been extensively used to simulate DIII--D experimental results. The robustness of the plasma to random fluctuations of the flux function, a measure of the plasma stability to external perturbations, is presented. In Section 2.13.2 the ideal MHD plasma stability issues for this equilibrium are investigated using the GATO code. Results on $n = 0$ and $n = 1$ modes are reported. Section 2.13.3 develops the optimum passive system requirements for the vertical stability of this design. The characteristics of a passive plate system capable of reducing the MHD growth rates to the L/R time constant of the passive system are presented. Finally, in Section 2.13.4, the characteristics of an active control system capable of responding on the time scale of the passive system are presented. This study was undertaken after the basic machine characteristics and design were frozen; no attempt was made to incorporate the concepts in the base design. The study does show that vertical stabilization of the plasma can be achieved with the incorporation of a minimum of additional hardware into the baseline design.

2.13.1 FREE BOUNDARY EQUILIBRIUM

The base equilibrium used in this study is computed using the full free-boundary MHD equilibrium code EFIT [1]. As in the DIII and DIII-D experiments, a set of flux loops is used to regulate the plasma shape. The equilibrium is calculated using the 16 external coils, as in the original LLNL TIBER II equilibrium, with their corresponding control flux loops located around the plasma boundary. These are shown in Fig. 2.13-1. The location of these flux

loops has not been optimized. The sensitivity of the computed equilibrium to the distance of the flux loop from the plasma surface is described in the next section. In the calculation, no up-down symmetry of plasma is assumed and finite-size coils are used. The toroidal current density is taken to have the form

$$J_T = J_0 \left[\beta_{p0} \frac{R}{R_c} + (1 - \beta_{p0}) \frac{R_c}{R} \right] (1 - \psi^n)^m.$$

Here J_0 , β_{p0} , R_c , n , and m are parameters describing current profile. In the base case, R_c is taken to be the major radius of the vessel center and n is adjusted to give an axial safety factor value $q(0)$ of 1.05. The remaining current parameters and the required flux values at the control flux loops are chosen to match as closely as possible the shape and the major plasma parameters in the original LLNL TIBER II equilibrium.

The equilibrium plasma flux surfaces, the toroidal current density J_T along the horizontal magnetic plane, and the safety factor profile for the base case are shown in Fig. 2.13-1. The major plasma parameters and the required currents and resulting poloidal flux at the external coils are summarized in Tables 2.13-1 and 2.13-2. Comparisons between the EFIT and LLNL equilibrium are also shown. Figure 2.13-2 compares the plasma boundary between the two equilibria. The dashed line is the EFIT equilibrium; the solid line and the separatrix X-point are for the LLNL equilibrium. Agreement is reasonable.

2.13.1.1 Equilibrium Sensitivity Study

This equilibrium is insensitive to the current initialization used and can be obtained over a wide range of initial plasma current distribution. However, since the plasma shape in this case has a high elongation of about 2.7 at the separatrix, the equilibrium plasma shape, particularly its vertical position, may be sensitive to the location of the control flux loops and the noises in the loops. The results of the sensitivity study are summarized in this section.

Figure 2.13-3 shows the equilibrium plasma surfaces for a case where the control flux loops are located approximately halfway between the vessel and equilibrium field coils. As the flux loops are moved closer to the coils, the

Figure 2.13-1. EFIT baseline equilibrium.

Table 2.13-1. Equilibria comparisons.

Parameter	EFIT	LLNL
R (m)	3.00	3.00
a (m)	0.84	0.86
κ	2.70	2.74
δ	0.82	0.75
R _x (m)	2.31	2.34
Z _x (m)	2.26	2.34
$\langle\beta_T\rangle$ (%)	5.97	5.96
q(0.95)	3.32	--
$\phi(0)$ (V-s)	19.1	--
$\phi(1)$ (V-s)	0.33	--
I (MA)	10.0	10.0
B _T (T)	5.60	5.55

Table 2.13-2. Coil currents and flux.

Coil no.	Current (MA)	Flux (V-s)
1	-11.9	-25
2	-7.2	-15
3	0.85	-1.0
4	4.1	5.6
5	0.98	6.8
6	3.60	29
7	-3.7	-141
8	-4.0	-161

Figure 2.13-2. Comparison of EFIT (dashed) and LLNL (solid) equilibrium.

Figure 2.13-3. Example of plasma shift with flux loop location.

plasma tends to shift vertically. This is further illustrated in Fig. 2.13-4, where the equilibrium plasma vertical position is plotted against the normalized distance of the flux loops from the first wall:

$$\rho = \frac{(R_c - R_{c0})}{(R_{c1} - R_c)},$$

where R_c , R_{c0} , and R_{c1} are the major radius of the flux loop and the corresponding ones at the first wall and at the coils. As the flux loops are moved beyond the midpoint between the first wall and coils, the plasma tends to shift vertically. This may be interpreted physically as the plasma cannot be passively stabilized with a conducting wall at more than halfway from the first wall. This result is consistent with that obtained using the GATO stability code as described in the next section.

The sensitivity of the plasma vertical position to the noises in the flux loops is illustrated in Fig. 2.13-5, where the standard deviation of the plasma vertical position is plotted against the maximum random perturbation introduced into the poloidal flux values at the control flux loop locations. Each data point is computed using 50 perturbed equilibria. The noises in the flux loops must be less than 1% for the plasma to remain vertically stable.

2.13.2. GATO IDEAL STABILITY RESULTS

Axisymmetric ($n = 0$) and nonaxisymmetric ($n = 1$) stability calculations were performed for the given TIBER II equilibrium described in Subsection 2.13.1. For the stability calculations, the separatrix was moved in from the diverted surface by 0.1% as measured in flux, to ensure that the last surface was closed.

The stability code GATO [2] is a variational code that computes the unstable eigenmodes and eigenfrequencies of a given equilibrium. The calculation is done on a flux coordinate mesh (ψ, χ), where χ is an angular poloidal coordinate and ψ is a radial coordinate, and the eigenmodes are expanded on this mesh as a sum of finite hybrid elements. Two choices of poloidal coordinate χ are possible, equal arc length or the PEST (straight field-line) coordinate defined by the Jacobian $J = \tau^2 q / R B_\phi$.

An important point to note is that, even though variational principles provide optimistic stability predictions, the finite hybrid element implementation of the variational principle is, in most cases, destabilizing.

Figure 2.13-4. Plasma vertical position variation with normalized flux loop position.

MAGNETIC COILS VARIATION

Figure 2.13-5. Plasma position variation with random perturbation of flux.

This is in contrast to the PEST-I code [3], which is always optimistic. In the limit of infinite resolution, the eigenvalues and eigenmodes converge to their true values, but, for finite resolution, stable modes can be destabilized. These modes must be eliminated either by a convergence study--that is, by extrapolating results to infinite resolution if the convergence properties are smooth--or from observation of the structure of the mode.

GATO treats either an infinite vacuum or a finite infinitely conducting continuous wall surrounding the plasma. The wall position can be defined by inputting a series of points (τ_{wi}, z_{wi}) , or conformally by

$$\rho_w(\chi) = \rho_p(\chi) + (R_w - 1)\rho_p(\chi = 0) \quad .$$

where ρ is a polar wall coordinate, the subscripts w and p represents the wall and plasma boundary, and R_w is a parameter. Passive or active stabilization by finite conducting plates cannot yet be treated. Since the TIBER II design calls for passive stabilizing plates with no effective wall, the stability results reported below give the maximum distance from the plasma that the passive plates can be placed to ensure stability.

2.13.2.1 Axisymmetric Stability

For the TIBER II design equilibrium, Fig. 2.13-6 shows the square of the growth rate γ^2 versus normalized wall position R_w . The equal arc length coordinate was used in this analysis since it produces the best results for axisymmetric stability. Comparisons using the PEST angular coordinate would place the marginal wall position 20% closer than that shown in the figure, although convergence studies result in complete agreement. These results were obtained with a fairly coarse $N_\psi/N_\chi = 30/60$ mesh resolution, but convergence studies were done near the marginal point, and the change was found to be small. A continuous wall therefore needs to be placed at about 1.58 times the plasma minor radius to ensure stability to $n = 0$. Passive stabilizing plates in place of the wall would consequently need to be closer than that.

The mode structure is shown in Fig. 2.13-7 corresponding to the case just above the marginal point at $R_w = 1.70$. Note that the assumption of a rigid vertical displacement is reasonable.

GATO $n=0$ Stability
Growth Rate vs Wall Position

Figure 2.13-6. Growth rate versus wall position for $n = 0$.

Figure 2.13-7. $n = 0$ modal shape.

2.13.2.2 Nonaxisymmetric Kink Stability

Figure 2.13-8 shows the growth rate (squared) γ^2 versus wall position R_w for $n = 1$ modes. For these modes, the PEST coordinate is the better representation, and all calculations were done with this coordinate. The safety factor on axis is $q_0 = 1.05$ and $\bar{\beta}$ is 6.2%. Resolution problems prevent meaningful results from being obtained with a coarser mesh than $N_\psi/N_\chi = 40/80$. Time and memory constraints prevent the resolution being increased above 60/120. Convergence studies are therefore not meaningful, but the unstable modes are well resolved by the 60/120 mesh. Flux surfaces are packed around integer q values to properly resolve the mode structures. Stable modes can then be conveniently eliminated by varying the packing to determine if the displacements become singular or not with increased packing. This procedure seemed to be sufficient, though it requires considerable familiarity with the code.

The marginal wall position for $n = 1$ modes for this particular equilibrium is at $R_w = 1.26$. Higher resolution is unlikely to greatly modify this conclusion by very much. The mode structure is shown in Fig. 2.13-9 for $R_w = \infty$. All indications are that the modes are typical of high- $\bar{\beta}$ driven kink instabilities with a significant internal kink component inside the $q = 2$ surface. The internal kink component is a result of the high elongation; elongated ellipses are well known to be susceptible to destabilization by finite $\bar{\beta}$, and, for this reason, wall stabilization is relatively ineffective. An attempt was made to scale the equilibrium at constant $\bar{\beta}_p$ to decrease $\bar{\beta}$ to 5.7% and to 4.7% by scaling q_0 from 1.05 to 1.10 and 1.20. This was largely unsuccessful since increasing q_0 also further destabilizes the internal kink, even at low $\bar{\beta}$. At 4.7%, the growth rate with a wall at infinity drops by an order of magnitude, but is still unstable. To find the true limiting $\bar{\beta}$, a full optimization over various current profiles is necessary.

In conclusion, the design equilibrium studied requires a continuous wall at 1.26 times the plasma radius to stabilize $n = 1$ high- $\bar{\beta}$ driven kink modes. The alternative is to lower the $\bar{\beta}$ value at constant $q_0 \sim 1.05$, but axisymmetric stability will still require the equivalent of a conducting wall at 1.58 times the plasma minor radius.

GATO n=1 Stability
Growth Rate vs Wall Position

Figure 2.13-8. Growth rate versus wall position $n = 1$.

Figure 2.13-9. $n = 1$ modal shape (wall at ∞).

2.13.3 PASSIVE STABILITY

High plasma elongation and high plasma triangularity beneficial for high-beta operation greatly decreases the vertical stability of the plasma. TIBER II, with an elongation of 2.4 (95% flux), is vertically unstable. Without passive stabilization the plasma would displace vertically on the MHD time scale (sub-millisecond). Passive conductors are required to increase this time constant to a range where active control of the plasma is possible. Present experiments typically rely on a toroidally continuous, tightly fitting vacuum vessel to provide this passive stability. The external vacuum vessel and other toroidally continuous components are located far from the plasma in the TIBER II design, and other means must be used to provide vertical stability. The following section describes several options available for passively stabilizing TIBER II. The final design adopted relies on several plates located in optimal locations around the perimeter of the plasma. Toroidally continuous conductors are recommended; however, up/down anti-symmetric conductors forming saddle-currents on the face of each sector are also capable of providing stabilization, albeit with an increase in growth rate.

A number of assumptions and simplifications are used in this and in the following section on active stability. The EFIT equilibrium described in the previous section is used to describe the plasma shape and current distribution. The differences between this and the baseline equilibria should not greatly influence the results. The plasma is assumed to be a massless, rigid body moving in a set of resistive, toroidally continuous conductors. The plasma shape and current distribution are simulated by a large number of current elements, each of which must have the same vertical velocity and position vector (i.e., rigid body motion). This is a good approximation for the $n = 0$ modal shape, as can be seen in the MHD ideal stability analysis. Other assumptions include: axisymmetric system, no induced plasma currents, and negligible stabilization from the external equilibrium field coils. Stabilizing saddle currents within structural components are neglected. Small displacement about the neutral position is assumed; linear perturbation theory is used. The problem is formulated in the form of an eigenvalue problem and solved using the PSTAB computer code.

Figure 2.13-10 shows the vacuum field flux contour for the TIBER II equilibria (EFIT version). The convex nature of the flux contours owing to the quadrupole field required for high elongation is evident. The configuration is

TIBER-II VERTICAL STABILITY ANALYSIS

Flux (Vs)
Equilibrium Field Only, No Plasma Contribution

Figure 2.13-10. Vacuum field flux plot.

one of neutral stability; a small vertical displacement of the plasma will cause it to accelerate in the same direction. The destabilizing force gradient for this configuration is:

$$F'_d = \frac{dF_d}{dz_p} = 55.6 \times 10^6 \text{ N/m} .$$

A vertical displacement of 1 cm (0.01 m) produces a vertical force of 0.556 MN on the plasma from the equilibrium field coils. The passive stability system must be capable of providing a comparable but opposite force. The plasma motion induces current in the passive conductors which oppose plasma motion. These currents decay on the L/R time scale of the passive system and, accordingly, the active system must provide stabilization on this time scale.

2.13.3.1 Stability Parameter

The stability parameter (f) for the passive system is the ratio of the stabilizing force gradient to the destabilizing force gradient acting on the system.

$$f = - \frac{F'_s}{F'_d} = - \frac{\dot{F}_s}{\dot{F}_d} ,$$

where F_s and F_d are the stabilizing and destabilizing forces acting on the plasma, respectively. The ($\dot{}$) is differentiation with respect to time, and ($'$) is differentiation with respect to the vertical coordinate (z). The destabilizing force is caused by EF coils, and the stabilizing force is caused by eddy currents induced in the passive conductors. The stability parameter can be expressed as a combination of the plasma and conductor system mutual-inductance matrix as:

$$f = \frac{[I_p]^T [M'_{p,s}] [M_{s,s}]^{-1} [M'_{s,p}] [I_p]}{[I_p]^T [M'_{p,ef}] [I_{ef}]}$$

where: (I)_s are current vectors and (M)_s are inductance matrices. The subscripts refer to: (p), plasma; (s), stabilization conductors; and (ef), equilibrium field coils.

The stability parameter separates plasma motion into two distinct time-scale regions. If the stability parameter is less than one, ($f < 1$), plasma mass is important, and inertia terms dominate plasma displacement growth rate. Resistance of the passive elements is negligible, and the conductors act like superconducting elements on the time-scale of plasma motion (see Section 2.13.2). Plasma movement would occur on an Alfvén time scale, and the active system would be required to respond on the sub-millisecond time-scale. This leads to unacceptable requirements on the active control power supply.

For systems with stability parameter greater than one, ($f > 1$), the displacement growth rate is dominated by the L/R of the passive system, and the plasma inertia term can be neglected (i.e., massless plasma). The passive system time-constant is primarily a function of the conductor location and resistance and, accordingly, is controllable within certain limits. The growth rate can be expressed in terms of the coil time-constant and the stability parameter for a single-coil, single-conductor system:

$$\gamma = \frac{1}{\tau_{\text{passive}}} \frac{1}{(f - 1)}, \quad (f > 1),$$

where: γ is the plasma growth rate, and τ_{passive} is the L/R time constant of the passive system. The growth rate becomes infinite as the stability parameter approaches zero. In the more general eigenvalue problem defined below, there exists a single positive eigenvalue that approaches infinity as f approaches one. Essentially, the massless plasma solution is not valid in this limit, and the inertia terms must be included in the analysis; the ($f = 1$) essentially represents the MHD, $n = 0$ stability limit. For practical designs, the stability parameter must be greater than one for the entire plasma discharge. Since the stability parameter is critically dependent on plasma elongation, it is important to have sufficient margin above this limit to ensure stability under all discharge conditions.

The optimum conductor configuration was selected using the stability parameter as a performance parameter. The displacement growth rate is another parameter that can be used as a performance parameter; however, the stability parameter defines a limit below which plasma stability cannot be assured, while growth rate, as determined from the passive system L/R characteristics,

affects only the active system power requirements; i.e., without $f > 1$ for all plasma configurations through which the discharge will pass, it will become vertically unstable.

Figure 2.13-11 shows a contour plot of the stability parameter for a passive system consisting of a coil pair symmetrically located above and below the midplane. The coils are moved over the r - z plane, and the stability parameter associated with the coil location is plotted as the contour value. The coil size is a 10 x 10 cm square and is shown in the upper right-hand corner of the figure. The separatrix, 95% flux surface, and a line representing the first wall boundary are shown in the figure.

The contour plots essentially define the optimum location for passive conductors. Increasing contour values represents a better passive conductor location. As would be expected, the optimum location is inside the plasma boundary--and, if this were possible, the design would be passively stable. At this location the currents induced by vertical motion of the plasma have the largest restoring force on the plasma. The optimum location on the first wall boundary is at the $r = 3.7$ m, $z = 1.3$ m location. This optimum position for passive conductors is typical of highly elongated plasmas. Variations in plasma shape would only slightly change the contour shape, although the magnitude of the contour levels is highly dependent on shape and elongation. This location is also optimal for placement of the active coils.

Figure 2.13-12 shows the effects of placing a coil set at this optimum conductor location and moving a second conductor set around the r - z plane. The stationary coil is located approximately on the $f = 0.75$ contour of Fig. 2.13-11 and, accordingly, the minimum stability parameter for the second coil located far from the plasma is 0.75. The mutual inductance between the two systems requires that the stability parameter be less than the sum of the contributions of each separate system. Placing the coils close together is not the optimum solution. This figure shows that a stability parameter approaching one can be achieved if the additional conductors are located above and below the optimum position.

Figure 2.13-13 shows the stability parameter for a system containing a solid plate, represented by a number of ring conductors, located on the outer first wall. The plate thickness is 2 cm. As with the other figures, a pair of movable conductors are used to map out the optimum location of additional conductors. The stability parameter for this configuration (minimum contour

TIBER-II VERTICAL STABILITY ANALYSIS

f-Stability Parameter = $F_{\text{stable}}/F_{\text{destable}}$ Single Passive Coil Set $\Delta r = \Delta z = 1\text{m}$

Figure 2.13-11. Passive stability parameter for single coil pair.

TIBER-II VERTICAL STABILITY ANALYSIS

 f -Stability Parameter = $F_{\text{stable}}/F_{\text{destable}}$ Two Passive Coil Sets: 1st fixed @ $r=3.68$ $z=1.34$ m

Figure 2.13-12. Passive stability parameter for two coil pairs.

TIBER-II: PSTAB With Outer Passive Wall
f Stability Parameter

Figure 2.13-13. Passive stability parameter for Design 1.

level in the figure) is 1.16, or the stability parameter is 16% above the MHD stability limit. The contours indicate that the inner wall position around the $z = 1.5$ region is the next best location for conductors. Figure 2.13-14 shows a design with plates located in this region. The stability parameter increases to 1.35 for this design.

Little can be gained by adding conductors outside the first wall boundary. From a stability point of view, the divertor plates are located in an optimal location. Figure 2.13-15 shows the contours associated with adding conductors at this location. The minimum stability parameter is 1.51. The back surface of the divertor plate would have to be continuous in the toroidal direction. In this design, the divertor plate structure would have to be designed to carry the large electromagnetic loads associated with large plasma displacements.

Toroidally continuous plates impact the inductive start-up of the plasma and are difficult to implement in a toroidally segmented design like TIBER II. Continuous toroidal currents can be achieved by toroidal segmented rings which are connected top-to-bottom with vertical conductors or legs. Figure 2.13-16 shows a method for connecting the plates in an up/down anti-symmetric method. This construction method has the advantage of allowing for toroidal segmentation and produces an up/down anti-symmetric conductor system which is decoupled from the equilibrium field system. Flux swing of the equilibrium field coils to induce plasma current will not induce currents in these passive conductors. The conductor currents would be saddle currents flowing in the top and bottom rings in opposite directions. The up/down currents in alternate legs cancel, and only toroidal current would influence plasma position. The added resistance in the system associated with the vertical legs would increase the growth rate but has no effect on the stability parameter. The net result would be an increase in the active stability power requirements.

A multi-turn coil could be used instead of the plate and ring concepts described above. Figure 2.13-17 shows an example of a configuration utilizing a single large coil with many turns. The stability parameter for this case is 1.3. This design requires substantially more material than a plate design for a similar margin-to-MHD stability. The proximity to the plasma would make design and maintenance of this coil difficult. Additionally, the segmented

TIBER-II: PSTAB--Outer & Inner Passive Wall

$$t\text{-Stability Parameter} = F_{\text{stable}}/F_{\text{destable}}$$

Fixed Wall (Z= 2-2.0, 1.5-1.9) Movable= 10X10cm

Figure 2.13-14. Passive stability parameter for Design 2.

TIBER-II: PSTAB--Outer, Inner & Divertor Wall
f-Stability Parameter = $F_{\text{stable}}/F_{\text{destable}}$
Fixed Conductors: (R=9-2.5, 1.5-1.9) Movable = 10X10cm

Figure 2.13-15. Passive stability parameter for Design 3.

Figure 2.13-16. Toroidally segmented ring concept.

TIBER-II: PSTAB-- Single Large Passive Coil

 f -Stability Parameter = $F_{\text{stable}}/F_{\text{destable}}$ Growth Rate = 1.3 s^{-1} , Movable = 10×10

Figure 2.13-17. Passive stability parameter for Design 4.

nature of the TIBER II design would require up/down connection of the coils. As pointed out in the next section, this design achieves the lowest growth rate of any of the passive designs.

2.13.3.2 Growth Rate

For a multifilament, multi-conductor system, the growth rate is represented by the solution to an eigenvalue problem with the circuit equations, constrained by the force-balance equation, as the characteristic equation. The circuit equation for the passive system is represented as

$$[M_{s,s}] [\dot{I}_s] + [R_{s,s}] [I_s] = 0 \quad .$$

The force balance equation is

$$m_p v_p = 0 = F_d' z_p + [I_p]^T [M_{p,s}] [I_s] \quad ,$$

where m_p , z_p , and v_p are plasma mass, position, and velocity, respectively. $[M_{s,s}]$ and $[R_{s,s}]$ are the coil inductance and resistance matrices, respectively. The inertia term $m_p v_p$ has been set to zero in accordance with the assumptions of a massless plasma. The equations are solved using a generalized eigenvalue method (PSTAB). The eigenvalues represent the growth rates. The largest positive eigenvalue and the resulting eigenvector represent the plasma rigid-body motion and the modal current vector, respectively.

Table 2.13-3 summarizes the parameters associated with the previous designs. The growth rates and time constants reflect the use of copper ($2.0 \times 10^{-8} \Omega m$) as the plate material. The growth rate (inverse time constant) is directly proportional to the material resistivity. An order-of-magnitude increase in resistivity would proportionally increase the active power supply requirements.

2.13.4 ACTIVE STABILITY

In the resistive time scale the plasma growth rate must be stabilized with an active control system. As a minimum, a single anti-symmetric pair of actively controlled coils is required. Power and control considerations require that these coils should be located as close to the plasma as possible. The optimum passive conductor location is also the optimum active coil

Table 2.13-3. Passive conductor design summary.

Figure	Design	Configuration	Stability parameter (f)	Growth rate $\gamma(s^{-1})$	Growth rate time constant $\tau(ms)$
2.13-13	Design 1	Outer wall	1.16	13.71	73
2.13-14	Design 2	Outer and inner wall	1.35	6.41	160
2.13-15	Design 3	Divertor, outer, inner wall	1.52	2.93	340
2.13-17	Design 4	Single large coil	1.30	1.33	750

location. Figure 2.13-18 shows the geometry chosen for the active coil analysis, and passive Design 2 was chosen because of its simplicity. A pair of active coils was located symmetrically about midplane. A 20 x 20 cm coil was used in the analysis, although the final size will depend on structural and thermal considerations. Adding these coils as passive elements to the Design 2 system decreased the growth rate to $5 s^{-1}$. This new system is designated "Design 2a."

Similar analysis and assumptions used in the previous section are used in this section. For the active system, a voltage feedback term is added to the circuit equation, and an eigenvalue problem similar to that solved for the passive system provides the characteristics. The simplest assumption is to provide linear voltage gains on both plasma position and velocity. The circuit equation for this system is:

$$[M_{s,s}] \dot{[I_s]} + [R_{s,s}] [I_s] = - [G_z] z_p - [G_v] v_p ,$$

where $[G_z]$ and $[G_v]$ are the position- and velocity-gain vectors. In the present problem, only two components of the gain vectors are non-zero, and these are the gains which are applied to the active coils in the positive and negative plane. The gain terms allow energy to enter the system from an external source and mathematically provide a means for driving the growth rate negative. The system is actively stable when all real eigenvalues are

TIBER-II: Design 2a
Active Coil Geometry

Figure 2.13-18. Passive/active stabilization geometry, Design 2a.

negative; that is, any perturbation in the system decays in time. The solution produces real and complex eigenvalues, the latter representing oscillatory behavior in the plasma position.

The initial current vector of the active system, and the vector which determines the contributions of each eigenvector to the particular solution, is the largest positive eigenvector of the passive system. This essentially means the fast-time-scale perturbation associated with plasma motion is reacted by the passive system. The active system responds to the perturbed passive system on the time-scale of the combined passive/active system. This is similar to stating that the active system has a certain time delay, and during this delay the currents are set up in the system based on the passive system response. Following the time lag, the active system develops voltages that drive currents in the active coils to center the plasma. The solution, while complex in nature, is fully deterministic. The displacement, current, and power-time histories can be found. The solution is carried out using the active stability code ASTAB.

Figure 2.13-19 shows the plasma response to a 1-cm vertical plasma offset. The exponential growth of the passive system (i.e., $G_V = G_Z = 0$) is shown for reference. The remaining curves all have a position-gain factor (G_Z) of 1000 V/m (1 turn volts/centimeter). At this value of position gain, the largest growth rate is always negative; that is, theoretically the system is actively stable. However, without velocity gain the solution is oscillatory and decays at a very slow rate. The three other lines in the figure represent systems with increasing velocity gain. In the classical sense, these three lines represent systems that are: $G_V > 80$, under-damped; $G_V = 80$, critically-damped; and $G_V < 80$, over-damped. The critically damped solution represents the minimum time to bring the plasma back to the centered position without overshoot and is typically the optimum control point for the power system.

Figure 2.13-20 shows contours of the plasma growth rate as functions of the system-gain parameters (G_Z and G_V). The point solutions are shown in the previous figure. Negative values indicate exponential decay of any perturbation, a necessary condition for active stability. The elbow in the curves represents critically damped systems. This represents approximately the optimum design point of the active feedback system since it stabilizes the perturbation in the least amount of time.

Figure 2.13-19. Plasma displacement time history, $G_z = 1000 \text{ V/m}$, $Z_{p0} = 1 \text{ cm}$.

**ASTAB6: --TIBER-II Design 2a--
Maximum Oscillatory Growth Rate (s^{-1})

Figure 2.13-20. Growth rate contours in gain phase space.

The current in the upper active coil is shown in Fig. 2.13-21. Equal and opposite current flows in the lower active coil, and the two active coils theoretically could be connected in series to a single power supply. At $t = 0$, the currents are flowing in all elements of the passive and active system based on the distribution associated with a rapid vertical shift in the plasma. The gain terms in the equation set apply the appropriate voltage to the active coils at $t = 0$, and the active coil currents respond accordingly. The negative initial currents represent opposition to the vertical motion of the positive plasma current. The active coil drives the plasma back to the neutral $z = 0$ position and also induces currents in the passive system, which counteracts this motion. The peak current requirements are proportional to plasma displacement. The 40 kA associated with a 1-cm displacement would become 400 kA for a 10-cm displacement.

The power demand for these systems is shown in Fig. 2.13-22. The passive system power demand is, of course, zero. The peak power demand decreases with increasing velocity gain; without velocity gain the peak power is 1 MW. The critically damped solution reduces this to 0.4 MW. Additional power savings can be achieved with additional velocity gain (or velocity damping) at the expense of damping time. The power demand is a strong function of the passive system time-constant, plasma excursion, and active coil position. The powers in this figure are close to the minimum values achievable. The power demand depends on the passive system time-constant. Decreasing the resistivity of the system would decrease the time-constant and allow the system to damp more rapidly. This would also increase the power demand proportionally, as it depends on the square of the plasma displacement. Excursions as large as 10 cm are possible in the present design, and, accordingly, the power demands would increase from 0.4 MW to 40 MW for a 10-cm plasma excursion. Nonlinear effects could also become important at such large excursions. The proposed location of the active coil set is reasonably close to the plasma. Power demands are a strong function of coil/plasma proximity. Vertical control from coils located far from the plasma would dramatically increase power demands.

Other areas which were not included in this study but which impact the active control problem are the sensing and feedback control time-lag problem and nonlinear effects, such as plasma shape-change and non-ideal equilibrium displacement. It was assumed that the plasma position was prior knowledge and that no feedback time lag was introduced by the power supply. Additionally,

Figure 2.13-21. Active coil current time history, $G_z = 1000$ V/m, $ZP_0 = 1$ cm.

Figure 2.13-22. Active coil power time history, $G_z = 1000$ V/m, $ZP_0 = 1$ cm.

the power supply was assumed ideal in its power/frequency spectra. Problems associated with plasma position-sensing can make certain gain-control schemes unstable. Although these effects should be second-order relative to the linear rigid-body results of this study, a more detailed analysis is needed to quantify their effects on system design. The net result of these phenomena will be to increase the peak power demand on the active system. Substantial margin should be placed on the current, voltage, and powers defined in this study.

2.13.5 SUMMARY AND RECOMMENDATIONS

The system concept developed for the vertical stabilization of the TIBER II design utilizes both active and passive elements to control the vertical displacements within tolerable limits. A series of passive plates carrying toroidal current in continuous or toroidally segmented up/down anti-symmetric configuration is capable of providing 35% margin on MHD $n = 0$ stability. With the plates made of copper the passive system time-constant is 160 ms; other materials and/or a segmented design would have lower time constants. The active system is composed of two actively controlled coils located close to the plasma. Displacement and velocity gain are both required to achieve active stabilization of the plasma. Power levels as low as 0.5 MW are sufficient to stabilize anomalous vertical displacements of 1 cm. Larger displacements, a more resistive passive system, or movement of the active coils away from the plasma would dramatically increase this power demand.

Sufficient margin can be included in the design to ensure the system is stable for the different equilibria that may be encountered during the life of the machine, although no attempt was made to study system performance at these off-design conditions. Further optimization of the system is possible. The system is simple and can be included in the present design with a minimum of additional hardware. The study established the feasibility of the concept for implementation in the TIBER II design; however, significant detailed work remains to be done for full integration into the present design. At least conceptually, the implementation seems tractable.

A number of additional research areas remain to be studied before stability of the next generation, highly elongated tokamaks, can be assured. In particular, the need for close plasma/coil proximity and the control issues associated with detecting and correcting plasma position with active feedback

are open questions in the fusion community. The impact of real passive and active control surfaces on ideal MHD phenomena and β limits must be explored for these highly elongated shapes. As was pointed out in previous sections, the ideal MHD stability for higher modes and its role in limiting accessibility to certain plasma shapes must be further explored. Additionally, a number of non-linear phenomena influence plasma stability and control. Transport and edge control issues associated with highly elongated divertor plasmas influence plasma positional control and must be integrated into any stability analysis. Internal MHD phenomena--such as sawteeth, ELMs (Edge-Localized-Modes [4]), and minor plasma disruptions--also influence stability system issues. Finally, stability issues are intimately involved in determining machine performance and operating regimes. The ignition margin is ultimately constrained by plasma stability issues. For the next machine, the stability and control issues should be addressed synergistically with other machine design issues early in the design cycle.

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2.14 IMPURITY FLOW REVERSAL

2.14.1 INTRODUCTION

The purpose of this section is to consider neutral beam (NB) driven flow reversal and electron cyclotron resonance heating (ECRH) driven flow reversal as possible mechanisms for impurity removal in TIBER II.

Stacey and Sigmar [1,2] have shown that strong neutral beam co-injection can produce an outward flow of impurities in tokamak plasmas, while counter injection results in an enhanced inward flow. The Stacey-Sigmar theory is self-consistent and its predictions are supported qualitatively by a considerable experimental data base [3-5]. Detailed analysis of the experimental data from PLT [6] with tungsten limiters has shown that the theory can quantitatively explain the experimental data quite well.

Stacey has also shown that the poloidal variation in the electrostatic potential caused by ECRH [7] can produce an outward contribution to the impurity transport flux and even reverse the usual inward flow [8].

In the following sections, after a brief review of the relevant theories, both formalisms are applied to analyse the transport of helium ash and edge-originated impurities (carbon) in TIBER II.

2.14.2 NEUTRAL BEAM DRIVEN FLOW REVERSAL

The parallel and normal components of the momentum balance equation and the particle continuity and charge neutrality equations can be combined on a fast time scale to write the radial flux-surface averaged impurity flux as a sum of the contributions arising from different effects, that is:

$$\begin{aligned} \langle n_I v_{Ir} \rangle = & \langle n_I v_{Ir} \rangle_{PS} + \langle n_I v_{Ir} \rangle_{NC} + \langle n_I v_{Ir} \rangle_M + \\ & \langle n_I v_{Ir} \rangle_I + \langle n_I v_{Ir} \rangle_{\phi'} + \langle n_I v_{Ir} \rangle_{\phi}^{\sim} \end{aligned} \quad (2.14-1)$$

The explicit form of each term in Eq. (2.14-1) is given in Ref. [2].

The first two terms correspond respectively to the Pfirsch-Schluter and neoclassical fluxes of the usual transport theory, but now modified to account for the radial transfer of toroidal momentum and for the variation of impurity and ion densities over the flux surface. The third term is the transport flux resulting directly from the interaction of the beam ions with the plasma and

impurities. The fourth term results from the retention of the inertial term $nm(\vec{v} \cdot \nabla)\vec{v}$ in the momentum balance equation. The fifth term is the radial impurity flux driven by the linear component of the electric field, while the last term results from the poloidal variation of the electrostatic potential over the flux surface.

For strong co-injection, the inertial and radial electric field flux components are outward, while the neoclassical and PS terms are proportional to the difference in main ion and impurity pressure gradients and can give inward flux contributions. As a result of this competition between the different terms in Eq. (2.14-1), flow reversal can occur.

An important element in the impurity flow reversal theory is the radial transfer of the beam momentum across a flux surface. This was represented in the original theory by a "drag" term. The more recent Stacey-Sigmar gyroviscous theory [9] for momentum confinement has been shown [10] to be capable of accounting for the experimentally inferred radial momentum transport rates and to lead to a representation of the viscous force, which is consistent with the drag term form that was used in the impurity flow reversal theory [1,2].

According to gyroviscous theory [9], the drag frequency is given to within a constant of order unity by:

$$\nu_d = \frac{T_i}{2R^2ZeB} \quad (2.14-2)$$

where R is the major radius, Z the effective ionic charge, e the electronic charge, B the magnetic field, and T_i the ion temperature.

2.14.2.1 Application to Helium Transport in TIBER II

The removal of helium ash (thermalized alpha particles) is an important problem in the design of tokamak reactors. Accumulation of thermal alphas at the center can interfere with the burning process and degrade the performance of the reactor [11].

The Stacey-Sigmar theory is now used to study the transport of helium in TIBER II. The radial density and temperature profiles for the ions and electrons were assumed to be of the form

$$x(r) = x_0 [1 - (r/a)^2]^{\alpha_x} \quad (2.14-3)$$

where x represents either density or temperature and x_0 is the value at the center. The exponents α_n and α_T for the main ion density and temperature, respectively, were taken to be 1.02 and 0.58. The rest of the Tiber II parameters are summarized in Table 2.14-1.

The fast ion deposition profile needed for the calculation of the neutral beam momentum deposition was computed with the code PROCTR [12]. The beam parameters and the beamline geometry are the same as in the NB current drive study. The helium profile was also assumed to be of the form given in Eq. (2.14-3). Since the exact spatial distribution of the thermalized alpha particles is not known, three different values of the exponent α_n were considered: 1.0, 1.5, and 2.0. The central value of the helium density was calculated in such a way that the volume-averaged helium density was 5% of the volume-averaged ion density.

The results of the numerical simulation are shown in Figs. 2.14-1(b) - 2.14-1(d). In each figure, the flux-surface-averaged helium flux is plotted versus the minor radius for different beam powers for the three different helium profiles.

The results are sensitive to the assumed helium profile: flow reversal occurs at lower beam power when the profile is more peaked. This can be explained if we note that the diffusive impurity fluxes in Eq. (2.14-1)--which are driven by the main ion and impurity ion density gradients and are normally inward--for low-Z impurities such as helium and peaked impurity profiles relative to the ion profile can be outward. The same behavior was observed in the case of strong profile coupling where the thermal alpha profile was proportional to the fusion reaction rate $n_i^2 \langle \sigma_f v \rangle$, in agreement with similar observations made by Sigmar et al. [13].

If the helium profile was the same as the main ion profile (Fig. 2.14-1(b)), one would need 75 MW or higher of co-injected beam power to drive the impurities out of the center. For more peaked profiles, which are probably more realistic, 25 to 50 MW would be adequate to drive the helium out of the center (Figs. 2.14-1(c) and 2.14-1(d)).

Figure 2.14-1. Neutral beam deposition profile (a), and beam driven flow reversal (b-d) for different helium profiles in TIBER II with 500 keV beam energy.

Table 2.14-1 TIBER II parameters.

R_0	m	3.0
a	m	0.834
B_0	T	5.55
$\langle T_e \rangle$	keV	19.2
$\langle T_i \rangle$	keV	19.58
$\langle n_e \rangle$	10^{20} m^{-3}	1.060
$\langle n_{DT} \rangle$	10^{20} m^{-3}	0.796
Z_{eff}		2.085
$q(0)$		1.268
$q(a)$		3.945
E_b	keV	500

2.14.2.2 Application to Carbon Transport in TIBER

Here, the NB flow reversal theory was used to analyze the transport of carbon in TIBER II. The carbon profile was the same as the main ion profile, and the carbon concentration, $\alpha = Z^2 n_z / n_i$, was assumed to be 5%. The results are shown in Figs. 2.14-2(a) - 2.14-2(b). The flux-surface-averaged carbon flux is plotted versus the minor radius for different beam powers (Fig. 2.14-2[a]). It can be seen that 50 MW of co-injected beam power would be adequate to drive the carbon impurities out of the center.

To study the sensitivity of the calculations to the beam deposition profile, the tangency radius of the outer beamline was changed from 3.0 m to 3.2 m. This results in a less peaked deposition profile and hence in less momentum being deposited in the center. The results are shown in Fig. 2.14-2(b). Although the impurity fluxes are now smaller for the same beam power, flow reversal still occurs for co-injected beam power of the order of 50 MW.

2.14.3 ECRH DRIVEN FLOW REVERSAL

Cyclotron wave heating increases preferentially the perpendicular kinetic energy of the resonant particles, which leads to increased trapping of those particles at the outboard of the torus, resulting in a poloidal variation of the electrostatic potential of order $\epsilon = r/R$ [7,14].

Figure 2.14-2. Carbon flow reversal in TIBER II with reference profiles, at 500 keV neutral beam energy and beam tangency radius of (a) 3.0 m, (b) 3.2 m.

This poloidal variation in the potential will produce a poloidal variation in the main ion and impurity densities, which has been shown [2,15] to significantly affect impurity radial transport. In the case of ECRH, Stacey has shown [8] that the poloidal asymmetries can produce an outward contribution to the impurity transport flux, which could inhibit, possibly even reverse, the normal inward flow of impurities.

The radial flux-surface-averaged impurity flux is still given by Eq. (2.14-1) after omitting the beam-particle momentum exchange term. The last term in Eq. (2.14-1) is a flux resulting directly from the poloidal variation in the electrostatic potential and is given by

$$\langle n_I V_{rI} \rangle_{\phi}^{\sim} = \frac{n_I^0 m_I \nu_I^0 I_i}{e_I B_p^0} \epsilon^2 \beta_I \left(\frac{\phi^C}{\epsilon} \right) \left(\frac{E_r^0}{B_p^0} \right),$$

where ϕ^C is the cosine component of the poloidal variation of the electrostatic potential, and the rest of the variables are defined in Ref. [8].

2.14.3.1 Application to Helium Transport in TIBER II

The theory was applied to helium ash transport in TIBER II. The parameters and helium profiles of Section 2.14.2.1 were used. The results are shown in Fig. 2.14.3(a)-(c). The flux-surface-averaged helium flux is plotted versus ϕ^C/ϵ for different radial locations for the three different helium profiles.

Here, too, the results are sensitive to the helium profile: for the same radial location, flow reversal occurs for lower values of $|\phi^C/\epsilon|$ for more peaked profiles. This can be explained by using arguments similar to those in Section 2.12.2.1. In the case of the broader profile with exponent $\alpha_n = 1.0$, a value of $|\phi^C/\epsilon|$ of order 2 would be needed to drive the helium out of the center (Fig. 2.14.3[a]). For the more peaked profiles in Fig. 2.12.3(b)-(c), a value of $|\phi^C/\epsilon|$ of order 1 would be adequate.

Figure 2.14-3. ECRH drive helium flow reversal for different helium profiles in TIBER II.

2.14.3.2 Application to Carbon Transport in TIBER II

Here the theory is applied to carbon transport in TIBER II. The carbon concentration, $\alpha = n_Z Z^2 / n_i$, was assumed to be 5%. The results are shown in Fig. 2.14-4. The flux-surface-averaged carbon flux is plotted versus Φ^C/ϵ for different radial locations. It can be seen that a value of $|\Phi^C/\epsilon|$ of order 2 would be needed to drive the carbon impurities out of the center.

2.14.4 SUMMARY AND DISCUSSION

The neutral beam injection flow reversal theory and the ECRH flow reversal theory were applied to model the transport of helium ash and carbon in TIBER II.

The results of the neutral beam injection flow reversal theory were sensitive to the assumed helium profile, but it was found that for the more realistic profiles, 25 to 50 MW of co-injected beam power would be adequate to drive helium out of the center. In the case of carbon, it was found that 50 MW of co-injected beam power would be needed to remove the carbon impurities from the center. The theory has a well developed theoretical basis, and its predictions are supported by experimental results [3-6]. Thus we tentatively conclude that about 50 MW of co-injected 500-keV NB power would drive helium and carbon impurities out of the central part of the TIBER plasma.

The ECRH-flow-reversal theory also predicts an outward flow of helium for values of the poloidal variation of the potential $|\Phi^C/\epsilon|$ of order 1 and an outward flow of carbon for $|\Phi^C/\epsilon|$ of order 2. While preliminary calculations indicate that such values of $|\Phi^C/\epsilon|$ are obtainable [7,14], more work is needed in this direction. There is no experimental evidence yet for ECRH-driven flow reversal, and the theory, in its present form, is strictly valid only for collisional impurities; hence the results could be regarded as tentative and indicative of order of magnitude.

In conclusion, neutral beam injection flow reversal appears to be a feasible mechanism for impurity removal in TIBER II. ECRH flow reversal is a promising candidate, but additional research is needed to determine the power levels required to produce the desired poloidal variations in the electrostatic potential. For both methods, time-dependent impurity transport calculations must be carried out to provide a more accurate estimate of power requirements for impurity control.

Figure 2.14-4. ECRh driven carbon flow reversal in TIBER II.

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APPENDIX 2.A

A BRIEF DESCRIPTION OF MUMAK: A COMPUTER CODE FOR MODELING PLASMA POWER BALANCE AND CURRENT DRIVE IN TOKAMAKS

MUMAK [1] is a zero-dimensional (volume-averaged), time-independent power balance code for modeling plasma particle and energy flow in tokamaks. The equations describing the steady-state power balance include terms for alpha power from D-T burning plasmas, inductive (ohmic) as well as non-inductive current drive power, conduction power losses described by empirical energy confinement scaling laws, and radiation power losses. Radial profiles of plasma density and temperature are controlled by input parameters and are assumed fixed for the power balance calculations.

MUMAK also performs calculations of non-inductive current drive efficiencies for electron cyclotron harmonic (ECH), neutral beam injection (NBI), lower hybrid (LH) slow wave, and fast wave (FW) current drive. The NBI and LH current drive models are described briefly in Sec. 2.2 of this report. The ECH models include relativistic effects and degradation of the efficiency due to electron trapping effects. The NBI models also include the effects of electron trapping as well as the beam energy dependence. The LH models take into account both the accessibility limit and the condition for avoiding strong electron Landau damping of the wave to determine the radius of maximum penetration for the waves.

Ancillary calculations performed by MUMAK include the determination of volumetric neutron source strength from D-T burning plasmas, thermal stability analysis of the steady-state burn operating point, and numerous graphical diagnostics. Radial plots of cyclotron and plasma frequencies, optical depth for ECH waves of various polarizations, and the volumetric neutron source strength can be generated. Results from one- and two-dimensional scans of parameter space can also be plotted.

MUMAK also contains several models that address issues unique to steady-state tokamak operation. The first is the particle balance of fusion alpha particles. This particle balance is computed at the solution (operating) point and the required alpha particle confinement time necessary to maintain the steady-state thermal alpha concentration (user-specified) is determined.

Tokamak designs with much longer burn pulses must specify some mechanism by which the alphas are removed or the accumulation of alpha ash will eventually choke off the burn.

Another unique feature of steady-state tokamak burn modeled by MUMAK is that the current drive power P_{CD} is comparable to the alpha power. For fixed toroidal current, P_{CD} is also sensitive to the plasma density and temperature at the operating point. The overall power balance is therefore quite sensitive to the current drive efficiency scaling and to the partitioning of the current drive power between bulk electrons and ions.

2.A.1 METHOD OF SOLUTION

The power balance equations in MUMAK are formulated as a set of N non-linear equations in N unknowns ($N = 3$ or 4). The non-linear equation solver COSNCF from the NAG mathematics library [2] is used to obtain a solution (a zero) to these equations to a specified accuracy. Routines are also available to search for multiple solutions to the equations. This is done by renormalizing the power balance equations after the first solution is found so that a new set of equations is generated for which the first solution is no longer acceptable. This renormalization is done in such a way that if a solution to the new set of renormalized equations is found it will satisfy the original power balance.

Power balance solutions are found by varying plasma temperatures, densities, energy confinement time, and operating beta. A two-fluid model is used with separate power balance equations for electrons and ions including electron-ion drag. The electron power balance equation takes the form

$$P_{ae} + P_{ie} + P_{OH} + P_{CD,e} + P_{aux,e} = P_{c,e} + P_{rad} \quad (2.A-1)$$

where

- P_{ae} = alpha power deposited in bulk electrons
- P_{ie} = ion-electron drag power
- P_{OH} = ohmic heating power
- $P_{CD,e}$ = non-inductive current drive power to bulk electrons
- $P_{aux,e}$ = auxiliary bulk electron heating power
- $P_{c,e}$ = bulk electron transport power
- P_{rad} = bulk electron radiation power.

The electron radiation term includes bremsstrahlung and synchrotron power losses. The ion power balance equation is

$$P_{\alpha i} - P_{ie} + P_{CD,i} + P_{aux,i} = P_{c,i} \quad (2.A-2)$$

where

- $P_{\alpha,i}$ = alpha power deposited in bulk ions
- $P_{CD,i}$ = non-inductive current drive power to bulk ions
- $P_{aux,i}$ = auxiliary bulk ion heating power
- $P_{c,i}$ = bulk ion transport power.

Detailed expressions for the terms in Eqs. (2.A-1) and (2.A-2) are given in Ref. [1].

The independent variable corresponding to the electron power balance equation is the density-weighted, volume-averaged electron temperature, $T_e = \langle n_e T_e \rangle / \langle n_e \rangle$ (angle brackets denote volume averaging). Temperature and density profiles are taken to be of the form

$$X(r) = X(0) [1 - (r^2/a^2)]^{\alpha_X} \quad (2.A-3)$$

where $X = n, T$. Volume averaging over these profiles is done using a circular flux surface model. For the ion power balance the ion temperature, $T_i = \langle n_i T_i \rangle / \langle n_i \rangle$ or the gross energy confinement time, τ_E , parametrized by an H-mode enhancement factor, $H_f = \tau_E / \tau_E(\text{L-mode})$, can be varied. The third equation is a gross pressure balance constraint that requires that the total plasma pressure normalized to magnetic field squared be equal to a user-specified operating beta. Typically this equation is used to determine the volume-averaged electron density, $n_e = \langle n_e(r) \rangle$. For solutions with user specified n_e , this equation can also be used to specify the operating beta. Finally, a fourth equation is available to constrain the solution to a user-specified value of average neutron wall loading, $\langle \Gamma \rangle$. In this case, either the operating beta or the energy confinement enhancement factor, H_f can be varied providing it is not being used as the independent variable for one of the other equations.

Typically, the user specifies the total toroidal current, the empirical scaling law for energy confinement, and the type of current drive to be used. From the current and the plasma shape MUMAK derives the beta limit either from the Troyon [3] expression or the Yamazaki [4] formulas, and the density limit from the Murakami [5] or Greenwald [6] expressions. The user also specifies either the H-mode energy confinement enhancement factor, H_f or the fraction $f_b = \beta(\text{operating})/\beta(\text{critical})$ (or both if $\langle r \rangle$ is allowed to vary). Scenarios without current drive power, in which the plasma current is assumed frozen (long L/R decay time) can also be modeled. For some special cases, the current drive power may be input and the models will calculate the resulting toroidal current.

MUMAK can also perform Plasma Operating CONtour (POPCON) analysis. The user specifies NEQ = 2 and sets the density and temperature to the desired values. The auxiliary power required to satisfy power balance for the electrons and for the ions is calculated separately. The sum of the electron and ion powers can easily be plotted for each density and temperature in a density versus temperature space to generate a POPCON plot.

2.A.2 FUTURE UPGRADES

Though MUMAK contains all of the basic components for a complete zero-dimensional description of steady-state tokamak power balance, several improvements can still be made. The model that determines the core current fraction, in cases using multiple current drive schemes simultaneously, will be upgraded so that the current density profile data can be read in from external files. This will allow results from 2-D MHD calculations to be read by MUMAK without changing the source and recompiling. Results from Fokker-Planck and transport calculations that predict the bootstrap current density profile will be accessible by MUMAK in the same way. This should increase the accuracy of the partitioning of toroidal current into bootstrap and multiple current drive components calculated by MUMAK.

The NB current drive models will be improved to include ion trapping effects and the degradation of the efficiency due to toroidal plasma rotation. Preliminary estimates indicate these two effects may decrease the A/W conversion efficiency for NBI current drive by 10 to 30%. For NBI calculations, improvements are also required in the hot ion model. Particular attention will be given to improving the calculation of the average hot ion

energy. This may have an appreciable effect on the hot ion fraction of the operating beta. This is an important calculation since the performance of many ETR class tokamaks may be limited by achievable beta more than other factors.

The deleterious effects of beam ion instabilities on the NBI current drive efficiency must be modeled in MUMAK. For beams with particle velocities above the Alfvén speed in the plasma, instabilities may increase the beam ion slowing down [7], thereby significantly reducing the current drive efficiency. The beneficial effects of bulk plasma stabilization by beam ions must also be modeled. Preliminary theoretical estimates [8] indicate that the toroidal beta limit may be increased with beam injection.

Finally, the MUMAK power balance models may be expanded to treat the core and edge power balances separately. This upgrade would require models for energy transfer between regions, separate energy transport scaling for each region, and accurate partitioning of the current drive and auxiliary power to the regions in which the power is absorbed. This kind of calculation may be required to more accurately treat scenarios in which LH power is absorbed in the edge and heats electrons there while another form of current drive power is absorbed in the core.

2.A.3 DOCUMENTATION

A report [1] exists which is intended to serve both as a reference to the physics models incorporated in MUMAK and as a user's guide to running MUMAK. Since MUMAK is presently configured as a BASIS package [9-11], users will need to become somewhat familiar with BASIS to appreciate the full capability of the MUMAK package for parametric analysis. The user's guide in this document will give enough information on running MUMAK under BASIS to generate individual cases and simple parametrics. Readers interested only in using MUMAK as a "black-box" can skip Secs. 3-6 of Ref. [1] without loss of continuity.

Reference [1] is organized as follows: Section 2 presents an overview of the mathematical approach for solving the power balance equations, including a description of some unique features of steady-state tokamak operation that are modeled. The detailed power balance equations are described in Sec. 3. The steady-state current drive efficiency models are presented in Sec. 4. Several additional calculations performed by MUMAK such as the volumetric neutron

source calculation and the thermal stability analysis are described in Sec. 5. Sample results for the TIBER II configuration are given in Sec. 6. A self-contained user's guide is given in Sec. 7. A summary, including the present status of the MUMAK modeling effort and future plans, is presented in Sec. 8.

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APPENDIX 2.B.

LOWER HYBRID PARAMETERS FOR BASELINE CURRENT DRIVE SYSTEM

2.B.1 PARAMETER SELECTION

Physics codes can be used to model wave behavior and rf-current generation to determine the engineering features to be implemented in the rf hardware design. We describe in this section the choice of wave frequency, power spectrum, and launcher location needed for steady-state, high-density, lower hybrid current drive (LHCD), which is expected to be achievable in the outer regions of the high density plasma.

2.B.2 BENCHMARK COMPARISON OF CODES

Two codes have been used in portions of the LHCD design work. The first, a quasilinear damping code [1], most accurately models wave damping, whereas the second code [2] calculates linear Landau damping, but is capable of accurately computing a two-dimensional self-consistent MHD equilibrium from LHCD. The former code iterates the local distribution function on each flux surface while the latter code iterates the MHD equilibrium; both codes treat multiple ray spectra. Since many calculations reported here used the MHD code, it was desirable to first compare the results of both codes for a standard case, thus gauging the accuracy of the linear damping treatment.

The benchmark geometry for the comparison was an "equivalent circular" TIBER II with major radius $R_0 = 3.0$ m, minor radius $a = 1.24$ m, axis vacuum field $B_0 = 6.0$ T. The peak electron and ion temperatures were $T_{eo} = T_{io} = 30$ keV with a peak electron density $n_{eo} = 1.7 \times 10^{20} \text{ m}^{-3}$. The quasilinear code has an approximate (shifted circle) equilibrium with an effective minor radius coordinate r and specified profiles:

$$T_e = (T_{eo} - T_{eb})(1 - r^2/a^2) + T_{eb}$$

$$n_e = (n_{eo} - n_{eb})(1 - r^2/a^2)^{1.21} + n_{eb}.$$

The MHD code specifies profiles as functions of a normalized flux $\tilde{\psi}$

$$T_e = T_{eo} \tilde{r}^2$$

$$n_e = n_{eo} \tilde{r}^{1.2}$$

Figure 2.B-1 displays n_e and T_e in terms of the half-width of each flux surface for the quasilinear (Q) and linear-MHD (M) code inputs; the principal difference is that edge density and temperature are zero for the MHD code

Figure 2.B-2 compares ray trajectories in a poloidal projection for a mode with a toroidal mode number $n_t = 1.41$, frequency $f = 4.6$ GHz, and launched at the bottom. The paths are quite similar, although the Q model allows slightly deeper penetration. Damping and LHCD depend critically on the ratio $C/N_{\parallel} v_e$, where $C/N_{\parallel}$ is the wave's phase speed along a field line and v_e is the electron thermal speed. The value of $n_{\parallel}$ evolves along the ray path due to toroidal effects, and Fig. 2.B-3 shows that $n_{\parallel}$ varies along the path in a similar fashion for both codes

Both codes model wave damping by launching a collection of rays with different $N_{\parallel}$ values. The Q code specified a power distribution $P \propto \exp[-32 (N_{\parallel} - 1.5)^2]$, and this was simulated in the M code with an equivalent power distribution among six modes with $N_{\parallel} = 1.3, 1.4, 1.5, 1.6, 1.7$, and 1.8 . The flux surface average heating power density of electrons is compared in Fig. 2.B-4 with a total launched power of 45 MW. The results are strikingly similar although the Q code actually achieves penetration to higher T_e and n_e (due to the different profiles, as shown in Fig. 2.B-1). Peak heating occurs at 11.2 keV for Q and 9.0 keV for M; heating power vanishes at 17.5 keV for Q and 14.5 keV for M. Thus, the quasilinear damping allows heating and LHCD to slightly higher T_e values than a linear calculation, as expected from basic principles [3]. An approximate comparison of the LH current density profiles is given in Fig. 2.B-5; the results are not absolutely comparable as the Q model predicts a local value of $j_{\parallel}$ while the M code computes a flux surface averaged quantity, $G = \langle j_{\parallel} B \rangle$. Nevertheless, the LHCD predictions are similar for both codes.

This benchmark calculation was repeated for two other cases:

$T_{eo} = 20.0$ keV and $T_{eo} = 15.0$ keV. Both codes achieve deeper spatial penetration--peak power at $r = 91$ cm for Q and at $r = 83$ cm for M; again the Q code reaches a slightly higher T_e value.

Figure 2.B-1. Benchmark calculations for (a) quasilinear (Q) damping code and MHD (M) - linear damping code: density profile and (b) temperature profile.

Figure 2.B-2. Benchmark comparison: poloidal projection of ray trajectory with toroidal mode number 1.41; $f = 4.6$ GHz.

Figure 2.8-3. Benchmark comparison: evolution of $n_{||}$ along ray path. Abscissa is $n = (n_e - n_{\text{edge}})/(n_{e0} - n_{\text{edge}})$.

Figure 2.B-4. Benchmark comparison: electron heating power density.

Figure 2.B-5. Benchmark comparison: current density profile. Note that Q model achieves strong damping at somewhat larger electron temperature than M model.

In addition to comparing code results, these calculations serve to underscore a basic conclusion regarding LHCD. In all cases LHCD is only possible in the outermost plasma region when T_{eo} and n_{eo} are large (in the reactor range). Both codes contain additional features which model current density spatial diffusion, but at present this can not be relied upon to fill in current density near the magnetic axis.

2.B.3 RAY TRACING AND CURRENT DRIVE IN TIBER II GEOMETRY

A detailed study of LHCD was carried out with the MHD code. For these cases rays were launched into a plasma in MHD equilibrium with $R_0 = 3.0$ m, $A = 3.6$, $\kappa = 2.4$, and $d = 0.4$. This equilibrium, pictured in Fig. 2.B-6 and 2.B-7, has a pressure profile $p \propto n_{eo} T_{eo}^{1.4}$, total current $I_0 = 10.2$ MA, and $\beta = 0.0504$. The safety factor is $q_a = 1.0$ at the axis and $q_b = 3.35$ at the boundary. The plasma profiles corresponded to an earlier version of the baseline, with a flat density and peaked temperature:

$$T_e = 30 \text{ keV } (\propto r^{1.1})$$

$$n_e = 1.43 \times 10^{20} \text{ m}^{-3} (\propto r^{0.3}) .$$

The ions are a 50% DT mix, and $Z_{eff} = 1.5$ is assumed in the current drive calculation.

Three wave-design parameters are generally available in addition to the choice of polarization. These are frequency, $N_{\parallel}$ and θ , (the launch position). For a given plasma a three-parameter survey should be performed to optimize the LHCD system, but we are guided in choosing the bounds on the parameters. Thus, $f \lesssim 2.4$ GHz would risk ion power absorption, while $f \gtrsim 4.6$ GHz presents waveguide grill fabrication problems. For a given f , a minimum $N_{\parallel}$ is imposed by accessibility.

Figure 2.B-8 shows the ray path and $N_{\parallel}$ variation for $f = 4.6$ GHz, $N_{\parallel} = 1.5$, and $\theta = -0.25$ (bottom launch). The upshift of $N_{\parallel}$ permits low edge values of $N_{\parallel}$ and weak edge damping without violating accessibility. Similar calculations at $f = 2.47$ GHz, $N_{\parallel} = 1.5$, and $f = 1.23$ GHz, $N_{\parallel} = 1.9$ indicate that the LH resonance moves closer to the low density surface region as frequency lowers, leading to strong ion heating near the surface. The best frequency choice is a compromise, $f \simeq 3.7$ GHz, based on these considerations:

Figure 2.B-6. TIBER II poloidal flux contours: $A = 3.6$, elongation = 2.4, $d = 0.4$, $R_0 = 3.0$ m; baseline case with $\beta = 0.05$ and $I = 10$ MA.

Figure 2.B-7. Baseline plasma illustrating poloidal flux, pressure, toroidal current density, and safety factor in the midplane.

Figure 2.B-8. (a) Ray trajectory and index of refraction along path; LH wave launched near bottom, with $n_u = 1.5$ and $f = 4.6$ GHz. (b) Electron damping along ray path; ion damping is negligible. (c) $G = \langle J_{\parallel} B \rangle / \langle B^2 \rangle$, a measure of current density averaged over a flux surface. The abscissa labeled PSIHAT is the normalized poloidal flux: zero is the edge and unity is the magnetic axis. The G function actually generated is dotted, and the solid curve represents smoothing and averaging.

(a) preliminary engineering studies at 4.6 GHz showed difficulty in cooling the grill septa, which are too thin; (b) to guarantee LHCD capability for low n_e current ramp or transformer recharge a frequency $f \gtrsim 2$ GHz is desirable; and (c) from LHCD experience [4] it is felt that $f \simeq 3.7$ GHz could provide steady-state operation at $\bar{n}_e \simeq 0.8 \times 10^{20} \text{ m}^{-3}$ or on the surface of higher density tokamaks.

For the 3.7-GHz frequency we consider next the best launch location. The bottom position is inconvenient because of difficulties in locating the LH waveguides between the toroidal coils. As the starting point moves towards the midplane ($\theta = 0$) the degree of maximum wave penetration gradually decreases. Calculations show that a grill complex spanning $-0.15 \lesssim \theta \lesssim 0$ promises the best penetration. Figure 2.B-9 shows ray trajectories and $G = \langle j_{\parallel} B \rangle$ for a series of rays ($N_{\parallel} = 2.3, 2.2, 2.1, 2.0, 1.90$) starting at $\theta = -0.15$.

Finally, the power in each ray needs adjustment so that the function $G(\psi)$ duplicates the profile needed to achieve the desired equilibrium. Figure 2.B-10 shows the contribution to G from LHCD when 22.6 MW is divided among three rays, $N_{\parallel} = 3.5, 2.4$, and 1.9, assuming no bootstrap current. (The remaining current density, in the interior, must of course be provided by other means, e.g., by neutral beams or, as in this figure, by the fast wave.) When there is a bootstrap contribution the LHCD power decreases; Fig. 2.B-11 shows the self-consistent MHD equilibrium with LH surface currents when the bootstrap is present in our flat density baseline. For this case only 16.5 MW of LH power is needed.

Our calculations indicate significant further reductions in LHCD power are achieved if in fact the density profile becomes more peaked, as is the present TIBER II philosophy. Thus, 20 MW of power at 3.7 GHz appears adequate for the present operating mode options.

Figure 2.B-9. Ray paths (Z vs R) and current density (G vs Ψ) generated by wavespectrum, $2.3 > n_{\parallel} > 1.9$, at $f = 3.7$ GHz. Note LH waves do not penetrate deeply into plasma. Normalized diamagnetism function $F' = dF/d\Psi$ and its derivative are shown, as generated by LHCD. Ray spectrum must duplicate solid curve $G(\Psi)$ to duplicate baseline equilibrium.

Figure 2.B-10. Converged, self-consistent RFCD/MHD equilibrium. LH power = 23 MW is divided among three modes, which generate G up to $\tilde{\psi} \approx 0.3$. Some current driver other than LH waves (fast waves in this illustration) is needed in the interior. Note safety factor is well-behaved. Small scale variations in F' and toroidal current density can be removed by using a higher mode density.

Figure 2.8-11. Same equilibrium as Fig. 2.8-10, except bootstrap current is included. In this case only 17 MW of LH power is needed.

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APPENDIX 2.C

NEUTRAL-BEAM AND NEOCLASSICAL CURRENT FROM A FOKKER-PLANCK CODE

In this appendix we describe calculations of neutral-beam and bootstrap current carried out with the aid of the FPT (Fokker-Planck Transport) code [1]. This code numerically solves the Fokker-Planck equation in $v_{\perp}$ - $v_{\parallel}$ space for the distribution function, f_{FI} , of the fast ions deposited by the high-energy neutral beams. The solution is obtained as a function of radius for a circular plasma cross section in the approximation of infinite aspect ratio, i.e., no banana orbits and no bounce-averaging. Collisions of fast ions with background Maxwellian ions and electrons as well as self collisions are included. The deposition of beam ions is carried out with a Monte Carlo code which does include toroidal effects and allows for focusing of the neutral beams in the horizontal and vertical planes. Fixed radial temperature and density profiles and constant Z_{eff} are used, as described in Section 2.2. The current carried by the fast ions is then calculated from

$$j_{FI}(r) = e Z_b \int v_{\parallel} f(v_{\perp}, v_{\parallel}, r) d^3v \quad (2.C-1)$$

where $eZ_b = e$ is the charge of the deuterium ions deposited by the beam and $v_{\parallel}$ and $v_{\perp}$ are the components of the velocity parallel and perpendicular to the magnetic field.

Because of drag on the circulating fast ions, the thermal electrons will accelerate to partially cancel the fast-ion current. If the effective charge, Z_{eff} , of the background plasma differs from that of the beam ions, the drag on the thermal plasma will prevent this induced electron current from totally canceling the ion current. In addition, a fraction of the electrons are trapped (approximately $\epsilon^{1/2} = (r/R_0)^{1/2}$, where R_0 is the major radius), and cannot carry canceling current. Calculation of this electron current must be carried out with the electron Fokker-Planck equation. This has been done by Start and Cordey [2]. Their results can be presented as a correction factor to the fast-ion current,

$$j_{NB} = j_{FI} [1 - Z_b/Z_{eff} + g(\epsilon, Z_b/Z_{eff})] = F_e j_{FI}$$

where F_e has been tabulated in Ref. [2].

Because of the steep density gradients and high poloidal beta, we expect significant neoclassical currents to be generated in TIBER (see the discussion in Sect. 2.2). This can be computed simultaneously with the neutral-beam-induced current. Since the density and temperature profiles are fixed by assumption, it is necessary only to know the local poloidal magnetic field, B_p , to obtain the current due to neoclassical effects [see Eq. (2.2-25)]. B_p is found from Ampere's law in cylindrical geometry,

$$\frac{\partial}{\partial r} (r B_p) = \mu_0 r j \quad (2.C-3)$$

where $j = j_{NB} + j_{BS}$ is the total local current density. The bootstrap current density, j_{BS} , is computed from an expression [3] similar to Eq. (2.2-25) for plasmas with Z_{eff} differing from unity. One should also include in j the current generated by lower hybrid waves. Unfortunately, no code now exists which simultaneously computes both neutral-beam- and lower-hybrid-driven currents. Since the lower hybrid current is only in the outer edge of the plasma, it is this region only where the poloidal field, and thus the bootstrap current, might be in error.

These computations are dependant on the properties of the background plasma. The plasma properties are in turn dependant on the current-drive efficiency and the fraction of bootstrap current. It was therefore necessary to iterate between the properties computed with the MUMAK code (see Sections 2.2, 2.4 and Appendix 2.A) and the Fokker-Planck calculations described in this section until a self-consistent solution was obtained. Essentially, this amounts to adjustments of the coefficient f_0 in Eq. (2.2-24) for the current-drive efficiency and the leading coefficient in Eq. (2.2-26) for the fraction of current arising from bootstrap effects. The results for the baseline parameters of Sect. 2.4 are a bootstrap current fraction of approximately 0.29 (a slightly lower value of 0.27 was used in the actual baseline calculations since the iteration had not yet converged at the time the properties were frozen) and a neutral-beam current drive efficiency of

$\gamma = 0.33$. The latter leads to a coefficient $f_D = 3.4$ in Eq. (2.2-24). The profiles for j_{NB} and j_{BS} shown in Section 2.4 have been obtained from this computation.

Two effects that will decrease the efficiency have been neglected in these calculations: plasma toroidal rotation and banana trapping. Because of the momentum imparted to the plasma by the unidirectional beams, the plasma will rotate, thus effectively reducing the speed of the fast ions and lowering the efficiency. Some of the fast ions will collisionally trap in banana orbits. They can subsequently detrapp with equal probability either co- or counter to the current. Those that detrapp in the counter direction will subtract from the fast ion current, again lowering the efficiency. Both of these effects are estimated to require small corrections to the efficiency computed by the methods of this appendix. In order to be conservative, we have arbitrarily taken a reduced efficiency of $\gamma = 0.25$ for the baseline calculations of Sect. 2.4. The coefficient in Eq. (2.2-24) then becomes $f_D = 2.6$. In considering possible advanced operating modes in Section 2.5, a reduced "safety factor" was employed and we used $f_D = 3.1$.

Some additional calculations were carried out to verify the scaling of the bootstrap current with poloidal beta, β_p , shown in Eq. (2.2-26). Since the poloidal field at the edge of the plasma, $B_p(r=a)$, scales with the total plasma current, I_p , we can hold the plasma pressure constant and adjust the neutral-beam current to give varying values of the plasma current. The dependance of bootstrap fraction of the current, I_{BS}/I_p , on β_p was found to obey the linear relationship of Eq. (2.2-16) over the range studied: β_p from 0.73 to 1.35.

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APPENDIX 2.D

EDGE PLASMA TRANSPORT AND NEUTRAL PARTICLE SIMULATION IN ETR

2.D.1 EDGE PLASMA TRANSPORT

It is proposed to analyze scrape-off layer (SOL) and divertor plasma conditions for ETR using the ODESSA edge plasma transport code developed at UCLA. The code is quasi-two-dimensional [1,2] and provides axially averaged radial profiles of plasma density, temperature (ion and electron) and flow velocity in the upstream zone or SOL and downstream zone or divertor chamber (adjacent to the target). The model accounts for the important edge transport processes (e.g., flux-limited parallel conduction, anomalous cross-field transport, sheath effects, shear and parallel viscosity, etc.), as well as recycling at the target. A simplified but reasonably accurate neutral transport model is incorporated that provides radial profiles of particle sources and electron energy sinks due to recycling neutrals. With minor modifications the code can be readily iterated with a Monte Carlo neutral code to obtain more realistic source profiles.

ODESSA has been extensively used to study the stability of plasma conditions in the divertor as a function of heat and particle fluxes leaving the core and the neutral exhaust rate [3]. Results from the code have been compared against two-dimensional simulations of the NET divertor. Good agreement was found for both magnitudes and profile shapes of plasma parameters, indicating that the physics model in ODESSA is adequate. Supplemented with analytical modeling, numerical computations using ODESSA have established criteria that must be satisfied by input fluxes to achieve stable high recycling divertor conditions. This work has been successfully applied in the context of the TITAN RFP reactor study, currently approaching completion at UCLA, to the identification of parameter space (SOL width, neutral pumping time, etc.) that yields an adequate divertor design. In the process, unusual physics phenomena were encountered that result in profile shapes somewhat contrary to what one might expect from less detailed models. In particular, the occurrence of locally reversed flow (i.e., flow out of the divertor at certain locations in the SOL) results in greater than expected cross-field transport of particles and energy in the upstream region (around

the midplane) yielding relatively flat profiles. This raises the possibility of significant recycling at the first wall. These factors will clearly influence the location of the first wall relative to the separatrix.

The above phenomena are also expected to arise in the ETR divertor (in fact, certain features are universal to high recycling divertors), and a systematic study will clearly be necessary. Although ODESSA does not provide a detailed two-dimensional map of plasma parameters, it does yield the important profile information for estimating particle and heat loads on the target and first wall at a fraction of the computational cost typical of two-dimensional simulations. Very minor modifications to the code will make it relevant to the ETR divertor configuration.

2.D.2 NEUTRAL PARTICLE TRANSPORT

In order to establish the optimum location of the divertor target relative to the null-point and to assess the neutral exhaust efficiency of the divertor, it is necessary to carry out detailed Monte Carlo neutral transport calculations for a given divertor geometry. Results from these simulations will provide valuable information on the probability of neutral leakage back across the separatrix into the core (which must be minimized for high recycling), on the charge exchange flux and neutral particle energies striking the divertor target and first wall (to estimate erosion rates), and on the ionization sources necessary to iterate with ODESSA. It is proposed to perform these calculations using the Princeton DEGAS Monte Carlo code [4] with which extensive operating experience exists at UCLA. The code has been most recently applied to the complex 3-D geometry of the TITAN RFP divertor. Earlier experience with DEGAS consists of simulations of pump limiter experiments [5,6] and optimizing pump limiter designs for maximum helium enrichment. [7]

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APPENDIX 2.E

TADDDPAK - A COMPUTER CODE PACKAGE FOR MODELING THE THERMAL EFFECTS ON THE DIVERTOR IN TOKAMAKS DURING PLASMA DISRUPTIONS

2.E.1 DESCRIPTION

TADDDPAK [1] is a quasi-two-dimensional Fortran computer code package for modeling the time-dependent and positional-dependent thermal and mechanical effects on the divertor during plasma disruptions. This code package has been created by modifying and coupling several pre-existing computer codes.

The basic structure of TADDDPAK is illustrated in Fig. 2.E-1. DSTAR [2], a plasma simulation code written by B. Merrill and S. Jardin, calculates the power leaving the edge of the plasma as a function of time during a disruption. Either this power distribution, or a user-defined distribution based on current experimental data for disruptions, is used as input for HT-K76E [3]. W. Barr originally wrote HT-K76E as a power mapping code to assist in divertor design under steady-state conditions. His IBM Basic version has been translated to Fortran and adapted for the time dependent problem of a plasma disruption. This new version of the code, which is the first component of TADDDPAK, calculates the time-dependent power densities (MW/m^2) on a poloidal cross section of the inner, lower, and outer plates of the TIBER divertor.

Once these power densities are determined, they are then used as input by one of two thermal analysis codes. The thermal analysis codes are LESLIE [4], written by J. Haines, and SOAST [5], written by D. Croessmann. Both codes solve the moving boundary heat equation in one-dimension. LESLIE, in addition to providing one dimensional thermal analysis, also calculates mechanical effects to include stress and strain. SOAST, an abridged version of the comprehensive A*THERMAL Code [6], provides similar thermal analysis to LESLIE and offers a simple vapor shielding model to be used at the user's discretion. The selected code's output is then integrated into DISSPLA (Display Integrated Software System and Plotting Language) [7] which generates three-dimensional graphics of selected positional and time dependent variables.

An example of TADDDPAK's graphic output is illustrated in Fig. 2.E-2. This figure depicts the vaporization rate, total vaporization, and total mass ablated from the divertor's inner plate for either the upper or lower half

Figure 2.E-1. The module structure of the TADDPACK code.

Figure 2.E-2. An example of TADDPACK output for the vaporization rate and ablation on the outer divertor plate of TIBER II.

of the machine, assuming symmetry. Additional information and examples of TADDDPAK's results are available in Ref. [8].

2.E.2 OBJECTIVES

TADDDPAK was developed for several reasons. It was created to minimize the time requirement for parametric analysis of the divertor as compared to DSTAR, which uses 30 to 40 minutes of CPU time per calculation. Since TIBER is one of many future tokamak fusion reactor designs, TADDDPAK was designed to be generic in the sense that it can be adapted to any divertor/tokamak configuration. It was also designed to provide two-dimensional presentation and self-consistent peaking factors for the following positional- and time-dependent variables: heat flux, surface temperature, vaporization rate, total vaporization, and mechanical effects. It is the intention that this package will provide a convenient and rapid user-friendly medium to facilitate parametric analysis of disruption effects on divertors, and thereby provide relevant data for divertor design.

2.E.3 ISSUES UNDER INVESTIGATION

Although limited calculations were completed for TIBER II at the time of this report, complete calculations will be made for TIBER II as well as for the next generation of tokamak fusion reactor design. These calculations will be made in regard to two important issues. The first issue is the broad spectrum of thermal effects, and this spectrum's dependence on the heatflux pulse shape at the edge of the plasma during a disruption. The second issue is candidate material choices and how different materials respond under a disruption environment as compared to steady state. The materials of interest are carbon, carbon-carbon composites, beryllium, and tungsten. Parametric studies will be made to investigate these two issues.

2.E.4 FUTURE IMPROVEMENT TO TADDDPAK

Two features of TADDDPAK are earmarked for future modification. To improve the modeling ability of the thermal analysis portion of this code package, the currently existing vapor shielding model requires modification. A self-consistent approach is needed to make this improvement viable, perhaps based on the methodology of plasma vapor shielding for pellet fueling. The scrape-off layer model in the mapping portion of the code package also needs

improvement. Currently the assumption is made that the e-folding thickness remains constant during the thermal quench of the disruption. A time-dependent thickness would be more valid. A model for this phenomenon is needed.

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APPENDIX 2.F.

BEAT-WAVE CURRENT DRIVE WITH INTENSE PULSED FREE-ELECTRON LASERS

High-power free-electron lasers (FELs) make new methods possible for driving current in toroidal devices with electromagnetic waves. Earlier considerations of beat-wave current drive are generalized to a hot magnetized plasma and an arbitrary beat wave. Here the beating of two electromagnetic waves resonantly excites a low-frequency beat wave that accelerates and heats electrons and leads to a current. The absolute current-drive efficiency depends nonlinearly on the two pump-wave intensities and is constrained by the Manley-Rowe relations. Accessibility at high plasma densities is not a difficulty, but a degree of frequency tunability in the wave sources is required. Particle simulations indicate that there is good coupling to an electron velocity tail for a Langmuir beat wave with a phase velocity 2 to 4 times $(T_e/m_e)^{1/2}$, so that all of the high-frequency wave source is absorbed and the beat wave damps completely on the electrons. A realistic scenario for a beat-wave current-drive experiment in the Livermore Microwave Tokamak Experiment has been calculated using the TORAY toroidal ray tracing code, and the scaling to an Engineering Test Reactor plasma is described.

2.F.1. INTRODUCTION

The efficient production of microwaves with an intense pulsed free-electron laser (FEL) has recently been demonstrated [1]. Such intense wave sources are attractive for plasma heating [2,3] and current-drive applications in toroidal confinement devices. Here we analyze a specific current-drive scheme that uses the nonlinear beat of two high-frequency FEL sources to resonantly excite a longitudinal wave in the plasma that accelerates and heats electrons leading to a current. This process is closely akin to the beat-wave acceleration of electrons by electromagnetic waves proposed some years ago by Dawson and Tajima [4] for advanced accelerators, for which some experimental data base now exists [5]. The analytical and numerical calculations presented here generalize and extend an earlier calculation of beat-wave current drive [6]. A realistic beat-wave current-drive scenario has been calculated for the

Microwave Tokamak Experiment (MTX), and the scaling of the beat-wave current-drive efficiency to an ETR is described. Beat-wave current drive has the advantage that, given no limit on the frequencies of the two wave sources, the entire plasma is accessible and both the spatial interaction region and the region of electron velocity space accelerated can be controlled. It has the disadvantage that the scheme is essentially nonlinear; there are well-defined intensity requirements for efficient current drive; and competition with other, possibly deleterious nonlinear effects must be avoided.

The discussion here is organized as follows. In Sec. 2.F.2 the general beat-wave coupling equations are given in the eikonal limit for two high-frequency electromagnetic pump waves resonantly coupling to a low-frequency electrostatic wave in a hot, magnetized, nonuniform plasma. A summary of recent particle simulation work on beat-wave current drive is presented in Sec. 2.F.3. The simulations demonstrate the importance of having the beat-wave phase velocity fall in the optimal part of the electron velocity distribution so that there are an adequate number of electrons to efficiently dissipate the beat wave and also to form a high-energy electron tail. Section 2.F.4 contains a discussion of the current-drive efficiency to be expected from the beat-wave mechanism. The current-drive efficiency is constrained by the Manley-Rowe relations, i.e., the beat wave can carry off at most a fraction of the energy given by the ratio of its frequency to the higher frequency pump wave if only a single resonant three-wave interaction occurs. The nonlinear conversion rate of photons to the beat wave also enters. The constraints on frequency and wavenumber matching and on the beat-wave phase velocity relative to the electron velocity distribution determine a relation between the electron temperature, the plasma density, and the frequencies of the pump waves. Section 2.F.5 concludes with a description of a beat-wave current-drive scenario for MTX and how this scales to the Engineering Test Reactor (ETR).

2.F.2. BEAT-WAVE COUPLING FORMALISM

In this section we present envelope equations describing the nonlinear coupling of two high-frequency transverse waves to a lower-frequency electrostatic wave in a hot, nonuniform, magnetized plasma. Steady-state analytical solutions describing the transfer of photons (and, hence, energy and momentum) from the high-frequency transverse wave to the lower-frequency

transverse wave are reviewed for the cases of parallel and anti-parallel propagation of the transverse waves. From conservation of energy and momentum, one can immediately deduce the energy and momentum available to the beat wave. Our eikonal formulation of beat-wave coupling assumes that the plasma and magnetic field scale lengths are long compared to the wavelengths and admits an arbitrary polarization of the low-frequency electrostatic beat wave.

Envelope equations are straightforwardly derived from the perturbation theory for parametric instabilities presented in Refs. [7] and [8]. The dispersion equation for a transverse pump wave with frequency and wavenumber $(\omega_2, \underline{k}_2)$ coupled to a low frequency wave $(\omega_3, \underline{k}_3)$ and a finite-amplitude pump wave $(\omega_1, \underline{k}_1)$ is [7,8]

$$D_2 - \frac{\chi_e (1 + \chi_i)}{\epsilon} \frac{k_3^2 u_1^2}{4} |\hat{k}_2 \times \hat{u}_1|^2 = 0 \quad (2.F-1)$$

where $D_2 = (-k_2)^2 c^2 - (-\omega_2)^2 \epsilon(-\omega_2, -\underline{k}_2)$, $\epsilon(\omega_3, \underline{k}_3)$ is the longitudinal dielectric, $\epsilon = 1 + \chi_e + \chi_i$, $\chi_e(\omega_3, \underline{k}_3)$ and $\chi_i(\omega_3, \underline{k}_3)$ are the linear electron and ion susceptibilities, $\underline{u}_1$ and $\underline{u}_2$ are the linear electron quiver velocities in the transverse waves, $\hat{k}_{1,2}$ and $\hat{u}_{1,2}$ are unit vectors, and there is a restriction on the relative polarizations of the pump waves. The corresponding equation for the transverse pump wave with frequency $(\omega_1, \underline{k}_1)$ is

$$D_1 - \frac{\chi_e (1 + \chi_i)}{\epsilon} \frac{k_3^2 u_2^2}{4} |\hat{k}_1 \times \hat{u}_2|^2 = 0 \quad (2.F-2)$$

where $D_1 = k_1^2 c^2 - \omega_1^2 \epsilon(\omega_1, \underline{k}_1)$. These equations are valid if $|u_{1,2}/c| \ll 1$, $\omega_{1,2} \gg \omega_3$, Ω_i , ω_{pi} and $k_{1,2} \rho_e \ll 1$, where ρ_e is the electron Larmor radius, v_e is the electron thermal velocity, ω_{pi} and Ω_i are the ion plasma and cyclotron frequencies.

The restriction on the relative polarization of the pump waves comes from the structure of the nonlinear coupling terms in the fundamental equations. The nonlinear wave equation for the transverse pump waves is [7,8]

$$D(\omega_t^{(1,2)}, \underline{k}_t^{(1,2)} - i\nabla) \underline{E}_t(\omega_t^{(1,2)}, \underline{k}_t^{(1,2)}) = \frac{i4\pi e}{\omega_t} \frac{n_e(\omega_3, \underline{k}_3)}{n_0} \underline{\Gamma}_e^* (\omega_t^{(2,1)}, \underline{k}_t^{(2,1)}) \quad (2.F-3)$$

where $\underline{\Gamma}_e$ is the high-frequency linear electron flux density, $\underline{E}_t$ is the transverse electric field, $\underline{k}_t^{(1,2)} = \underline{k}_1, -\underline{k}_2$ and $\omega_t^{(1,2)} = \omega_1, -\omega_2$ and $D(\omega_t, \underline{k}_t) = k_t^2 c^2 - \omega_t^2 \epsilon(\omega_t, \underline{k}_t)$. Thus, the coupling is most efficient when $\underline{E}_t^{(1,2)}$ and $\underline{\Gamma}_e^{(2,1)}$ have the same polarization. This is also manifested in the low-frequency ponderomotive force density [8]

$$\delta \underline{f}_e(\omega_3, \underline{k}_3) = -e \underline{\Gamma}_e(\omega_t^{(1,2)}, \underline{k}_t^{(1,2)}) \times \underline{B}(\omega_t^{(2,1)}, \underline{k}_t^{(2,1)})/c - m_e \nabla \cdot (\underline{\Gamma}_e^{(1,2)} \underline{\Gamma}_e^{(2,1)})/n_0 \quad (2.F-4)$$

where $\underline{B}(\omega_t, \underline{k}_t) = c \underline{k}_t \times \underline{E}_t(\omega_t, \underline{k}_t)/\omega_t$. Specific examples of beat-wave coupling illustrating the dependence of the coupling on the relative polarizations will be presented in the following.

Equations (2.F-1) and (2.F-2) were derived to describe the coupling of three normal modes. The coupling is efficient if the frequency and wavenumber matching conditions are satisfied:

$$\omega_1 = \omega_2 + \omega_3 \quad (2.F-5)$$

and

$$\underline{k}_1 = \underline{k}_2 + \underline{k}_3. \quad (2.F-6)$$

These phase-matching conditions are intimately related to energy and momentum conservation [6].

Consider Eqs. (2.F-1) and (2.F-2) in the steady-state limit, which should be appropriate for times longer than the inverse damping rate of the beat wave. Expanding $-\nabla^2 = k_t^2 - 2ik_t \cdot \nabla$ in the eikonal limit and introducing action fluxes in natural units, $\underline{J}_{1,2} = (k_{1,2}/2\pi) |u_{1,2}/c|^2$, for pump waves with identical polarization, Eqs.(2.F-1) and (2.F-2) become

$$\hat{k}_1 \cdot \nabla J_1 = \frac{-2\pi k_3^2}{k_1 k_2} J_1 J_2 \text{Im}[\chi_e(1 + \chi_i)/\epsilon] \quad (2.F-7a)$$

$$\hat{k}_2 \cdot \nabla J_2 = \frac{2\pi k_3^2}{k_1 k_2} J_1 J_2 \text{Im}[\chi_e(1 + \chi_i)/\epsilon] \quad (2.F-7b)$$

The linear susceptibilities and the longitudinal dielectric are evaluated at the low-frequency difference frequency and difference wavenumber in Eq. (2.F-7). Note that for coupling to electron normal modes $\chi_e \approx -1$ and $1 + \chi_i \approx 1$ to good approximation. Absolute action fluxes are given by multiplying the action fluxes in natural units by $(m_e c^2/e)^2$. The power and momentum fluxes in the transverse waves are given by the product (or dyad product) of the wave frequency and wave vector, respectively, with the absolute action fluxes [6]. For the special cases of parallel and anti-parallel propagation of the two transverse pump waves, the coupled one-dimensional differential equations describing action conservation result:

$$-\frac{\partial}{\partial x} J_1 = \pm \frac{\partial}{\partial x} J_2 = \frac{2\pi k_3^2}{k_1 k_2} J_1 J_2 \text{Im}[\chi_e(1 + \chi_i)/\epsilon] . \quad (2.F-8)$$

The upper/lower sign in Eq. (2.F-8) is for parallel/anti-parallel transverse waves. Thus, for every photon at frequency ω_1 that nonlinearly decays, a photon is nonlinearly emitted at frequency ω_2 . In consequence of Eqs. (2.F-5), (2.F-6), and (2.F-8), there is an amount of energy flux proportional to $\omega_3 \Delta J$ and momentum flux proportional to $k_3 \Delta J$ deposited in the beat wave. This energy and momentum is available for accelerating electrons and driving current. Conservation of energy, momentum, and wave action have a profound influence on any beat-wave acceleration or heating scheme [6].

Equation (2.F-8) was analytically solved for forward scatter (parallel propagation) and backward scatter (anti-parallel propagation) to obtain implicit relations for the relative action transfer $R_a = \Delta J/J$ given an input ratio $\rho = J_2^{\text{in}}/J_1^{\text{in}}$. J_1^{in} and J_2^{in} are the input action fluxes from the same side for forward scatter or opposite sides for backscatter. For relatively weak dissipation so that the beat wave is a well-defined normal mode and a narrow resonance zone, but wide compared to the wavelengths involved, the action

transfer from the first wave to the decay products is independent of the damping rate of the beat wave: $\int_{-\infty}^{\infty} dx \operatorname{Im} \epsilon^{-1} = -\pi [d\epsilon/dx]^{-1} = -\pi L$ for linear gradients [9] and assuming that there is no singular behavior in $\chi_e(1 + \chi_i)$ or $k_3^2/k_1 k_2$. This extends earlier results to the case of an arbitrary beat wave in a magnetized plasma. The solution for the action transfer for forward scatter is determined by [6]

$$\bar{J}_1 (k_3^2/k_1 k_2) \pi k_1 L |u_1^{\text{in}}/c|^2 = (1 + \rho)^{-1} \ln [(\rho + R_a)/(1 - R_a)\rho], \quad (2.F-9a)$$

and for backward scatter by [9]

$$\bar{J}_1 (k_3^2/k_1 k_2) \pi k_1 L |u_1^{\text{in}}/c|^2 = (1 - R_a - \rho)^{-1} \ln [(1 - R_a)(\rho + R_a)/\rho]. \quad (2.F-9b)$$

Numerical solutions of Eqs. (2.F-9a) and (2.F-9b) for R_a as a function of $\bar{J}_1$ and the ratio of input action fluxes ρ are displayed in Fig. 2.F-1 for forward and backward scatter. One notes that in either case there is efficient action transfer ($R_a > 0.1$) if $\bar{J}_1 \rho^{1/2} > 0.1$ and $R_a \rightarrow 1$ for $\bar{J}_1 \rho^{1/2} > 1$. However, for the same value of $|u_{1,2}/c|^2$ the value of $\bar{J}_1$ is typically much smaller for forward scatter where $k_3 = k_1 - k_2$ is small, than the value of $\bar{J}_1$ for backscatter where $k_3 = k_1 + k_2$. Thus, the coupling and the relative action transfer for backscatter is generally much larger than for forward scatter. There are two exceptions to this. First, if the lower frequency transverse wave is near cutoff, then k_3 is nearly the same for forward and backward scatter ($k_3 = k_1$); furthermore, $\bar{J}_1$ is significantly enhanced because $k_2 \rightarrow 0$. This is a vestige of an absolute instability for the linear parametric decay to a daughter transverse wave at its turning point [10], and our eikonal steady-state mode-coupling equations become invalid. A major practical difficulty in taking advantage of the enhancement of the coupling when $k_2 \rightarrow 0$ is that wave refraction for the second pump wave becomes a problem and the resonance zone in which there is good phase matching may be too small. The second instance in which the coupling for backscatter may be weaker than for forward scatter occurs when the beat wave for backscatter is severely damped because of its significantly smaller phase velocity.

Figure 2.F-1. The relative action transfer $R_a = \Delta J/J$ as a function of the dimensionless action flux density J and the input ratio $\rho = J_2^{in}/J_1^{in}$ for (a) forward and (b) backward scatter.

Thus far, we have not specified the type of beat wave, except that it is a longitudinal wave. In order to efficiently accelerate electrons, we have considered Langmuir waves propagating parallel to the magnetic field in the torus. Resonance with the electrons is the simple Landau type: the beat-wave phase velocity must match the parallel electron velocity. However, we can just as easily consider an obliquely propagating upper hybrid wave for which resonant electrons must satisfy cyclotron resonance. This offers the opportunity for cyclotron-auto-resonant acceleration [11] of the electrons, which could lead to much more energetic current-carrying electrons. Because more energetic current-carrying electrons would be less collisional, the resulting current-drive efficiency might be significantly improved. In the present discussion, we will only consider Langmuir beat waves; a treatment of a cyclotron-auto-resonant beat wave will be addressed in a future communication.

2.F.3. PARTICLE SIMULATIONS

Beat-wave current drive involves the nonlinear coupling of intense electromagnetic waves to normal modes in a hot magnetized plasma with the intent to deposit a finite amount of momentum and energy into resonant electrons. Nonlinear phenomena play a central role at several steps of this process, and the self-consistency of the waves and the concomitant plasma modification is important if high efficiencies are desired. These factors motivate the use of a self-consistent electromagnetic particle simulation as a means of studying some of the basic physics of beat-wave current drive. Because a first-principles solution of Maxwell's equations and the relativistic equations of motion for the plasma are contained in the simulation code, the particle code can test our analysis of beat-wave coupling. Furthermore, the simulation gives the complete time-dependent solution to the coupled equations and not just the steady state. In addition, the simulations give valuable insight on nonlinear mechanisms affecting the saturation of the resonantly excited beat wave, on competition with parasitic stimulated scattering instabilities in which the pump waves and the beat wave may participate, and on the role of ions in possibly diverting the beat-wave momentum from being deposited in an electron tail. Our principal simulation finding is that there is a range of beat-wave phase velocities, from approximately 2 to $4 \times (T_e/m_e)^{1/2}$ in a Maxwellian plasma, that facilitates

complete damping of a Langmuir beat wave on an electron tail with little or no parasitic ion damping nor unwanted stimulated scattering parametric instabilities.

The simulation results reviewed here focus on the coupling of a large-amplitude beat wave to the electron distribution. We are interested in conditions for which there can be complete pump depletion of the higher-frequency pump wave, i.e., the relative action transfer R_a approaches one. The important questions addressed in the simulations were the competition of the beat-wave coupling of the transverse pump waves with parasitic stimulated scattering instabilities [12] and the competition of the beat-wave damping on an electron tail with nonlinear damping involving ion motion. Earlier simulations of beat-wave heating with immobile ions confirmed the results of Eq. (2.F-9) describing action transfer for the beating of opposed transverse wave pumps in a nonuniform plasma with open boundary conditions [13]. The earlier simulation results and those reported here were produced with the one-dimensional, relativistic, electromagnetic particle code EMONE [13]. The recent simulations modeled both parallel and anti-parallel propagation of the pump waves in a uniform unmagnetized plasma with mobile ions ($m_i/m_e = 100$), initially equal electron and ion temperatures, and periodic boundary conditions.

The recent simulations indicate that with electron quiver velocities equal to a few percent of c (as proposed in MTX), ratios of $\omega_1/(\omega_1 - \omega_2) = 3$ to 5, and careful choice of the ratio of the beat-wave phase velocity $\omega_3/k_3 v_e = 2$ to 4, complete transfer of action (pump depletion) of the higher-frequency pump wave can be achieved and the beat wave is efficiently damped by an electron tail extending to $\sim 6v_e$. Furthermore, the momentum deposited in the plasma goes almost entirely into the electrons. In fact, the ions ultimately drifted backwards in a number of simulations, allowing the electrons to acquire additional momentum ($\leq 20\%$) and enhancing the current drive. Action, energy and momentum conservation were good. Large-amplitude beat waves that coupled well to the electrons were obtained for both parallel and anti-parallel propagating transverse pump waves (Figs. 2.F-2 and 2.F-3).

The beat waves exhibited a significant maximum amplitude in good agreement with analysis [14,15]

Figure 2.F-2. Particle simulation results for parallel propagating transverse waves with parameters $\omega_1 = 2.24\omega_{pe}$, $\omega_2 = 1.12\omega_{pe}$, $v_e/c = 0.22$, $m_i/m_e = 100$, $u_1^{in}/c = u_2^{in}/c = 0.04$, and $(\omega_1 - \omega_2)/(k_1 - k_2) = 3.4 v_e = 0.75 c$. (a) The higher frequency pump-wave energy density vs time. (b) The lower frequency pump-wave energy density components vs time. (c) The longitudinal field energy spectrum vs wavenumber at $\omega_{pe} t = 147$ showing the beat-wave peak. (d) The electron velocity distribution vs velocity parallel to the beat-wave phase velocity at $\omega_{pe} t = 980$ showing the formation of a tail $v_e \sim 0.6-0.7 c$. The trapping width in velocity was as large as $0.6 v_e$. There was 80% of the beat-wave momentum deposited in the electrons and 20% in the ions.

Figure 2.F-3. Particle simulation results for opposed propagation of the transverse pump waves with parameters $\omega_1 = 2.24$, $\omega_2 = 1.12$, $v_e/c = 0.13$, $m_i/m_e = 100$, $u_1^{in}/c = u_2^{in}/c = 0.04$, and $(\omega_1 - \omega_2)/(k_1 + k_2) = 3.5 v_e = 0.45 c$. (a) The higher frequency pump-wave energy density vs time. (b) The lower frequency pump-wave energy density components vs time. (c) The longitudinal field energy spectrum vs wavenumber at $\omega_{pe} t = 98$ showing the beat-wave peak. (d) The electron velocity distribution vs velocity parallel to the beat-wave phase velocity at $\omega_{pe} t = 980$ showing the formation of a tail. The trapping width in velocity was as large as v_e . There was 120% of the beat-wave momentum deposited into the electrons and -20% into the ions.

$$|k_3 \delta|_{\text{sat}} = \left| \frac{k_3^2 e \phi}{m_e \omega_{pe}^2} \right|_{\text{sat}} = \left| \frac{\omega_{pe} k_3^2 u_1 u_2}{4 \Delta} \right| \sim 0(10^{-2} \text{ to } 10^{-1}) \quad (2.F-10)$$

for our simulation parameters, where δ is the Lagrangian displacement amplitude of an electron in the beat wave, $\Delta = \omega_1 - \omega_2 - \omega_{BG}$, and ω_{BG} is the Bohm-Gross frequency. At these amplitudes the trapping velocity in the beat wave was substantial, $v_t = |e\phi/m_e|^{1/2} \sim 0.1c \sim v_e$. Because $|k_3 \delta|_{\text{sat}} \ll 1$, the beat-wave amplitude was well below the wave-breaking limit,

$$|k_3 \delta| = 1 \quad (2.F-11)$$

Furthermore, the limit on the beat-wave amplitude set by relativistic detuning [15] is

$$|k_3 \delta| = \left(\frac{k_3 c}{\omega_{pe}} \right)^{4/3} \left(\frac{16 u_1 u_2}{3 c^2} \right)^{1/3} \quad (2.F-12)$$

which is much larger than the value determined by Eq. (2.F-10) for our simulations.

Conservation of action also sets an upper limit on the beat-wave amplitude. In the simulation model considered here, the action density is conserved rather than the action flux density [6]. The beat wave cannot acquire more action than was initially resident in the first pump wave, hence

$$E_3^2 / \omega_{pe} \leq (E_1^{\text{in}})^2 / \omega_1 \quad (2.F-13)$$

or equivalently

$$|k_3 \delta| \leq \left(\frac{\omega_1}{\omega_{pe}} \right)^{1/2} \left| \frac{k_3 u_1^{\text{in}}}{\omega_{pe}} \right| \quad (2.F-14)$$

Because we established conditions such that the beat wave was strongly resonant, the beat-wave interaction was able to deplete the first pump wave; and the estimates of Eqs.(2.F-10) and (2.F-14) were nearly coincident in our simulations. In summary, the saturated beat-wave amplitude and the pump

depletion in the simulations were consistent with the theoretical expectations. Once the beat-wave amplitude peaked and the higher-frequency wave became severely depleted, the beat-wave amplitude decayed because of the strong damping by the resonant electrons.

We investigated the use of high phase-velocity beat waves that might accelerate more energetic and, hence, less collisional current-carrying electrons. The linear Landau damping rate decreases exponentially with increasing phase velocity. From Eq.(2.F-8), one expects that the rate of action transfer at resonance will increase with decreasing $\text{Im}\epsilon$. However, accompanying the increased action transfer rate is the smaller resonance region in a nonuniform plasma, proportional to $\text{Im}\epsilon L$. In our uniform-plasma simulations with damping rates less than the frequency mismatch, $|\text{Im}\epsilon| < |\text{Re}\epsilon|$ and the reduced beat-wave dissipation decreases the action transfer rate. The action transfer then occurs over a longer period of time. The reduced electron damping invites competition with other damping mechanisms, e.g., scattering by ion modes [16]. The weaker action transfer rate may also encourage the possibility of stimulated scattering instabilities.

We performed a number of simulations with $\omega_3/k_3 v_e > 6$, all of which exhibited a turbulent spectrum of longitudinal modes and nonlinear scattering by ions that disrupted the beat-wave coupling. Furthermore, the ions typically acquired as much momentum or more than that deposited in the electrons, the latter which was often directed backwards relative to the higher frequency wave source! Simply stated, there were insufficient electrons at high velocities to adequately damp beat waves with phase velocities exceeding $6 v_e$. As a result, the beat wave found other ways to couple to the plasma, to the detriment of the current drive.

One can also excite beat waves with phase velocities resonant with the bulk of the electron distribution, $\omega_3/k_3 < 2 v_e$. However, here the electron damping can become so strong that $|\epsilon|$ is not small compared to unity; and the beat wave is a heavily damped quasi-mode. Equations (2.F-1), (2.F-2), and (2.F-8) suggest that the beat-wave coupling will proceed more slowly in this limit. Furthermore, the bulk electrons are much more collisional than those in an energetic tail. In consequence, the current-drive efficiency will be significantly reduced in the absence of any other nonlinear effects.

2.F.4 CURRENT-DRIVE EFFICIENCY

We first review an estimate of the beat-wave current-drive efficiency given earlier in [6]. The beat-wave coupling produces a flux of electrostatic waves in the interaction region. These beat waves are efficiently damped by resonant electrons over a relatively short distance going around the torus. The resonant electrons carry away the beat-wave momentum and thus drive a current. Bunches of electrons are produced in pulses because the pump waves are pulsed FEL sources. The electrons collisionally scatter and slow down on a much longer time scale, and the equilibrium adjusts to the current on the L/R time scale.

On averaging over many pulses, the beat-wave coupling delivers momentum and current to the resonant electrons. The momentum transfer is constrained by the Manley-Rowe relations and depends on the average power in the pump waves. However, the fraction of power delivered to the beat wave is given by the product of the quantum efficiency q_e

$$q_e = \frac{\omega_3}{\omega_1(1 + \omega_2\rho/\omega_1)} \quad ,$$

where ρ is the ratio of input action intensities, times the relative action transfer R_a calculated earlier, which depends nonlinearly on the pump wave peak intensities. Thus, the current drive depends on both the peak power to achieve a satisfactory value of R_a and the average power to drive current. The toroidal current-balance equation deduced from momentum conservation is

$$\frac{d}{dt} j_t = \frac{\Delta j_e}{\tau} - \nu_r j_t^r - \nu_b j_t^b + \frac{e^2 n_e}{m_e} E_t \quad (2.F-15)$$

in the nonrelativistic limit, where E_t is the toroidal induction electric field, $j_t = j_t^r + j_t^b$ the sum of fast resonant and bulk current contributions, ν_r is the collisional slowing-down rate of a fast electron $\nu_r = 3.9 \times 10^{-6} s^{-1} \epsilon_e (eV)^{-3/2} \times Z^2 n_i \ln \Lambda$, ν_b is the slowing-down rate for bulk electrons ($\nu_b \gg \nu_r$), n_i is the ion density, ϵ_e is the energy of the resonant electrons, $\ln \Lambda$ is the Coulomb logarithm, τ is the toroidal transit time $\tau = 2\pi R_0 / v_r$, and $v_r = \omega_3 / k_3$ is the resonant electron velocity. The electron current density

introduced on each toroidal transit of the electrons is Δj_e and is equal to $-e/m_e v_r$ times their momentum flux density, which is calculable from the Manley-Rowe relations, i.e.,

$$\Delta j_e = -(e/m_e v_r) R_a k_3 j_1^{\text{in}} \quad (2.F-16)$$

where j_1^{in} is the input action flux density and R_a is the relative action transfer. The expression in Eq. (2.F-16) assumes that the beat wave is completely damped by the resonant electrons over a toroidal transit.

At steady state, $dj_t/dt = E_t = j_t^b = 0$ and Eq. (2.F-15) gives $j_t = \Delta j_e / \nu_r \tau$. From the ratio of the toroidal current to the input power in the high frequency waves $\omega_1 j_1^{\text{in}} + \omega_2 j_1^{\text{in}} = \omega_1 j_1^{\text{in}} (1 + \omega_2 \rho / \omega_1)$, one deduces the beat-wave current-drive efficiency η_{bw} ,

$$\eta_{bw} = |j_t / \omega_1 j_1^{\text{in}} (1 + \omega_2 \rho / \omega_1)| = q_e R_a \frac{e}{m_e \nu_r} \frac{1}{2\pi R_0 \nu_r} \quad (2.F-17)$$

A careful relativistic treatment of the Fokker-Planck aspects of rf current drive has been given by Fisch [17,18]. Fisch calculates a current-drive efficiency η_{cd} for an electrostatic wave that is completely damped by resonant particles

$$\eta_{cd} = \frac{I_{cd}}{P_{cd}} = \frac{31.2}{2n\lambda} \left(\frac{\hat{J}}{\hat{P}} \right) \frac{1}{R_0(m) n_e (10^{14} \text{cm}^{-3})} \quad (2.F-18)$$

where $\hat{J}/\hat{P}$ is a function of the momentum of the current carrying particles. In the nonrelativistic limit with $v_{\parallel} \gg v_{\perp}$, $\hat{J}/\hat{P} = 8x^2/(5 + z_{\text{eff}})$ and $x^2 = 2\epsilon_e(\text{keV})/511$. The beat-wave current-drive efficiency is reduced by the quantum efficiency q_e and the relative action transfer ratio R_a so that

$$\eta_{bw} = q_e R_a \eta_{cd} = 0.8 \times 10^{-2} q_e R_a \epsilon_e(\text{keV}) \frac{1}{R_0(m) n_e (10^{14} \text{cm}^{-3})} \quad (2.F-19)$$

in the nonrelativistic limit and with $8/(5 + z_{\text{eff}}) = 1$ and $31.2/2n\lambda = 2$. There is good agreement between Eqs.(2.F-17) and (2.F-19). It is important to

recall Fisch's result that at higher resonant velocities relativistic effects cause the current-drive efficiency to saturate with increasing electron energy: $\hat{J}/\hat{P} = O(1)$ for $x = p/mc \geq O(1)$ and $p \ll p$.

2.F.5 DISCUSSION

In this section we discuss the application of beat-wave coupling to the MTX experiment and the scaling to an ETR plasma. Conditions are described in which all of the required tokamak current is driven by beat waves. For simplicity the particle simulations described here were unmagnetized and focused on the coupling to the beat wave and its absorption in the plasma. In the proposed application to experiments, there will be a strong background magnetic field. Consider the coupling of parallel or anti-parallel propagating transverse waves to a Langmuir wave directed along the magnetic field. The transverse normal modes in this case are right- and left-circularly polarized electromagnetic cyclotron waves with local cold-plasma dispersion relations

$$k_R = \pm \frac{\omega}{c} \left(1 - \frac{\omega_{pe}^2}{\omega(\omega - \Omega_e)}\right)^{1/2}, \quad k_L = \pm \frac{\omega}{c} \left(1 - \frac{\omega_{pe}^2}{\omega(\omega + \Omega_e)}\right)^{1/2} \quad (2.F-20)$$

where ω is the wave frequency and $\Omega_e = |eB_0/m_e c|$ is the electron cyclotron frequency. The right circularly polarized wave rotates in the electron sense and propagates for $0 \leq \omega \leq \Omega_e$ and $\omega \geq \Omega_2$, where $\Omega_2 = (\Omega_e/2)[(1 + 4\omega_{pe}^2/\Omega_e^2)^{1/2} + 1]$. The left-circularly polarized wave rotates in the ion sense and propagates for $\omega \geq \Omega_1$, where $\Omega_1 = (\Omega_e/2)[(1 + 4\omega_{pe}^2/\Omega_e^2)^{1/2} - 1]$.

Vacuum propagation characteristics are recovered from Eq. (2.F-20) for $\omega_{pe}^2 \ll \omega^2$, and the ray paths become straight. Cyclotron harmonic absorption by the electrons is much reduced for left-circularly polarized waves, and toroidal ray-tracing calculations in TORAY [19] indicate that refraction is insignificant for $\omega_{pe}^2/\Omega_e^2 \leq 1/2$ and $\omega > \Omega_e$. Absorption of right-circularly polarized waves can be minimized by raising the wave frequency to a higher cyclotron harmonic number and arranging the cyclotron resonance to fall outside the plasma or at its edge. The ray paths for the proposed beat-wave coupling to a Langmuir wave are shown in Fig. 2.F-4. The ray paths of the transverse waves are nearly parallel to the magnetic field near the magnetic axis and are obliquely incident at the plasma edge. The plasma normal modes

Figure 2.F-4. A toroidal sector showing ray paths for beat-wave coupling of O-modes in MTX calculated in the TORAY code. The ports are 60° apart, and there is negligible attenuation of the pump waves from cyclotron damping. The O modes become almost exclusively left circularly polarized in the interaction region.

are elliptically polarized extraordinary (X) and ordinary (O) modes. The X mode becomes the right-circularly polarized wave near the point of tangency with the magnetic field, and the O mode becomes left-circularly polarized. The time-averaged beating of a right circularly polarized wave with a left circularly polarized wave is zero: only waves with the same polarization will beat-wave couple in Eqs.(2.F-3) and (2.F-4). (In the particle simulations, which were unmagnetized, the electric field vectors of the transverse pump waves were linearly polarized parallel to one another to give maximal coupling.)

When oriented parallel to one another, the transverse wave wavenumbers vectorially subtract to give a small difference wavenumber. For opposed propagation, the beat wavenumber is the sum of the wavenumber magnitudes. Hence, opposed propagation leads to relatively smaller beat-wave phase velocities, $\omega_{pe}/2\omega_1 < \omega_{pe}/(k_1 + k_2)c < 1/\sqrt{3}$ and a parallel orientation produces higher beat-wave phase velocities, $1/\sqrt{3} < \omega_{pe}/(k_1 - k_2)c < 1$. Thus, a range of beat-wave phase velocities can be selected extending nearly to the speed of light. This is important, because we wish to match the beat-wave phase velocity to 2 to 4 v_e . If the frequency ω_1 is fixed by the technology available, then matching the beat-wave phase velocity to 2 to 4 v_e determines a relation between the local plasma density and electron temperature compatible with the beat-wave interaction. This constraint relation does not prove to be too restrictive for the MTX and ETR examples described in the following. Frequency-tunable high-frequency wave sources would allow greater flexibility in the density and temperature operating conditions, and would be an advantage.

A beat-wave current-drive scenario has been calculated for MTX parameters [2]. Consider two pump waves with frequencies equal to 250 GHz and 175 GHz, a peak plasma density of 0.7 to $0.8 \times 10^{14} \text{ cm}^{-3}$, and a magnetic field on axis equal to 3.9 T. This corresponds to $\omega_1 = 2.3\Omega_e(0)$, $\omega_2 = 1.6\Omega_e(0)$, and $\omega_{pe}^2(0)/\Omega_e^2(0) = 1/2$. Two ordinary modes could be injected directed at one another from two adjacent ports 60° apart as pictured in Fig. 2.F-4. In the beat-wave interaction region near the magnetic axis, the two O modes are almost completely left circularly polarized waves. The ratio of the beat-wave phase velocity to the electron thermal speed $\omega_3/k_3v_e = 3.5$ for $T_e = 1.5 \text{ keV}$. The TORAY code indicates that the pump waves will be negligibly absorbed by direct resonant interaction with the electrons. The beat-wave current-drive

can be calculated from Eq. (2.F-19). In MTX, $u_1/c = 0.015$ for 8 GW peak power and a 3cm x 8cm illuminated cross-section. The wavelength $\lambda_1 = 0.1$ cm, and for a density scale length along the ray paths we assume $L_n \sim 2a = 30$ cm. Then $J_1 = 5$ and $R_a = 0.8$ for $\rho = 0.1$ and $R_a = 0.4$ for $\rho = 0.01$, where ρ is the ratio of the incident action fluxes in the pump waves. For $\rho < 0.1$ the quantum efficiency is approximately $q_e = 1/3$, and for $R_0 = 0.6$ m the current-drive efficiency $\eta_b = 0.06 R_a$ amps/watt. In terms of $\eta_b = n_e IR_0/P = n_e R_0 \eta_{bw}$, $\eta_{bw} = 0.03 R_a A \cdot 10^{20} \text{ m}^{-2}/\text{W}$. The 8 GW peak power corresponds to 2 MW average power, which leads to 120 R_a kA of beat-wave current. For reference, the value of the rotational transform q at the plasma surface is given by [2,20]

$$q = \frac{B_z a}{B_\phi R} = 5 \frac{B_z (G)}{I(\text{amps})} \frac{a^2(\text{cm})}{R_0(\text{cm})},$$

where B_z and B_ϕ are the toroidal and poloidal magnetic fields. A current of < 100 kA will maintain $q > 7$ at the edge of the plasma. The Murakami empirical limit on density for Ohmic and rf-heated discharges is [2,20]

$$n_0(\text{m}^{-3}) = 10^{20} B_z(T)/qR(\text{m}),$$

which is approximately $2 \times 10^{14} \text{ cm}^{-3}$ for MTX parameters. Porkolab reports that Alcator-C was able to operate in the regime of density, temperature, and current described in this MTX example [20].

In assessing the possibility of beat-wave current drive in MTX, there are stringent requirements on the pump waves. Two fairly intense wave sources are required. Recall that the action transfer ratio R_a is quite sensitive to the peak power in the two sources. Furthermore, we have considered a particular choice of polarizations, G mode, which is dominantly left-circularly polarized. However, the Livermore FEL experiments reported in [1] used a linear wiggler magnet set and produced dominantly linearly polarized microwaves, which are a mix of right- and left-circularly polarized waves. If half the power is lost to the beat-wave coupling, because it has the wrong polarization, the action transfer ratio R_a and the average wave power will each drop by approximately one-half, resulting in a four-fold reduction in the current driven. The other half of the average power in the pump waves is accessible to beat-wave current drive if both transverse-wave pumps have the

same mix of polarizations, so that there are two sets of beat-wave interactions occurring simultaneously. However, the reduction in R_a from its value when all the pump-wave power is in one polarization state persists. Furthermore, the magnetic field and peak density must be slightly reduced for the same pump frequencies from their values in the example described for MTX. This is necessary in order that the X-mode components can propagate into the heart of the plasma. However, this is somewhat counter-productive, because it reduces the quantum efficiency.

We next present a possible scenario for beat-wave current drive in an ETR plasma. Assume first that the pump-wave frequencies, the peak powers, the equilibrium magnetic field, and the density remain unchanged from the MTX example. However, let the electron temperature be $T_e = 25$ keV on axis and assume a five-fold increase in the microwave pulse rate from that in MTX so that 10 MW average power is provided in the first pump wave. Consider a major radius $R_0 = 3.0$ m and a minor radius $a = 0.8$ m. To match the beat-wave phase velocity to $4 v_e \sim 0.9 c$ requires a parallel orientation of the pump waves.

In this orientation and with a scale length $L_n = 2a$, $\bar{J}_1 = 1$. If the ratio of input actions is as large as $\rho = 0.3$, then we deduce that $R_a = 0.4$ from Fig. 2.F-1(a) and $q_e = 1/4$. The beat-wave current-drive efficiency determined by Eq. (2.F-18) and with $\hat{J}/\hat{P}$ given in [17] is $\eta_{bw} = 0.1$ A/W and $\hat{\eta}_{bw} = 0.2A \cdot 10^{20} m^{-2}/W$. The total driven current is 1 MA, which would give $q = 5$ at the edge of a plasma with a circular cross-section. This should prove adequate for sustaining the plasma. Elongation of the plasma will allow more current for the same q at the edge of the plasma and require more average power in the beat wave to drive the current. A better current-drive efficiency could be attained at higher peak pump-wave powers so that R_a could be improved to a value closer to unity. With an upgrade of the FEL to 24 GW peak power (corresponding to electron beam energies > 20 MeV and currents > 3 kA), $\bar{J}_1 = 3$ and $R_a > 0.8$ for $\rho = 0.1$. This would increase $\hat{\eta}_{bw}$ to $\hat{\eta}_{bw} = 0.4 A \cdot 10^{20} m^{-2}/W$.

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APPENDIX 2.6

HARMONIC OVERLAP IN ELECTRON-CYCLOTRON CURRENT DRIVE AT HIGH T_e

The relativistic wave-particle resonance condition imposes a harmonic-overlap constraint on electron-cyclotron current drive. Specifically, for fundamental cyclotron resonance the current-drive efficiency cannot be increased by raising the minimum resonant energy ϵ_1 above about 115 keV. A limiting efficiency corresponding to an even lower ϵ_1 follows for microwave beams with realistic width and angular divergence.

The tokamak-fusion-reactor concept would be significantly enhanced by a workable scheme for non-inductive current drive. Injection of electron-cyclotron waves is one of the candidate schemes and has some attractive features: simple launching structure away from the plasma edge, modest port-size requirements, and localizable power deposition.

Electron-cyclotron current-drive efficiency is highest at low density and high electron temperature T_e , according to non-relativistic analysis [1-3]. At reactor-relevant values of T_e (> 10 keV), relativistic effects must be included in the analysis of electron-cyclotron heating and current drive. Fisch found that the relativistic electron dynamics reduces current-drive efficiency to half its non-relativistic value at an electron energy of about 170 keV and leads to a saturation of the efficiency at 800 keV [4]. In this note, we consider another relativistic effect, the modification of the wave-particle resonance condition, and the consequences of the relativistic resonance condition on current-drive efficiency. For reactor-relevant values of T_e we show that, at a given spatial position, different groups of electrons can resonate with the wave at different harmonics of the cyclotron frequency. This competition between harmonics significantly limits current-drive efficiency and is important at lower electron energies than the effects considered by Fisch [4].

We treat an arbitrary region, small compared to the dimensions of the plasma, in which the rest-mass electron-cyclotron frequency is Ω , the wave frequency is ω , and the parallel component of the index of refraction is $N_{\parallel}$.

The relativistic resonance condition can be expressed in terms of the parallel and perpendicular components of dimensionless momentum $u_{\parallel} = p_{\parallel}/m_e c$ and $u_{\perp} = p_{\perp}/m_e c$ by

$$\gamma - N_{\parallel} u_{\parallel} = \ell \Omega / \omega \quad (2.6-1)$$

where $\gamma^2 = 1 + u_{\perp}^2 + u_{\parallel}^2$, and ℓ is the cyclotron-harmonic number. For $|N_{\parallel}| < 1$, the locus of points satisfying Eq. (2.6-1) is an ellipse with center at $u_{\perp} = 0$ and $u_{\parallel} = (\ell \Omega / \omega) N_{\parallel} / (1 - N_{\parallel}^2)$. The semi-minor radius along $u_{\perp}$ is

$$\Delta u_{\perp} = \left[\frac{(\ell \Omega / \omega)^2}{(1 - N_{\parallel}^2)} - 1 \right]^{1/2}.$$

The semi-major radius along $u_{\parallel}$ is

$$\Delta u_{\parallel} = \frac{\Delta u_{\perp}}{(1 - N_{\parallel}^2)^{1/2}}.$$

There are no resonant electrons at all for harmonic ℓ if $(\ell \Omega / \omega)^2 < 1 - N_{\parallel}^2$. Note that the resonance ellipses for every harmonic number ℓ are determined by just two parameters, $N_{\parallel}$ and Ω / ω . Note also that, for $\ell \Omega / \omega < 1$, the resonance ellipse is entirely of one sign in $u_{\parallel}$, whereas for $\ell \Omega / \omega > 1$, the ellipse straddles the $u_{\perp}$ axis.

It is convenient to parametrize the resonance ellipses by means of their intercepts with the $u_{\parallel}$ -axis. The more important intercept is the one at the minimum resonant energy $\epsilon_{\ell} = m_e c^2 (\gamma_{\ell} - 1)$, where

$$\gamma_{\ell} = \frac{(\ell \Omega / \omega) - |N_{\parallel}| [(\ell \Omega / \omega)^2 - 1 + N_{\parallel}^2]^{1/2}}{1 - N_{\parallel}^2} \quad (2.6-6)$$

In the following, we consider overlap of harmonics $\ell = 1$ and $\ell = 2$, for which the constraints on current-drive efficiency are least restrictive.

For the moment, we ignore the $\ell = 2$ harmonic and restrict attention to $\Omega < \omega$. At the moderately relativistic energies of interest for a tokamak reactor, the contribution of individual resonant electrons to current-drive

efficiency is an increasing function of energy [1,4]. Hence, to maximize current-drive efficiency, we would choose γ_1 as high as possible, consistent with the existence of a sufficient number of resonant electrons to absorb the wave on a single pass (in order to avoid canceling currents from absorption on high- and low-field sides of resonance). It follows from Eq. (2.G-2) that we can increase γ_1 while maintaining a non-zero population of resonant electrons by increasing $|N_{\parallel}|$ and decreasing Ω/ω . The solid contours of γ_1 in Fig. 2.G-1 illustrate this fact. Nevertheless, increasing γ_1 decreases the absorption because of the falloff of the distribution function with energy. The requirement of single-pass absorption translates approximately to the condition that the (resonant-energy-dependent) absorption coefficient times the minor radius a be greater than or of order unity; for a Maxwellian distribution (weak heating) in a reactor-grade plasma, we estimate, taking the fundamental ordinary-mode absorption coefficient from Bornatici et al. [5], that $\epsilon_1 \approx \eta T_e$, where $\eta \approx \ln[(2\pi T_e/mc^2)^{1/2} (a/\lambda) (\omega_p^2/\Omega^2) g] \approx 5$. Here, λ is the wavelength, ω_p is the electron plasma frequency, and g is a propagation-angle-dependent function of order unity. For heating strong enough to form a significant tail in the electron distribution, η is likely to be appreciably larger. We adapt the rough guideline $\eta \approx 6$.

We now recognize the existence of electrons resonant with the $\ell = 2$ harmonic. The minimum energy of those electrons, $\epsilon_2 = m_e c^2 (\gamma_2 - 1)$, is indicated by the dashed contours of γ_2 in Fig. 2.G-1. To increase γ_1 , we see that we are forced to decrease γ_2 , which increases the number of electrons resonant at $\ell = 2$ and hence the power absorbed. For $0.5 < \Omega/\omega < 1$, the resonance ellipse for $\ell = 2$ crosses the line $u_{\parallel} = 0$, with the lowest-energy resonant electrons having $u_{\parallel}$ opposite in sign from the $u_{\parallel}$ of electrons resonant at $\ell = 1$. Therefore, for $\gamma_2 < \gamma_1$ and $0.5 < \Omega/\omega < 1$, the $\ell = 2$ harmonic absorbs power comparable to or greater than that of the $\ell = 1$ harmonic while contributing current that is always small [6] and is usually in the wrong direction. (For $\Omega/\omega < 0.5$ the problem is shifted to higher harmonics.)

The preceding discussion implies that, to maximize current drive, we want to operate with γ_1 as large as possible consistent with γ_2 being still larger. Examination of Fig. 2.G-1 shows that, at any given γ_1 , the largest γ_2 occurs on the curve $(\Omega/\omega)^2 = 1 - N_{\parallel}^2$. This fact can be understood easily from inspection of Eq. (2.G-2) for $u_{\parallel} = 0$ and $\ell = 2$:

Figure 2.G-1. Dependence on $|N_{||}|$ and Ω/ω of the minimum resonant energies ϵ_l for the $l = 1$ and $l = 2$ harmonics, shown by means of contour plots of $\gamma_l = 1 + \epsilon_l/m_e c^2$. Solid contours show γ_1 , and dashed contours show γ_2 . Each γ_l -contour ends on the curve $(\Omega/\omega)^2 = 1 - N_{||}^2$ at $\Omega/\omega = 1/\gamma_l$. To the left of that curve, there are no electrons resonant for $l = 1$.

$$\gamma_2 = (1 + u_{\parallel}^2)^{1/2} = N_{\parallel} u_{\parallel} + \Omega/\omega \quad (2.G-3)$$

For $\Omega/\omega > 0.5$ and $N_{\parallel} u_{\parallel} < 0$, it is clear from Eq. (2.G-3) that raising Ω/ω or lowering $|N_{\parallel}|$ increases γ_2 . We thus find the largest γ_2 for a given γ_1 at the highest Ω/ω or lowest $|N_{\parallel}|$ for which the $\ell=1$ resonance exists, i.e., along $(\Omega/\omega)^2 = 1 - N_{\parallel}^2$. On this curve, γ_1 reaches its maximum value ω/Ω , either as a function of $N_{\parallel}$ at fixed Ω/ω or as a function of Ω/ω at fixed $N_{\parallel}$ [6]. The maximum γ_1 consistent with $\gamma_2 \geq \gamma_1$ is thus found by solving $\gamma_2 = \gamma_1 = \omega/\Omega$, which leads to the limit $\gamma_1 < (3/2)^{1/2} \approx 1.22$ [or, equivalently, the requirement $\Omega/\omega > (2/3)^{1/2}$. For $\gamma_1 < (3/2)^{1/2}$ there are values of $|N_{\parallel}|$ and Ω/ω for which $\gamma_2 > \gamma_1$, while for $\gamma_1 > (3/2)^{1/2}$ there are not. The limit on the minimum resonant energy for the $\ell = 1$ harmonic is $\epsilon_1 < 115$ keV. Therefore, the current-drive efficiency cannot be increased above that achieved when the minimum energy of resonant electrons is about 115 keV. Also, the guideline $\epsilon_1 \approx 6T_e$, which leads to the highest achievable current-drive efficiency for each value of $T_e < 115 \text{ keV}/\eta \approx 20 \text{ keV}$, cannot be extended to $T_e > 20 \text{ keV}$, and efficiency cannot be increased above that at $T_e \approx 20 \text{ keV}$.

We next generalize our analytic estimate of the constraint imposed by harmonic overlap to an arbitrary pair of harmonics ℓ and $\ell + 1$. We consider current drive by selective heating of high-energy electrons at cyclotron harmonic ℓ . The largest $\gamma_{\ell+1}$ for any given γ_{ℓ} is found for $\gamma_{\ell} = \omega/\ell\Omega$ under the condition $(\ell\Omega/\omega)^2 = 1 - N_{\parallel}^2$, which we insert into Eq. (2.G-2). Equating the resulting expressions for γ_{ℓ} and $\gamma_{\ell+1}$ yields the harmonic-overlap constraint in the form of an upper limit on the minimum resonant energy given by

$$\gamma_{\ell} < \left(1 + \frac{1}{2\ell}\right)^{1/2}.$$

An approximate limit for ϵ_{ℓ} is $m_e c^2/4\ell$, which shows how the harmonic-overlap constraint becomes more restrictive as ℓ increases.

The sensitivity of γ_1 to $|N_{\parallel}|$ and Ω/ω , which is illustrated in Fig. 2.G-1, suggests difficulties in realizing the maximum current-drive efficiency consistent with the constraint $\gamma_1 < 1.22$. A practical scheme for propagating power into the plasma involves a finite spatial extent for the beam (and thus a finite range of Ω/ω) and a finite range of $N_{\parallel}$. Ranges of only a few percent in either parameter imply significant variation of γ_1 .

Since the maximum value of $\gamma_1 = 1.22$ with $\gamma_2 > \gamma_1$ is achieved only for one specific pair of values for $N_{\parallel}$ and Ω/ω , it is clear that the current-drive efficiency averaged over the extent of the beam is generally not as high as the maximum possible local value.

Electron-cyclotron current drive is subject to a harmonic-overlap constraint, which follows from the relativistic wave-particle resonance condition. This conclusion arose from analysis of the ranges $\ell < \omega/\Omega < \ell+1$ for the cases $\ell \geq 1$. An alternative current-drive scheme would use $\omega/\Omega < 1$ (absorption on the high-field side of the fundamental cyclotron resonance). If it is possible to produce a wave in the plasma with $|N_{\parallel}| \gg 1$, then for $\omega/\Omega \ll 1$ current drive becomes quite efficient. Unfortunately, conventional launching and propagation of extraordinary or ordinary modes generally leads to the limit $|N_{\parallel}| < 1$. Furthermore, schemes with $|N_{\parallel}| < 1$ and $\omega/\Omega < 1$ suffer from very low efficiency because the resonance ellipse crosses the line $u_{\parallel}=0$ and electron currents are driven in opposite directions, resulting in low net currents [6,7]. In addition, at high power the electron distribution is enhanced in the trapped part of velocity space, a process that drives current opposing the Fisch-Boozer current [6,8,9]. We therefore conclude that current-drive schemes utilizing absorption of electron-cyclotron waves as envisioned in Ref. [2] have efficiencies that do not scale favorably as the minimum resonant energy ϵ_1 increases above 115 keV. In practice, this limit on ϵ_1 constrains current-drive efficiency, even for an ideal microwave beam, to that achieved at $T_e \approx 20$ keV.

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APPENDIX 2.H

COMBINED FAST WAVE/LOWER HYBRID CURRENT DRIVE

2.H.1 INTRODUCTION AND SUMMARY

The main limitation of lower hybrid current drive (LHCD) is its lack of penetration into a hot, high-density plasma. For this reason the TIBER II baseline current drive scenario calls for high-energy neutral beam current drive (NBCD) near the magnetic axis, combined with LHCD near the surface. The drawbacks of this approach are related to the NBCD apparatus, which allows tritium and neutron streaming back through the beamline. Extensive shielding and source protection may be needed, and the large size and location of the beam line could hamper access and maintenance of the tokamak. In addition, MeV negative ion beam technology will require considerable development.

An alternative current driver is available in principle. This technique uses the fast wave polarization to achieve current generation near the magnetic axis. In theory the fast wave has a large radial group velocity, and at low frequencies and low $N_{\parallel}$ it is possible to achieve wave damping at high temperature and density. Fast wave ion heating has an extensive data base: the well-known techniques of minority heating and second harmonic cyclotron heating have been frequently demonstrated in the ion cyclotron range of frequencies (ICRF) with the fast wave on tokamaks such as PLT and JET. For electron heating and current drive the experimental results to date are inconclusive and the fast wave method is, therefore, still speculative. Table 2.H-1 summarizes the experimental situation. Briefly, high harmonic (GHz) fast wave experiments are hampered by problems of mode conversion to slow waves and antenna coupling, while conditions in the ICRF have been unfavorable to driving currents. A definitive experiment to demonstrate fast wave current drive in the ICRF has yet to be carried out.

The fast wave current drive (FWCD) method is an innovation which could provide numerous attractions. The RF equipment available today could be used for TIBER II with very little development effort. The transmission line could be coaxial with bends and windows in order to contain tritium and neutron radiation within the toroidal chamber. By slightly tuning the frequency, f , and properly phasing the antenna, the same system could switch between ion

Table 2.H-1. Experiments with fast wave current drive.

Experiment	Status	Comment
ACT-1 (Princeton, 1985)	Successful, 30A	
JIPP T-11 U (Nagoya, 1986)	Partial success, 15 kA (LH range) 2 kA (9th harmonic)	
Irvin Torus (U.C.I, 1986)	2 A	Antenna optimization
CCT (UCLA, 1986)	2 kA	Loop antennas
Versator 11 (MIT, 1986)	Partial success, 8 kA	Slotted waveguide antenna
PLT (PPPL, 1986)	Unsuccessful	Coupling to slow wave
JFT-2M (JAERI, 1987)	Partial success, 30 kA	

heating (for ignition heating) and FWCD (electron heating) if the ICRF is chosen. Also, with ICRF three distinct methods can be used to drive current. The simplest physics, direct electron Landau damping and transit time magnetic pumping of the launched waves, can provide current via the same mechanism which acts for LHCD; this first technique has received the most attention [1-4] and will be the subject of our present investigation. The second FWCD technique [5,6] involves mode conversion of a launched fast wave into an ion Bernstein mode (provided $N_{\parallel}$ is sufficiently low, $\lesssim 2$) which is strongly Landau damped and should generate current density across the mode conversion surface. This technique may be tested first on JET. The third ICRF alternative would employ minority ion heating to establish a circulating high energy ion beam in situ, without an external neutral beam line [7-9]. Finally, we point out that

a wide range of frequencies is available for FWCD. If, for example, ICRF is impractical it may still be possible to generate centrally peaked currents even with FWCD in the GHz range.

In the remainder of this section we present results of combined fast wave/lower hybrid current drive for a variety of physics assumptions. In all cases we calculate self-consistent, noninductively generated two-dimensional MHD equilibria. The main conclusions are as follows:

- For high-frequency FWCD an extensive survey was performed over a large range of geometries and plasma parameters. For $\beta = 0.06$ and $T_E = 14.3$ keV, we expect the current drive figure of merit to be $n_e R_O I / P_{CD} = \gamma = 0.32$ ($10^{20} \text{ AW}^{-1} \text{ m}^{-2}$). With this combination of drivers one would require $P_{CD} = 106$ MW for $\langle n_e \rangle = 1.13 \times 10^{20}$.
- Using the ICRF fast wave one finds $P_{CD} = 114$ MW is necessary to achieve equilibrium.
- If a bootstrap current is present [10] much of the toroidal current is provided due to classical effects. With a flat density profile we find $P_{CD} = 94$ MW is still needed, with the ICRF driver.
- A more peaked density profile augments the bootstrap effect, and we find $P_{CD} \leq 60$ MW only is needed for equilibrium. The ICRF fast wave provides a seed current near the magnetic axis, where plasma gradients are small. Little or no power is needed for higher frequency (lower hybrid) waves, depending on details of the profiles.

2.H.2 HIGH-FREQUENCY FWCD

An extensive series of rf-generated MHD equilibria has been calculated with multiple-wave power spectra. By properly selecting mode frequency, polarization, and power it is possible to generate axisymmetric toroidal currents which are consistent with MHD equilibrium. The density, N_i temperature, and, hence, pressure profiles are chosen ad hoc to be reasonable, simple, functions of poloidal flux, ψ , and the total current drive power, P_{CD} , absorbed in the plasma is computed for plasma beta near the Troyon limit. The

principal result is the ratio $\gamma \sim \langle n_e \rangle R_0 I / P_{CD}$, where $\langle n_e \rangle$ is the volume-averaged electron density, R_0 is the major radius, and I is the toroidal current.

For this study the safety factor is chosen at the magnetic axis and vacuum boundary as $q_a = 1.0$ and $q_b = 2.0$, respectively; $q(\psi)$ is monotonic. We surveyed a wide range of tokamak parameters: aspect ratio, $3.6 \leq A \leq 8.0$; elongation/triangularity, $(1.0, 0.0) - (2.0, 0.75)$; size $3.0 \text{ m} \leq R_0 \leq 7.0 \text{ m}$; density, $0.7 \times 10^{20} \text{ m}^{-3} \leq \langle n_e \rangle \leq 2.8 \times 10^{20} \text{ m}^{-3}$; temperature, $10 \text{ keV} \leq \langle T_e \rangle \leq 30 \text{ keV}$, and $11 \text{ keV} \leq \langle T_e \rangle \leq 25 \text{ keV}$; effective ion charge, $1.0 \leq Z_{eff} \leq 2.0$; beta, $0.02 \leq \beta \leq 0.08$; and temperature and density profile widths.

This investigation described in this subsection considered only the combination of lower hybrid and high-frequency fast waves, the former providing current density in the low density surface plasma region and the latter waves (at $\omega \gg \omega_D$) providing the bulk of the current in the high density plasma interior. The waves originate at the outboard midplane and are followed by a ray tracing routine which evolves the wave vector along the ray path, including the change in $N_{||}$. The calculation assumes weak damping (linear Landau damping) with single pass wave absorption, and rf current generation is computed according to the prescription given in Ref. [11]. The calculation assumes steady state current drive in the absence of an electric field. There is no suprathermal electron radial transport and no $N_{||}$ upshift due to effects such as wave scattering off edge density fluctuations.

We define

$$\gamma = (\langle n_e \rangle / 10^{20} \text{ m}^{-3}) (R_0 / \text{m}) (I / \text{MA}) / (P_{CD} / \text{MW})$$

and found a general expression which characterizes all our results, when $Z_{eff} = 1.5$:

$$\gamma / (\langle T_e \rangle^{0.77}) = 0.034 + 0.196 \beta, \quad (2.H-1)$$

where $\langle T_e \rangle$ is the volume-averaged electron temperature in keV. This formula is unreliable for $\langle T_e \rangle \lesssim 10 \text{ keV}$, and it tends to underestimate the value of γ obtained in experiments at $\pm 1 \text{ keV}$. Our detailed calculations produced γ values within $\pm 10\%$ of the curve given by this formula (see Fig. 2.H-1). For TIBER, a typical current drive efficiency, for $Z_{eff} = 1.5$, $\langle T_e \rangle = 14 \text{ keV}$ and

Figure 2.H-1. Current drive figure of merit calculated for high frequency fast wave current drive. All equilibria are in first stability regime.

$\beta = 0.06$, is $\gamma = 0.32$. Our detailed calculations usually required lower hybrid waves at a few gigahertz carrying about one per cent of the total power and fast waves ranging from ~ 0.4 GHz to ~ 3 GHz with the bulk of the power.

When $Z_{\text{eff}} \neq 1.5$ the value of γ changes in a ratio approximately proportional to

$$0.08 \left\{ \frac{32}{(5+Z)} + 2 + \frac{12(6+Z)}{(5+Z)(3+Z)} + \frac{3.76}{Z} \right\} \quad (2.H-2)$$

Thus $\gamma(Z_{\text{eff}} = 1.5)$ is smaller than $\gamma(Z_{\text{eff}} = 1.0)$ in the ratio 0.85. Our formulas are most accurate for $Z_{\text{eff}} \lesssim 2$.

A study of γ dependence on T_i , pressure profile, and temperature and density profiles showed a relatively small sensitivity to these variations.

Our work ignored a few physics issues which might affect the value of γ to be achieved in a reactor. At high temperatures the fast wave is expected to be damped by transit-time magnetic pumping (TTMP). If included in the current drive calculations, TTMP would increase γ , perhaps by $\sim 20\%$. Also, our theory assumes linear Landau damping so a broad, raised plateau in $v_{\parallel}$ does not occur in the electron distribution function. A broad plateau, extending to $v_{\parallel} \rightarrow c$ could increase γ , especially for $\langle T_e \rangle \lesssim 10$ keV. Likewise, in this calculation we have ignored any possible benefit gained from the presence of a bootstrap current. One shortcoming of this work is the disregard of high-energy fusion alpha particle damping of the fast wave. If this damping mechanism competes strongly with electron damping it could motivate further studies of alternative wave drivers for the plasma interior. Candidates would include the low frequency fast wave, described below.

A more detailed account of this work is available in Ref. [1].

2.H.3 ICRF CURRENT DRIVE SCENARIO

Fast wave current drive in the ion cyclotron range of frequencies (ICRF) was investigated by Perkins [12] for the FED-A device. In this scheme the toroidally unidirectional wave is absorbed by electrons via Landau damping and TTMP processes, in which the resonant electron parallel speed $v_{\parallel}$ satisfies $v_{\parallel} = \omega/k_{\parallel}$, where ω is the wave frequency and $k_{\parallel}$ is the parallel wavenumber. These resonant electrons then carry the plasma current, which is dissipated through collisional processes with the thermal particles. It follows that accelerating higher energy resonant electrons results in higher current drive

efficiency, implying that for fixed ω , a low $k_{\parallel}$ wave is desirable. To facilitate coupling to the plasma core, a low $k_{\parallel}$ spectrum is also required. However, the electron damping decrement γ_e scales as

$$\gamma_e = 0.25\sqrt{\pi}\omega\beta_e\epsilon^{1/2}e^{-\epsilon}/|v_g|$$

where $\epsilon = 1/2(\omega/k_{\parallel}v_e)^2$, v_g is the wave group velocity, v_e is the electron thermal speed, and β_e is the electron beta. Note that as $k_{\parallel}$ becomes smaller, the damping gets weaker, and competition from other resonance processes involving ions and alpha particles may become dominant. Also wall losses may play a significant role as the wave reflects off the first wall.

One way to avoid ion cyclotron damping with the fuel ions is to design a reactor with an aspect ratio $R/a > 5$ such that the $\omega = 2\omega_{CD}$ and $\omega = 2\omega_{CT}$ resonance surfaces can be placed outside the plasma. Since the average alpha energy is in the MeV range, residual damping from these particles may compete with electron absorption. Higher harmonic cyclotron absorption of impurity ions may also be a factor.

Using FED-A parameters with $R = 4.2$ m, $a = 0.92$ m, $n_0 = 2 \times 10^{14} \text{ cm}^{-3}$, $T_{e0} = 15$ keV and $k_{\perp} = 2.6 \text{ m}^{-1}$, Perkins [12] concluded that it takes 20 radial passes for the wave energy to get totally absorbed through electron TTMP. Damping of the wave by alphas was presumed to be weak, since $\omega = 2\omega_{Ca} = 2\omega_{CD}$ resonance is now located on the outboard edge of the plasma.

Requiring the aspect ratio $R/a > 5$ for avoiding ion and alpha absorption is an extremely stringent design criterion to fulfill. We now propose to relax this condition and place the $2\omega_{CD}$ resonance on the low field edge of the plasma, with the $2\omega_{CT}$ surface on the high field side of the magnetic axis inside the plasma. One scheme is to launch the fast wave from the top or bottom of the Dee-shaped plasma such that the energy is focused towards the axis and, by refractive effect, completely away from the $2\omega_{CT}$ surface. Such a wave trajectory can indeed be obtained [13], as shown in Fig. 2.H-2 for an FED-A device, using the same parameters cited earlier [12]. Here the incident toroidal wavenumber is $k_{\phi 0} = 2.6 \text{ m}^{-1}$ and the initial poloidal wavenumber is $k_{\theta 0} = 0$. The wave actually bounces a few times radially within a relatively narrow vertical strip about the magnetic axis, and does not cross the $2\omega_{CT}$ surface at all. In this case, using a relatively low $k_{\parallel}$, it will take at

Figure 2.H-2. Fastwave (ICRF) trajectory in FED-A with top launch.

least several bounces for the wave to be absorbed, but the resultant current drive efficiency should be attractively high. It should be noted that the ray trajectory is sensitive to the density profile, the shape of the plasma flux surfaces, and the incident $k_{\parallel}$ spectrum, which may necessitate an intricate launcher design.

For the TIBER II reactor, the plasma is highly elongated and the electron beta is also substantially higher than that of FED-A. Thus the above CD configuration may be more difficult to achieve. An alternative and more straightforward approach is to launch a rather high $k_{\parallel}$ wave from the outboard edge along the equatorial plane such that maximum electron damping occurs on the initial transit of the wave energy through the plasma core, before the $2\omega_{CT}$ resonance is encountered. However, more RF power may be required to drive the same amount of current since the resonant electron energy is now closer to the thermal range. In the following subsections, this scheme will be studied, while the top launch option will be left for future investigation.

2.H.3 ICRF CURRENT DRIVE WITHOUT BOOTSTRAP

The goal of this calculation is to identify the wave frequency and power spectrum needed to duplicate the TIBER II equilibrium. For this purpose we selected a target equilibrium (Figs. 2.H-3a and 2.H-3b) in the TIBER II geometry: $A = 3.6$, $\kappa = 2.4$, $\delta = 0.4$, and $R_0 = 3.0$ m. The pressure profile is $p(\psi) = p_0 \tilde{\psi}^{1.4}$, where $\tilde{\psi}$ is the normalized flux function. We chose an initial diamagnetism function $F(\psi) = RB_t$ which results in $\beta = 0.050$ with $I = 10.1$ MA and $q_a = 1.0$ and $q_b = 3.38$ (the axis and boundary safety factors), when $B_0 = 5.55$ T. In the present subsection we report results for the early TIBER II baseline (flat density): $T_{eo} = 30$ keV, $n_{eo} = 1.43 \times 10^{20} \text{ m}^{-3}$, $\langle T_e \rangle \approx 14.3$ keV, and $\langle n_e \rangle = 1.1 \times 10^{20} \text{ m}^{-3}$. We assume $T_{i0} = 25$ keV and $Z_{eff} = 1.5$, with a DT ion mix.

For current generation at the surface, LHCD is appropriate, and we investigated a variety of $N_{\parallel}$ modes and launch positions for the frequency $f = 3.7$ GHz. Figure 2.H-4 displays the shallow ray trajectories of slow waves with $N_{\parallel} = 2.3 - 1.9$, starting below the midplane. The function $G = \langle j_{\parallel} B \rangle / \langle B^2 \rangle$ is a measure of the parallel current density averaged over a flux surface. The launch position illustrated yields the deepest current generation with LHCD.

PSI CONTOURS

Figure 2.H-3. (a) Location of harmonic ion resonance for $B_0 = 5.55$ T and low frequency (ICRF) fast wave at $f = 66$ MHz. (b) Baseline TIBER II equilibrium with $b = 0.05$ and $I = 10$ MA.

T2LH92 *** TEST CASE FOR HI FREQ FAST WAVE *** SEVEN RAYS
 T2LH92 TIME-12:52:32 DATE-06/06/87 MACHINE-D SUFFIX-A

Figure 2.H-4. LHCD at 3.7 GHz; note only surface current drive is achieved at high density ($n_{eo} = 1.4 \times 10^{20} \text{ m}^{-3}$, $T_{eo} = 30 \text{ keV}$).

For current drive at the axis, the ICRF fast wave frequency is adjusted to permit maximum electron damping with minimal ion losses. Deuterium resonances are particularly dangerous, as hot alpha damping will be extremely strong at these locations. We remove all deuterium and alpha resonances by selecting $f = 66$ MHz. With $B_0 = 5.55$ T the most serious cold ion loss is the second tritium harmonic, which lies inboard of the magnetic axis at $R = 2.57$ m (see Fig. 2.H-3(a)). Next, a series of launch positions was tested. The best location, just below the midplane, is illustrated in Fig. 2.H-5. If other launch locations are chosen, it is difficult to control the evolution of $N_{||}$; either $N_{||}$ increases too quickly, leading to off-axis current drive, or $N_{||}$ drops and even changes sign, leading to reversed current density. As shown, the fast wave travels directly through the central region. By the time the ray (with $N_{||} = 3.2$) reaches the second tritium harmonic layer 76% of its power has been dissipated on electrons. Then, at $R = 2.57$ m, the remaining power is lost to the tritium; see Fig. 2.H-6. Note, for $f = 66$ MHz we would expect loop antennas to have a narrow poloidal extent ~ 35 cm, and it appears possible to model the ICRF as radiating from this single point (although the code is equipped to include a finite poloidal array of modes).

Finally, in order to duplicate an arbitrary equilibrium it is necessary to have some waves that drive current in the medium density region. For this purpose we somewhat arbitrarily chose the (high frequency) fast wave at $f = 0.8$ GHz, launched on the midplane. Figures 2.H-7 and 2.H-8 show the trajectory and off-axis damping of a typical mode, with $N_{||} = 2$. Figure 2.H-9 shows the trajectories and $G(\psi)$ for a uniform power spectrum spanning $2.6 \geq N_{||} \geq 1.1$.

The code calculation attempts to optimize the power spectrum to duplicate the $G(\psi)$ of the initial target equilibrium, using a given set of rays. We selected a set of rays from the candidates discussed above, and the code adjusted the power spectra to yield the equilibrium shown in Fig. 2.H-10. We smoothed the G function slightly to eliminate sharp peaks from individual ray contributions. More smoothing or additional rays (more continuous spectrum) would even out the $j_t(R)$ profile. The safety factor, $q(\psi)$, is an integral and hence appears smooth. The converged equilibrium has $\beta = 0.052$, $I = 10.2$ MA, $q_a = 1.2$, and $q_b = 3.36$. Beta poloidal (β_D) is about unity. Note that the macroscopic parameters (β , I , etc.) are very close to the target values of the initial equilibrium, which were to be duplicated. The required power is:

Figure 2.H-5. Low-frequency ($f = 66$ MHz) fast wave current drive; ray trajectory and index of refraction are displayed. At this launch location the parallel index remains near its edge value, $n_{||} = 3.2$.

Figure 2.H-6. Electron and ion damping for $f = 66$ MHz; tritium damping is complete at $\omega = 2 \Omega_T$.

Figure 2.H-7. High-frequency fast wave (800 MHz) trajectory.

Figure 2.H-8. High frequency wave (800 MHz) at $n_u = 2$ and fast wave polarization n^+ penetrates well into plasma.

T2LF03 *** TEST CASE FOR HI FREQ FAST WAVE *** SEVEN RAYS
 T2LF03 TIME-20:59:13 DATE-06/08/87 MACHINE-C SUFFIX-B

Figure 2.H-9. Ray paths and current density generation, G , for fast waves at 800 MHz and a spectrum $2.6 \geq n_{\parallel} \geq 1.1$.

Figure 2.H-10. Self-consistent MHD equilibrium with combination of LH and fast waves. Total power of 114 MW is needed to generate $I = 10.2$ MA and $b = 0.052$. Only 20% of power is LHCD.

22.6 MW, LHCD, 3.7 GHz, $3.5 \geq N_H \geq 1.9$

63.0 MW, FWCD, 800 MHz, $3.0 \geq N_H \geq 1.2$

28.8 MW, FWCD, 66 MHz, $N_H = 3.2$.

The total power is $P_{CD} = 114$ MW. It is likely that P_{CD} could be reduced; for example, q_a could be raised to >1.2 , and a broader j_t profile might be considered. It is also likely that power could be transferred from the 800-MHz to the 66-MHz fast waves, if desired; different equilibria would result. For this equilibrium we find 20% of the ICRF ray power, or a total of 5.8 MW, is lost to ion heating.

2.H.4 ICRF CURRENT DRIVE WITH BOOTSTRAP

Our first study including the bootstrap current assumed flat density and peaked temperature profiles, as in the previous subsection. Since $\beta_p \approx 1$, and the plasma is in the banana regime, the bootstrap effect is expected to be present. We included this classical contribution to current in the manner outlined in Ref. 14 to obtain the self-consistent MHD equilibrium.

The resulting equilibrium is shown in Fig. 2.H-11, where the bootstrap current is broad and averages about 20% or less of the current density. This converged equilibrium has $\beta = 0.052$, $I = 10.2$ MA, $q_a = 1.1$, $q_b = 3.33$, and $\beta_p \approx 1$, as desired. The RF power required is as follows:

16.5 MW, LHCD, 3.7 GHz, $3.5 \geq N_H \geq 1.9$

50.4 MW, FWCD, 800 MHz, $3.0 \geq N_H \geq 1.2$

27.2 MW, FWCD, 66 MHz, $N_H = 3.2$.

The total power is 94 MW. Thus, less than a 20% power reduction occurs with a flat density profile.

2.H.5 ICRF/BOOTSTRAP WITH PEAKED DENSITY PROFILE

Since the bootstrap current density adds a term $H(\psi)$ to $G(\psi)$, with $\langle B^2 \rangle H(\psi) = \langle j_{\parallel} B \rangle \sim -K_{13}(dp/d\psi) - |C|(dT/d\psi)$ [10], we find that, for a fixed $p(\psi)$, the bootstrap term increases as $|dT/d\psi|$ decreases and $|dn/d\psi|$ increases. This was verified quantitatively with the revised baseline plasma parameters, for which $T_e(\psi) = T_{eo} \tilde{\psi}^{0.3}$, $n_e(\psi) = n_{eo} \tilde{\psi}^{1.1}$, $T_{eo} = 24.7$ keV = T_{io} , $\langle T_e \rangle \approx 19$ keV, $n_{eo} = 2.14 \times 10^{20} \text{ m}^{-3}$, $\langle n_e \rangle \approx 1.02 \times 10^{20} \text{ m}^{-3}$, and $Z_{eff} = 1.5$.

T2LF07 *** TEST CASE FOR HI FRED FAST WAVE *** SEVEN RAYS
 T2LF07 TIME=20:52:54 DATE=06/09/87 MACHINE=C SUFFIX=A

Figure 2.H-11. Equilibrium with $b = 0.052$ and $I = 10.2$ MA including bootstrap contribution to current density. For a flat density profile a total power of 94 MW is needed.

A number of differences are found. With the broad temperature profile we find the 66-MHz fast wave is absorbed over a wider region, so this $G(\psi)$ contribution is broadened (cf. Fig. 2.H-12 with Fig. 2.H-6). This is partly due to the evolution of $N_{\parallel}$, which does not increase along the ray trajectory (cf. Fig. 2.H-13 with Fig. 2.H-5). The $N_{\parallel} = 3.2$ fast wave at 66 MHz is totally absorbed at the $2\omega_{CT}$ layer, but now only 11% of the power is lost to ion heating, a considerable improvement. The converged equilibrium, Fig. 2.H-14, shows roughly half the current density is due to the bootstrap effect. At the edge ($\tilde{\psi} = 0$) almost all the current is provided by the bootstrap, and the required LHCD power is negligible. At the center ($\tilde{\psi} = 1$), where $dp/d\psi \rightarrow 0$, the bootstrap vanishes, so the FWCD is required to provide the "seed current" at the magnetic axis. Interestingly, the 800-MHz power demand is considerably reduced. The spectral requirement is:

$$0.1 \text{ MW, LHCD, } 3.7 \text{ GHz, } 3.5 \geq N_{\parallel} \geq 1.9$$

$$9.3 \text{ MW, FWCD, } 800 \text{ MHz, } 3.0 \geq N_{\parallel} \geq 2.0$$

$$51.6 \text{ MW, FWCD, } 66 \text{ MHz, } N_{\parallel} = 3.2.$$

The converged equilibrium has $\beta = 0.051$, $I = 10.2 \text{ MA}$, $q_a = 1.2$, $q_b = 3.36$, and $\beta_p \simeq 1$. The total RF power is only $P_{CD} = 61.0 \text{ MW}$.

Since so little power is needed at 800 MHz we considered TIBER II equilibria generated without this intermediate frequency. We would expect less ability to duplicate arbitrary equilibria without these rays which generate G peaked at the mid-values of $\tilde{\psi}$. Our initial studies in fact have had difficulty finding high beta with $q_a \gtrsim 1.0$. Yet, by adjusting relative power to LHCD and FWCD a variety of innovative equilibrium current profiles can be created. Figure 2.H-15 shows the result with

$$0.4 \text{ MW, LHCD, } 3.7 \text{ GHz, } N_{\parallel} = 1.9$$

$$58.8 \text{ MW, FWCD, } 66 \text{ MHz, } N_{\parallel} = 3.2$$

The total $P_{CD} = 59.2 \text{ MW}$ in this case creates an equilibrium with $\beta = 0.045$, $I = 9.2 \text{ MA}$, $q_a = 0.74$, and $q_b = 3.66$. As a final illustration, Fig. 2.H-16 shows an equilibrium with

Figure 2.H-12. Broad damping of $f = 66$ MHz fast wave with peaked density ($n_0 = 3.20$) and broad temperature profile (baseline TIBER II).

Figure 2.H-13. Ray path and index of refraction for $f = 66$ MHz fast wave ($n_g = 3.20$) in baseline plasma.

T2LF09 *** TEST CASE FOR HI FREQ FAST WAVE *** SEVEN RAYS
 T2LF09 TIME-14:41:49 DATE-06/13/87 MACHINE-D SUFFIX-A

Figure 2.H-14. Converged equilibrium ($\beta = 0.051$, $I = 10.2$ MA) for baseline plasma (peaked density profile). Total power of 61 MW includes only 100 kW of LHCD; 85% of the power resides in the 66 MHz fast wave which provides a seed current for the bootstrap contribution.

T2LF11 *** TEST CASE FOR H1 FREQ FAST WAVE *** SEVEN RAYS
 T2LF11 TIME-21:14:58 DATE-06/15/87 MACHINE-C SUFFIX-A

Figure 2.H-15. Converged equilibrium without high-frequency fast wave, but including bootstrap current (TIBER II baseline): total power of 59 MW consists of only 0.4 MW LHCD, and the remainder is 66 MHz FWCD. Result has $\beta = 0.045$, $I = 9.2$ MA, and $q_a < 1.0$ at the magnetic axis.

T2LF13 *** TEST CASE FOR HI FREQ FAST WAVE *** SEVEN RAYS
 T2LF13 TIME-00:55:02 DATE-06/16/87 MACHINE-C SUFFIX-A

Figure 2.H-16. Converged equilibrium with higher LHCD power (1.6 MW) and less FWCD power (55 MW) at 66 MHz, including bootstrap current. Total power of 57 MW generates $\beta = 0.046$ and $I = 11.0$ MA. Note $q(\psi)$ is monotonic, although there is a sharp surface current density.

1.6 MW, LHCD, 3.7 GHz, $N_H = 1.9$

55.3 MW, FWCD, 66 MHz, $N_H = 3.2$

This case has higher current, $I = 11.0$, and $\beta = 0.046$, even though less power, $P_{CD} = 56.9$ MW, is used. It has $q_a = 0.78$, $q_b = 2.56$, and $q(\psi)$ is monotonic in ψ .

2.H.6 FUTURE WORK

The $2\omega_{CT}$ resonance remains a problem power sink. If $A > 3.6$ were possible, this would move the resonance away from the hot plasma region.

Nonthermal alpha damping needs study for the FWCD approach. Preliminary studies [12] and our own calculations indicate that alpha damping can be a small loss mechanism compared to electron damping, but detailed calculations should be performed.

In fact, work is now being carried out to develop a model for the alpha damping of the fast wave and an approximate scheme to take into account the residual absorption of the impurity ions via higher harmonic cyclotron damping. In particular, the concentration of the alpha density over the plasma volume and the local velocity distribution function need to be calculated, requiring a knowledge of the alpha transport. For the case of a weakly damped fast wave, a model for the wall absorption also needs to be implemented. All these damping processes will be incorporated into the ray tracing/MHD equilibrium code [3] to rigorously assess the viability of the proposed FW/LH current drive scenario for TIBER II.

Antenna design details need to be considered. Typical issues revolve around the choice of loops, ridged, reentrant, or slotted waveguides.

Finally, the most promising study area may involve questions of how to further optimize FWCD/LHCD wave spectra in order to tailor the current profile and provide better plasma stability and/or confinement.

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APPENDIX 2.I

NEUTRAL BEAM CURRENT DRIVE

This is an alternative to the baseline study of neutral beam injection (NBI) current drive for steady-state operation of TIBER II [1].

Present current-drive theories do not account for all the physical phenomena that affect the beam driven current. In particular, effects of radial momentum transfer, which analysis of unbalanced NBI rotation experiments [2] has shown to be very important in explaining the behavior of the observed toroidal velocities, are not included.

A theory is under development to take into account in a self-consistent way the effects of radial momentum transfer in the calculation of the neutral-beam-driven current. The predictions of the theory will be compared with data from TFTR and other experiments.

Meanwhile, to establish a basis for comparison with the existing current drive theories, current drive calculations have been done for TIBER II, using standard theory. The net beam driven current density is given by,

$$J_{\text{net}} = \left[1 - \frac{Z_b}{Z_{\text{eff}}} (1 - G) \right] J_f, \quad (2.I-1)$$

where the trapped electron correction, G , depends on the inverse aspect ratio, $\epsilon = r/R_0$, and on Z_{eff} . The finite aspect ratio calculation of G is approximated by [3,4]

$$G \approx \left(1.55 + \frac{0.85}{Z_{\text{eff}}} \right) \sqrt{\epsilon} - \left(0.20 + \frac{1.45}{Z_{\text{eff}}} \right) \epsilon \quad (2.I-2)$$

In Eq. 2.I-1, J_f is the fast ion current density given by,

$$J_f = n_b \tau_s Z_b e \epsilon_0 v_b^0 K_e, \quad (2.I-3)$$

where n_b is the fast ion source rate, τ_s is the fast ion slowing down time, ϵ_0 is the cosine of the initial pitch angle, v_b^0 is the initial velocity of the

beam particles, and K_e is the fraction of the initial fast ion toroidal momentum deposited in the plasma electrons.

The flux-surface-average fast ion source rate is

$$n_b = \frac{I}{eV_p} H(\rho) \quad (2.I-4)$$

where V_p is the plasma volume, I/e is the atomic current of the beam, and $H(\rho)$ is the neutral beam deposition profile as a function of the midplane minor radius ρ . $H(\rho)$ is computed with the code PROCTR [5] and the calculation is based on an analytical model described by Rome et al. [6]. The beam current is taken to be circular in cross section with a Gaussian profile. Plasma flux contours are shifted ellipses with constant elongation.

The momentum transfer function K_e , is also computed with PROCTR which uses a moments solution of the fast ion Fokker-Planck equation [7].

In Fig. 2.I-1, the fast ion current density j_f , and the net current density j_n , are plotted versus the minor radius. Two identical co-injected 500-keV negative ion neutral beams were assumed, with tangency radii of 2.6 m and 3.0 m respectively, in accordance with the baseline design. The two beams carry a total of 45 MW of power. A total beam driven current of 4.8 MA is obtained, resulting in an efficiency of 0.107 A/W and a figure-of-merit gamma equal to $0.34 \cdot 10^{20}$ A/W m².

To check the sensitivity of the current drive efficiency to the beam energy, the beam driven current was calculated for different beam energies, keeping the rest of the parameters constant and equal to their reference values. The results are shown in Fig. 2.I-2, where the current drive efficiency is plotted versus the beam energy. It can be seen that the efficiency has a maximum around 1 MeV and then it drops while the beam shinenethrough increases to approximately 17% at 2 MeV.

Figure 2.I-1. Neutral beam driven current in TIBER II for reference parameters.

Figure 2.I-2. Current drive efficiency in TIBER II for different beam energies.

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