OR-486

ORNL/TM-5573

ORNL/TM-5573

Oak Ridge Tokamak Experimental Power Reactor Study—1976

Part 2

Plasma Engineering in a Deuterium-Tritium Fueled Tokamak

D. G. McAlees

F. B. Marcus

S. E. Attenberger

R. N. Cherdack

G. D. Kerbel

J. R. McNally, Jr.

Y-K. M. Peng

G. Schilling

D. J. Strickler

N. A. Ukan

MASTER

OAK RIDGE NATIONAL LABORATORY

OPERATED BY UNION CARBIDE CORPORATION FOR THE ENERGY RESEARCH AND DEVELOPMENT ADMINISTRATION

DISTRIBUTION OF THIS DOCUMENT IS UNLIMITED

DISCLAIMER

This report was prepared as an account of work sponsored by an agency of the United States Government. Neither the United States Government nor any agency Thereof, nor any of their employees, makes any warranty, express or implied, or assumes any legal liability or responsibility for the accuracy, completeness, or usefulness of any information, apparatus, product, or process disclosed, or represents that its use would not infringe privately owned rights. Reference herein to any specific commercial product, process, or service by trade name, trademark, manufacturer, or otherwise does not necessarily constitute or imply its endorsement, recommendation, or favoring by the United States Government or any agency thereof. The views and opinions of authors expressed herein do not necessarily state or reflect those of the United States Government or any agency thereof.

DISCLAIMER

Portions of this document may be illegible in electronic image products. Images are produced from the best available original document.

Printed in the United States of America. Available from
National Technical Information Service
U.S. Department of Commerce
5285 Port Royal Road, Springfield, Virginia 22161
Price: Printed Copy \$6.00 ; Microfiche \$2.25

This report was prepared as an account of work sponsored by the United States Government. Neither the United States nor the Energy Research and Development Administration/United States Nuclear Regulatory Commission, nor any of their employees, nor any of their contractors, subcontractors, or their employees, makes any warranty, express or implied, or assumes any legal liability or responsibility for the accuracy, completeness or usefulness of any information, apparatus, product or process disclosed, or represents that its use would not infringe privately owned rights.

Contract No. W-7405-eng-26

Thermonuclear Division

PLASMA ENGINEERING IN A DEUTERIUM-TRITIUM FUELED TOKAMAK*

D. G. McAlees (ed.)
Exxon Nuclear Co., Inc.
F. B. Marcus (ed.)

S. E. Attenberger	Y-K. M. Peng
R. N. Cherdack	G. Schilling
Burns & Roe, Inc.	D. J. Strickler
G. D. Kerbel	N. A. Uckan
J. R. McNally, Jr.	

*Part 2 of Oak Ridge Tokamak Experimental Power Reactor Study 1976 (ORNL/TM-5572-5577).

Date Published: October 1976

OAK RIDGE NATIONAL LABORATORY
Oak Ridge, Tennessee 37830
operated by
UNION CARBIDE CORPORATION
for the
ENERGY RESEARCH AND DEVELOPMENT ADMINISTRATION

MASTER

THIS PAGE
WAS INTENTIONALLY
LEFT BLANK

ABSTRACT

Plasma Engineering includes both basic investigations and the application of research results to fusion system design. Plasma Engineering provides fusion physics input to various CTR programs so that the effects of system operating limitations and requirements can be considered in the design process.

In this report a broad scope of Plasma Engineering information is presented. For example, tokamak dimensions and scaling, magnetic field requirements, impurity control and wall material criteria, poloidal field system design, effects of changing the number and shape of the toroidal field coils, plasma ionization and breakdown, fueling, plasma heating by neutral beam injection, vacuum requirements and components, low-neutron-yield testing and prospects for clean fuels are discussed.

The analyses and techniques developed are specifically applied to the Oak Ridge Experimental Power Reactor (EPR).

THIS PAGE
WAS INTENTIONALLY
LEFT BLANK

CONTENTS

ABSTRACT	iii
LIST OF TABLES	ix
LIST OF FIGURES	xi
1. INTRODUCTION	1
2. PLASMA SIMULATION MODELS	5
2.1 Physics Models in the Codes	5
2.1.1 Beam-Plasma Fusion Power	7
2.1.2 Neutral Beam Injection Heating	8
2.1.3 Containment Times	8
2.1.4 Alpha Energy Deposition	9
2.1.5 Impurity Radiation Effects	9
2.2 Point Modeling — Static Code	10
2.3 Reactor Modeling — Time-Dependent SYSTEMS Code	14
3. IMPURITY MODELING AND EFFECTS ON SYSTEM DYNAMICS	25
4. POLOIDAL FIELD SYSTEM DESIGN	35
4.1 Equilibrium and Stability Requirements	35
4.1.1 Equilibrium Requirements	35
4.1.2 Stability Requirements	36
4.2 Shielding Concept and Computing Method	36
4.2.1 Channeling the OH (Ohmic Heating) Flux	41
4.2.2 Shielding Against the Plasma Flux	41
4.2.3 Tailoring the Vertical Field	41
4.2.4 Method of Poloidal Field Coil Design	43
4.2.5 Optimal Current Distribution	43
4.2.6 Approximation with Discrete Coils	43
4.3 Coil Locations and Circuits	44
4.3.1 OH Coils	44
4.3.2 S-VF Coils	44
4.3.3 T-VF Coils	46
4.3.4 Decoupling Coils	47
4.4 Poloidal Field Shielding	47
4.4.1 Time Behavior of Coil Currents	47
4.4.2 Poloidal Fields	49
4.5 Volt Seconds and Power Supply	57
4.6 Other Considerations	57
4.6.1 Divertor Configuration	60
4.6.2 Noncircular Cross Section	60
5. TOROIDAL FIELD RIPPLE	64
5.1 Introduction	64
5.2 Transport	66

5.3	Fast Ion Losses	68
5.4	Relaxation of Toroidal Flow by Ripple Viscosity	70
6.	START-UP: BREAKDOWN, PRE-IONIZATION EFFECTS, RUNAWAY PRODUCTION, COUPLED RATE EQUATIONS	74
6.1	Introduction	74
6.2	Gas Discharge Parameters	74
6.3	Plasma Buildup Rate	75
6.4	Choice of Pressure in EPR	75
6.5	Runaway Electrons	76
6.6	Pre-ionization by Microwave ECRH	76
6.7	Heating Volume	77
6.8	Microwave Pre-ionization Rate Equations	77
6.9	Coupled Rate Equations	79
6.10	Neutral Temperature Equations	80
6.11	Analytic Model of Minimum E/p for Complete Ionization	80
6.12	Numerical Results: Application to EPR	81
7.	FUELING	92
7.1	Plasma	92
7.2	Neutral Gas Blanket	92
7.3	Pellets	93
7.3.1	A Steady-State Formulation of the Pellet-Plasma Interaction	93
7.3.2	Results of the Pellet-Plasma Interaction Model	101
7.3.3	Pellet Acceleration Techniques	103
7.3.4	Mechanical Properties and Acceleration Limit of Solid Hydrogen	106
8.	NEUTRAL BEAM INJECTION SYSTEM	109
8.1	Definition and Requirements	109
8.2	Consideration of Possibilities	109
8.3	Recommendations for Reference Design	112
8.4	Principal Development Requirements	112
9.	VACUUM SYSTEMS	115
9.1	Vacuum Systems for EPR	115
9.2	Main Vacuum System	115
9.2.1	System Description	116
9.2.2	Performance Evaluation	117
9.3	Beam Vacuum System	117
9.3.1	System Description	118
9.3.2	Performance Analysis	118
9.4	Extrapolation of EPR Design to Larger Reactors	119

10. ADVANCED FUEL PROSPECTS	121
10.1 Use of DD and D ³ He in Testing	121
10.2 Thermal Runaway Without the Dissipative Trapped-Ion Mode Losses	121
10.3 DD Advanced Fusion Fuel Prospects	121
10.4 Noncircular Tokamaks	124

THIS PAGE
WAS INTENTIONALLY
LEFT BLANK

LIST OF TABLES

2.1	EPR machine parameters	18
2.2	Steady-state plasma parameters and power balance	21
3.1	Peak sputtering yields	28
4.1	Locations and currents of the OH coils	46
4.2	Locations and currents of the S-VF coils	46
4.3	Locations and peak currents of the T-VF coils	47
4.4	Locations and peak currents of the decoupling coils	47
4.5	Per (turn) ² inductances (μH) of the poloidal field systems shown in Fig. 4.6	58
4.6	Comparisons between magnetic shielding and conventional poloidal field systems	59
5.1	Typical parameters	65
6.1	Tokamak parameters	74
6.2	Ionization time	75
6.3	Maximum electron energies	77
6.4	Rate coefficients $S(\text{m}^3 \text{ sec}^{-1})$ and processes	79
6.5	EPR start-up parameters	85
8.1	Representative design values for four possible types of EPR neutral beam injectors	111
8.2	Advantages and disadvantages of the four possible types of EPR neutral beam injectors	112
9.1	EPR, Demo, CPR vacuum systems	119
9.2	Pump-down time and neutral density	119

THIS PAGE
WAS INTENTIONALLY
LEFT BLANK

LIST OF FIGURES

2.1	Beam power per unit volume and the ratio of Maxwellian alpha power to total alpha power as a function of N_T for several $\beta_{\theta e}$ at $T = 8$ keV	11
2.2	Plasma radius required to attain temperature of 8 keV for different trapped particle loss coefficients	12
2.3	Beta poloidal and plasma radius as a function of N_T for different toroidal magnetic field strengths	13
2.4	Minimum minor radius required for ignition at given B_T and T ($1/10 \times T.P.$ losses) . . .	15
2.5	The density-confinement time product as a function of temperature for several plasma radii, and the ignition condition	16
2.6	Beam power per unit volume and the ratio of Maxwellian fusion power to total fusion power as a function of N_T for several Z_{EFF} at $T = 8$ keV	17
2.7	Beam power per unit volume and the ratio of Maxwellian fusion power to total fusion power as a function of N_T for several temperatures at $Z_{EFF} = 2$	19
2.8	EPR steady-state operating points, $P_{thermal}$ vs N_e	20
2.9	EPR steady-state operating points, $\beta_{\theta e}$ vs N_e	22
2.10	EPR steady-state operating points, T_i vs N_e	23
3.1	Line and recombination radiation losses (integrated over θ and ϕ) vs r	26
3.2	Line and recombination radiation losses (integrated over θ and ϕ) vs r	27
3.3	Incorporation of sputtering and impurity effects in the plasma simulation model	29
3.4	Time dependence of temperature for various impurity control conditions	31
3.5	Sputtering coefficient as a function of incident deuteron energy and time dependence of temperature for two sputtering models	32
3.6	Time-dependent fusion thermal power and density for two assumptions on alpha-particle containment	34
4.1	Free-boundary equilibria with fixed I_p and $n = 0.5$	37
4.2	Dependence of the plasma shape on the decay index	38
4.3	The critical β_{q_a} and q_0 as a function of β_p for stability of ideal local modes . . .	39
4.4	Idealized model of poloidal field shielding (with up-down symmetry)	40
4.5	Shielding (eddy) current $J_S (=J_{OH} + J_{VF})$ on a toroidal, circular, conducting shell . .	42
4.6	Poloidal field system (a) with magnetic shielding and (b) without	45
4.7	The assumed behavior of the amp-turns in the OH, S-VF, D, T-VF, and VF coils	48
4.8	The current and pressure profiles used in the equilibria at $\beta_j = .4$ and 1.3	50
4.9	"Snapshots" of field due to currents and corresponding to those of Fig. 4.3	51
4.10	Average poloidal field over the TF coils	56
4.11	Divertor, magnetic limiter configuration	61
4.12	(a) Typical free-boundary, D-shaped plasma flux plot. (b) The vacuum equilibrium field	62
5.1	$G(\alpha, \beta)$ versus α	67
5.2	Heat conduction parameters versus plasma radius for the ORNL EPR design parameters . .	69
5.3	Toroidal flow velocity versus plasma radius for the ORNL EPR design parameters	72
6.1	EPR start-up: $T_{10,20} = 0.1$ eV, $n_{20} = 0.6 \times 10^{19} \text{ m}^{-3}$	82
6.2	EPR start-up: $T_{10,20} = 0.1$ eV, $n_{20} = 0.1 \times 10^{19} \text{ m}^{-3}$	83
6.3	EPR start-up: $T_{10,20} = 0.1$ eV, $n_{20} = 0.1 \times 10^{19} \text{ m}^{-3}$, $v_1/v_0 = 4$	84
6.4	EPR start-up: $T_{10,20}$ variable, $n_{20} = 0.1 \times 10^{19} \text{ m}^{-3}$	86
6.5	EPR start-up: $T_{10,20}$ variable, $n_{20} = 0.3 \times 10^{19} \text{ m}^{-3}$	87
6.6	EPR start-up: $T_{10,20}$ variable, $n_{20} = 0.6 \times 10^{19} \text{ m}^{-3}$	88
6.7	EPR start-up: $T_{10,20}$ variable, $n_{20} = 0.12 \times 10^{19} \text{ m}^{-3}$	89

7.1 Pellet geometry	96
7.2 The steady ablation of a pellet encountering gradual changes in plasma conditions . .	102
7.3 The dynamic ablation of a pellet whose shield cannot keep up with the changes in plasma conditions	104
8.1 Model of neutral beam injector systems	110
8.2 EPR reference design neutral beam injector system	113
10.1 Schema of burning characteristics of DT systems	122
10.2 Radiation-dominated DT reactor (pulsed burn, closed system)	123
10.3 Fractional burn-up in percent in steady-state catalyzed DD reactors	125

1. INTRODUCTION

Plasma Engineering (PE) involves the use of the results from applied physics analyses in the design of fusion systems. In some cases the physics research must be carried out by PE, while in others analysis techniques are available and only need to be implemented. The design process is made difficult by the fact that some of the most critical aspects of the system and its behavior are not well understood at present. Thus the procedure involves making plausible estimates and retaining a flexibility in the design work so that it can always accommodate change. This approach is motivated by the need for design criteria which permit engineering and technological development.

In this report, several of the main topics and the logic used to determine fusion system parameters are outlined. Most of the information here was developed in support of the Experimental Power Reactor (EPR) and Demonstration Power Reactor (Demo) studies at ORNL. These specific systems are described. However, it is important to note that the models and analysis techniques are generally applicable to other fusion systems and will be of immediate use in their design.

There are nine main sections in this report. They are summarized in paragraphs A-I below. The sections address the problems that are among the most critical to reactor design. The document provides information on problems related to tokamak dimensions, magnetic field strength, burn time, requirements on impurity control and wall materials, position and current waveforms of the poloidal field (PF) system, external pulsed magnetic fields, number and shape of toroidal field (TF) coils, plasma ionization and breakdown requirements, refueling of a burning plasma, power and energy needs for plasma heating, vacuum pumping, low-neutron-yield plasma testing, and prospects for clean fuels. Specific answers to these problems are required, and an attempt has been made to provide results based on the best available knowledge.

Section 2 describes the physics and application of plasma system simulation models which facilitate the establishment of reference parameters for a particular device, and the examination of its behavior under varying input conditions. Section 3 discusses impurity control separately since the influence of impurities on the detailed energy balance is critically important. The next two sections, Sects. 4 and 5, discuss aspects of magnetic design. The Poloidal Field System and its implicit relationship to plasma equilibrium and stability are treated first, and then the importance of ripples in the toroidal magnetic field structure and their relationship to the design of the Toroidal Field Coil System are discussed.

Section 6 describes the start-up physics associated with the breakdown and current-rise phases of a plasma. Section 7 discusses the fueling of an operating plasma, wherein the analysis is restricted to pellet ablation physics and acceleration techniques. Neutral beam injection is assumed to be the source of supplemental heating needed to achieve ignition in the EPR and Demo. The characteristics and research requirements of high-energy high-power beam systems are described in Sect. 8. Vacuum system features, criteria, and components are given in Sect. 9. Section 10 describes the use of advanced fuels in the testing phases of a reactor system and outlines the advantages which would be enjoyed if advanced fuels can ultimately be contained and burned.

A. Plasma Simulation Models

The plasma simulation models include: Maxwellian and beam-plasma fusion; fast-ion heating; refueling; plasma conduction and convection; bremsstrahlung, synchrotron, and line radiation; neutron, ion, neutral, chemical, and self-sputtering; and charge exchange. A point-static code is

used to scan parameter space and examine scaling sensitivities. The basic EPR parameters are: $a = 2.25$ m, $R_0 = 6.75$ m, $B_T = 4.8$ T, and $I = 7.2$ MA. (For definition of parameters, see Tables 2.1 and 2.2 in Sect. 2.) A time-dependent code SYSTEMS, which includes the profile effects of impurities and charge-exchange neutrals, is used to find the equilibrium operating point. The thermal power output is highly sensitive to density. Even if 50 MW of neutral beam injection power is applied, low-density plasmas will not ignite. This result is due to the scaling assumptions made. High-density ignited plasmas produce thermal powers in excess of 250 MW. In the reference design, $N_e = 7.4 \times 10^{19} \text{ m}^{-3}$, $\beta_0 = 1.9$, and $P_{\text{thermal}} = 410$ MW, during the steady-state phase.

B. Impurity Modeling and Effects on System Dynamics

The SYSTEMS code is used to model the time-dependent behavior of impurity and alpha particle buildup. The impurities result from sputtering by a variety of processes. The dominant sputtering term for an edge temperature of 100 eV is due to charge exchange. The charge-exchange flux is computed from radially dependent neutral and plasma densities. Line radiation from the sputtered impurities is computed from radially varying impurity, density, and temperature profiles. Plasma behavior is computed for a variety of assumptions about impurity diffusion, divertor efficiency, wall-reflux efficiency, neutral density, and wall material. If the impurity diffusion time is similar to that of the plasma ions, a steady state burn requires 90% divertor efficiency and 90% shielding of incoming impurities. More stringent conditions limit the burn to less than 50 sec. The results are sensitive to sputtering coefficients. In the absence of wall impurities, alpha particle buildup could limit the burntime to 120 sec.

C. Poloidal Field System Design

The poloidal field system with magnetic shielding for the ORNL EPR is described. It is consistent with a long-pulse high- β_0 tokamak with superconducting TF coils. The system is nearly optimal in the sense that it minimizes the number of volt-seconds required to set up the plasma current to permit increased pulse length and duty cycle. It minimizes the pulsed field (a reduction by a factor of 6) at the superconducting TF coils to permit attainment of the highest toroidal field with the least risk of superconducting quenching. It has equilibrium field coils which are closely coupled to the plasma and respond rapidly to large increases in plasma β . An intrinsic advantage of such a system is a substantial reduction in supplied power (by 41%) and stored energy (by 34%), compared to these designs without magnetic shielding. A detailed description of the magnetic field shielding system and a comparison of it with a conventional poloidal field system without shielding are presented. In view of the large uncertainties in the development of large superconducting high-field coils (TF) and pulsed high-field coils (OH-primary), the magnetic field shielding system allows use of simpler superconductor coil designs for future tokamak reactors.

D. Toroidal Field Ripple

In the design of tokamak devices, the choice of the number of TF coils depends on the toroidal field ripple that can be tolerated by plasma physics considerations and the space needed for access to the device. For a given set of parameters, typical of EPR-size plasmas, the calculated enhanced

heat loss due to ripples will be less than the total heat conduction loss due to neoclassical and trapped particle modes at all radii, provided that the magnetic field ripple at the plasma boundary is less than about 2.5%. Fast-ion losses due to ripple trapping might have no significant effect on power balance, but nonetheless need to be recognized. The magnetic field ripple has the beneficial effect of producing a viscous drag on toroidal rotations, except near the center of the plasma. The momentum input can be made small near the plasma center by balancing co- and counter-injected beams. For this purpose a higher magnetic field ripple near the plasma center might be useful.

In the ORNL EPR Reference Design with 20 TF coils, the maximum ripple of about 2.2% at the outer edge of the plasma column has been found to be tolerable according to the calculations described above.

E. Start-up: Breakdown, Pre-Ionization Effects, Runaway Electron Production, Coupled Rate Equations

The plasma start-up problem is considered on the basis of several effects. At very low pre-ionization density, the breakdown is described by the gas discharge parameters, E/p , α/p , V_D , and T_e . The effect of preionization density upon volt-second consumption is calculated. At low electric fields and high pressures, up to 10% of the EPR volt-seconds may be consumed. During this period runaway electrons may be produced, and E/p thresholds for this process are calculated. One method of creating a pre-ionized plasma is by resonant electron cyclotron heating, and the required input powers are determined. The plasma start-up during the period of partial to full ionization is described by a set of coupled rate-equations for the densities and temperatures of H_1^0 , H_2^0 , H_1^+ , H_2^+ , e , or their isotopes, and current. The numerical simulations show that wall effects can be very important during start-up. At a voltage of 50 volts in EPR, rapid ionization occurs at initial densities of H_2^0 below $0.6 \times 10^{19} \text{ m}^{-3}$, leading to an electron density far below the desired level. Small-radius start-up schemes are considered.

F. Fueling

In the most general terms, fueling a fusion reactor means maintaining the density of fuel ions in the reactor core at the operating state, while keeping a clean core, $Z_{\text{eff}} \approx 1$, and simultaneously removing the ash created by fusion.

The question of fueling is addressed directly and that of keeping the core clean, indirectly. There are two methods of accomplishing the first in the EPR: (a) a high density can be maintained at the plasma edge and dependence placed on diffusive mixing to maintain the central density; (b) neutral atoms can be injected directly into the core. Approach (a) appears to be the simpler of the two methods, but it seems that impurity buildup in the core may invalidate a fueling system based on this approach. Neutral atom injection (b) is assumed to be a feasible means of fueling, and the physics of the pellet-plasma interaction is discussed.

The existence of an ablation layer of dense gas which forms a shield around a frozen fuel pellet is reported, and its properties with respect to its effect on pellet lifetimes and penetration depths are examined. Research presently directed toward the development of pellet acceleration devices is described. In some areas, convenient assumptions have been made and are explained. Some further work which is needed has been indicated and discussed.

G. Neutral Beam Injection System

The function of the neutral beam injection system is to contribute to start-up heating of the plasma and, if ignition is not reached, to provide supplementary heating for achieving useful power

levels in a beam-driven EPR. The magnitude of the injected power, 50 MW, makes it of the highest importance to efficiently convert primary electric power into 200-keV neutral beam power. Four possible injection systems have been studied and their associated efficiencies, η , calculated: (1) positive ions, $\eta = 0.15$; (2) positive ions with direct recovery of the unneutralized ion energy, $\eta = 0.36$; (3) negative ions, $\eta = 0.50$; and (4) negative ions with direct recovery of the unneutralized ion energy, $\eta = 0.69$. The negative ion system without recovery (3) appears to offer moderate electrical efficiency as well as requiring fewer injectors and beam lines than the positive ion systems. It is adopted as the reference injection system.

H. Vacuum Systems

The vacuum systems pump the plasma exhaust, impurities, beam lines, and outgassing products. This is accomplished by a combination of cryopanel, cryosorption pumps, Roots Blowers, and mechanical pumps. In the absence of a divertor, the most critical uncertainties are the pumping of plasma exhaust during operation, and boundary neutral density. The reference system should achieve a base pressure of 2×10^{-10} torr in the presence of outgassing, and should reduce the deuterium/tritium (DT) pressure in the chamber by 10^3 during the 5-sec interval between pulses.

In the beam line enormous pumping speeds are required to achieve low densities downstream from the stripping cell, necessitating a large array of cryopanel, with attendant problems of neutron and gamma-ray heating.

Comparisons are made between EPR, Demo, and CPR (Commercial Power Reactor) vacuum systems.

I. Advanced Fuel Prospects

DT is the "classical" fusion fuel for ignition experiments and first-generation fusion reactors. The advanced fuels D^3He and DD offer important prospects for exploratory plasma studies prior to DT ignition with its potentially severe thermal runaway. These studies will permit essential evaluations of plasma behavior and plasma space parameters. Follow-on studies of the properties of an ignited DT system ($T_0 = 6$ keV, $n\tau > 8 \times 10^{14} \text{ cm}^{-3} \text{ sec}$) may reveal less severe losses than the dissipative trapped-ion mode losses ($n\tau \propto T_e^{-11/2}$), thermal runaway to high ion temperatures ($T_i > 50$ keV), and large $n\tau$ values (2×10^{15}). Attainment of such energetic operating conditions would demonstrate the feasibility of the advanced fusion fuel reactors utilizing DD or D^6Li in a catalyzed burner mode.

2. PLASMA SIMULATION MODELS

At the present state of knowledge, the design of fusion reactors is based on scaling which has not been fully verified experimentally. Various choices of scaling laws can be made, leading to different results. The following studies demonstrate the sensitivity of the design to the scaling laws assumed. Models suitable for simulating the energy and particle balances of a tokamak plasma, for surveying the range of parameters possible during operation, and the sensitivities of the parameters to system changes, are summarized in this section. Physical models for transport losses, beam-plasma interaction, noncircular plasma behavior, sputtering, impurity effects, and reactor control are included. The plasma simulation models summarized here are described in detail in a separate study.¹ The physics models in the codes, the point model, and the SYSTEMS code are briefly outlined here, together with typical applications for the ORNL EPR² and similar devices. The limitations on the operation of such devices due to constraints such as beta limits, sputtering, and impurity buildup are of special interest.

2.1 Physics Models in the Codes

The particle and energy balance equations used in the code include deuterium, tritium, alpha, carbon (to represent low-Z impurities), and iron density and electron and ion temperature. The equations are:

$$\frac{d}{dt}(N_D) = S_{D_c} - N_D N_T \langle \sigma v \rangle_{DT} - \frac{N_D}{\tau_D} + S_D - \frac{S_D E_b}{17600} \frac{N_T}{N_e} F(E_b, T_e) \quad (1)$$

$$\frac{d}{dt}(N_T) = S_{T_c} - N_D N_T \langle \sigma v \rangle_{DT} - \frac{N_T}{\tau_T} - \frac{S_D E_b}{17600} \frac{N_T}{N_e} F(E_b, T_e) \quad (2)$$

$$\frac{d}{dt}(N_\alpha) = N_D N_T \langle \sigma v \rangle_{DT} - \frac{N_\alpha}{\tau_\alpha} + \frac{S_D E_b}{17600} \frac{N_T}{N_e} F(E_b, T_e) \quad (3)$$

$$\begin{aligned} \frac{d}{dt}(N_c) = & -\frac{N_c}{\tau_c} + f(T) A_c d_c \left[\frac{\gamma_{cD} N_D}{\tau_D} + \frac{\gamma_{cT} N_T}{\tau_T} \right. \\ & \left. + \frac{\gamma_{c\alpha} N_\alpha}{\tau_\alpha} + \frac{\gamma_{cC} N_C}{\tau_C} + \frac{\gamma_{cFe} N_{Fe}}{\tau_{Fe}} \right] \\ & + d_{cn} A_c \gamma_{cn} \left(N_D N_T \langle \sigma v \rangle_{DT} + \frac{S_D E_b}{17600} \frac{N_T}{N_e} F(E_b, T_e) \right) \\ & + d_{cn} A_c \frac{(2.19 \times 10^5)}{a^2} \frac{N_e(0)}{N_e} \bar{\sigma}_{cx} \\ & \times \int_0^a f(T(r)) \frac{N_e(r)}{N_e(0)} N_o(r) \sqrt{T(r)} r dr \\ & \times \left[\frac{\gamma_{cD} N_D}{\sqrt{2}} + \frac{\gamma_{cT} N_T}{\sqrt{3}} \right] + d_c \gamma_{ch} \left[\frac{N_D}{\tau_D} + \frac{N_T}{\tau_T} \right] \end{aligned} \quad (4)$$

$$\begin{aligned}
\frac{d}{dt}(N_{Fe}) = & -\frac{N_{Fe}}{\tau_{Fe}} + f(T) A_{Fe} d_{Fe} \left[\frac{\gamma_{FeD} N_D}{\tau_D} + \frac{\gamma_{FeT} N_T}{\tau_T} + \frac{\gamma_{Fe\alpha} N_\alpha}{\tau_\alpha} + \frac{\gamma_{FeC} N_C}{\tau_C} \right. \\
& \left. + \frac{\gamma_{FeFe} N_{Fe}}{\tau_{Fe}} \right] + d_{Fen} A_{Fe} \gamma_n \left(N_D N_T \langle \sigma v \rangle_{DT} + \frac{S_D E_b}{17600} \frac{N_T}{N_e} F(E_n, T_e) \right) \\
& + d_{Fen} A_{Fe} \frac{(2.19 \times 10^5)}{a^2} \frac{N_e(0)}{N_e} \frac{1}{\sigma_{cx}} \left[\frac{\gamma_D N_D}{\sqrt{2}} + \frac{\gamma_T N_T}{\sqrt{3}} \right] \\
& \times \int_0^a f(T(r)) \frac{N_e(r)}{N_e(0)} N_0(r) \sqrt{T(r)} r dr
\end{aligned} \tag{5}$$

$$\begin{aligned}
\frac{d}{dt} \left(\frac{3}{2} N_e T_e \right) = & N_D N_T \langle \sigma v \rangle_{DT} U_{\alpha e} - \frac{\left(\frac{3}{2} N_e T_e \right)}{\tau_{e,E}} - 0.30 \times 10^{-20} Z_{eff} N_e \sqrt{T_e} \\
& - \frac{1.5 \times 10^5}{\sqrt{a}} \sqrt{N_e} B_T^{5/2} (1 - \beta)^{5/4} T_e^{11/4} (1 - R_e) \left(1 + \frac{T_e}{204} \right) \\
& + 1.5 \times 10^{-19} \frac{N_e}{T_e^{3/2}} \ln \Lambda \sum_j \left(\frac{Z_j N_j}{A_j} \right) (T_j - T_e) + n J^2 \\
& + S_D G_{eD} + S_T G_{eT} + \frac{S_D E_b}{17600} \frac{N_T}{N_e} F(E_b, T_e) U_{\alpha e} - P_{LINE}
\end{aligned} \tag{6}$$

$$\begin{aligned}
\frac{d}{dt} \left(\frac{3}{2} N_I T_I \right) = & N_D N_T \langle \sigma v \rangle_{DT} U_{\alpha I} + S_D G_{ID} + S_T G_{IT} \\
& - \frac{\left(\frac{3}{2} N_I T_I \right)}{\tau_{I,E}} - 1.5 \times 10^{-19} \frac{N_e}{T_e^{3/2}} \ln \Lambda \sum_j \left(\frac{Z_j^2 N_j}{A_j} \right) (T_j - T_e) \\
& + S_{Dc} E_{Dc} + S_{Te} E_{Tc} + \frac{S_D E_b}{17600} \frac{N_T}{N_e} F(E_b, T_e) U_{\alpha I}
\end{aligned} \tag{7}$$

The parameters used in the equations are defined as follows:

S_D, S_T	Deuterium and tritium source rate from fast ions, e.g., due to neutral injection.
$N_D, N_T, N_\alpha, N_C, N_{Fe}$	Average deuterium, tritium, alpha, carbon, and iron densities.
N_I, N_e	Average ion and electron densities.
$\tau_D, \tau_T, \tau_\alpha, \tau_C, \tau_{Fe}$	Particle confinement times of deuterium, tritium, alpha, carbon, and iron ions.
S_{Dc}, S_{Tc}	Deuterium and tritium source rate from fueling, e.g., due to injected recycling, injected pellets, or gas blanket.
E_b	Neutral beam injection energy.
$F(E_b, T_e)$	Ratio of beam-plasma fusion power to injected beam power, for D into pure T.
$\langle \sigma v \rangle_{DT}$	D-T reaction rate parameter.

$f(T)$	Sputtering yield as a function of temperature.
A_c, A_{Fe}	Fractional abundance of carbon and iron in wall.
$d_c, d_{cn}, d_{Fe}, d_{Fen}$	Impurity control inefficiency for sputtering of carbon and iron by charged and neutral particles.
$\gamma_{cD}, \gamma_{cT}, \gamma_{c\alpha}, \gamma_{cc}, \gamma_{cFe}, \gamma_{FeD}, \gamma_{FeT}, \gamma_{Fe\alpha}, \gamma_{Fec}, \gamma_{FeFe}$	Peak sputtering yield of carbon and iron by incident deuterium, tritium, alpha, carbon, and iron.
$\gamma_{cn}, \gamma_{Fen}$	Sputtering yield of carbon and iron by incident neutrons.
γ_{ch}	Chemical sputtering coefficient of carbon.
a, R	Plasma minor, major radius.
$N_e(r)$	Radially dependent electron density. Peak value is $N_e(0)$.
$\overline{\sigma v_{cx}}$	Charge exchange reaction rate.
$T_e(r)$	Radially dependent electron temperature. Peak value is $T_e(0)$.
$N_0(r)$	Radially dependent neutral atom density.
T_i, T_e	Average ion and electron temperature.
$U_{\alpha i}, U_{\alpha e}$	Fraction of fusion produced alpha energy delivered to the ions and electrons (function of T_e).
$G_{iD}, G_{iT}, G_{eD}, G_{eT}$	Fraction of injected deuterium and tritium beam power delivered to the ions and electrons.
$\tau_{i,E}, \tau_{e,E}$	Ion and electron energy containment time.
$\ln \Lambda$	Coulomb logarithm.
Z_j, A_j	Charge and atomic weight of species j.
E_{Dc}, E_{Tc}	Energy of cold deuterium and tritium fuel.
Z_{eff}	Effective ion charge ($\sum_j Z_j^2 N_j / N_e$).
B_T	Toroidal magnetic field (on axis).
β	Ratio of plasma pressure to toroidal magnetic field pressure.
R_e	Reflectivity coefficient for synchrotron radiation.
η	Plasma resistivity.
J	Plasma current density.
P_{LINE}	Impurity line and recombination radiation.

In order to identify the physics included in the code and the models used, a brief description of selected terms is given below. All terms are described in Ref. 1.

2.1.1 Beam-Plasma Fusion Power

The fusion power produced by beam-plasma interactions is expressed as a multiple, F , of the injected beam power as discussed by Dawson et al.³ For the case of deuterium injection,

$$F = F(E_b, T_e) N_T / N_e$$

where

$$F(E_b = 150 \text{ keV}, T_e) = -0.12759 + 0.27353 T_e - 0.00879 T_e^2.$$

For 200 keV, the coefficients change only slightly, to -0.12893, 0.2747, and -.00752, with $F = 2.3$ for $T_e > 18$ keV. F is taken to be zero for $F < 0$ and $F = 2$ for $T_e \geq 15$ keV. The fusion energy is 17.6 MeV and the 3.52-MeV alpha particles deliver their energy to the plasma as described above. The density change due to beam-plasma fusion is given by

$$= \frac{S_D E_b}{17600} \frac{N_T}{N_e} F(E_b, T_e) .$$

2.1.2 Neutral Beam Injection Heating

The energy delivered to the background plasma electrons and ions by fast injected deuterium and tritium beams is given by the functions G_e and G_I .⁴ The critical energy, E_c , i.e., the energy at which the fast ion slowing down rate on background electrons and ions is equal, is given by

$$E_c = 14.8 T_e \left(\frac{M_f}{M_H} \right)^{1/3} \left(\frac{M_f(Z)}{M_i} \right)^{2/3}$$

$$(Z) = \frac{\sum_j N_j Z_j^2 (M_i/M_j)}{\sum_j N_j Z_j} .$$

This quantity and the approximate fit to the G_e , G_I curves, neglecting charge exchange,

$$G_I = 0.5 \exp(-0.0916 E_b/E_c) \\ + 0.5 \exp(-0.635 E_b/E_c) ,$$

$$G_e = 1 - G_I$$

permit the necessary fractions to be computed.

2.1.3 Containment Times

The particle and energy containment times used are critical in determining the detailed energy balance for a system. Theoretical estimates must be used in this regard since at present scaling laws appropriate to large, hot plasmas are not known. Here scaling laws are taken from Ref. 5. They include both convection and conduction losses. The containment times are derived from the appropriate transport coefficients as follows:

$$\tau_d = a^2/4D ,$$

and

$$\tau_x = a^2/4\chi ,$$

where the density scale length r_N and the temperature scale length r_T have been set equal to the minor radius a . Quasi-neutrality requires that the convection times for ions and electrons be equal. Thus, $\tau_d = \tau_{d,e} = \tau_{d,i}$. The electron and ion thermal conductivity times, $\tau_{x,e}$ and $\tau_{x,i}$, may be different. The table below indicates the transport modes included in the analysis, and the symbols used to represent them.

	<u>Ion Conduction</u>	<u>Electron Conduction</u>	<u>Convection</u>
Pseudoclassical	—	$\tau_{X,e}^{PS}$	τ_d^{PS}
Neoclassical	$\tau_{X,i}^{NEO}$	—	—
First Trapped Electron Mode	—	$\tau_{X,e}^{TEM 1}$	$\tau_d^{TEM 1}$
Second Trapped Electron Mode	—	$\tau_{X,e}^{TEM 2}$	$\tau_d^{TEM 2}$
Trapped Ion Mode	$\tau_{X,i}^{TIM}$	$\tau_{X,e}^{TIM}$	τ_d^{TIM}

Reference 1 includes a detailed discussion of each mode. The formula used to represent each mode is given as well as the appropriate correction for cases where the plasma is noncircular. In order to use the confinement time formulations in the energy balance, it is necessary to combine the various modes into confinement times for particles, ion energy, and electron energy. This is accomplished as outlined in the following:

$$\frac{1}{\tau_{CONV}} = \frac{1}{\tau_d^{PS}} + \frac{1}{\tau_d^{TEM}} + \frac{1}{\tau_d^{TIM}}$$

$$\frac{1}{\tau_{COND,i}} = \frac{1}{\tau_{X,i}^{NEO}} + \frac{1}{\tau_d^{TIM}}$$

$$\frac{1}{\tau_{COND,e}} = \frac{1}{\tau_{X,e}^{PS}} + \frac{1}{\tau_d^{TEM}} + \frac{1}{\tau_d^{TIM}}$$

The containment times used in the rate equations are:

$$\tau_d = \tau_D = C_T \tau_T = C_\alpha \tau_\alpha = C_c \tau_c = C_{Fe} \tau_{Fe} = \tau_{CONV}$$

$$\frac{1}{\tau_{E,i}} = \frac{1}{\tau_{COND,i}} + \frac{T^*}{\tau_{CONV}}$$

$$\frac{1}{\tau_{E,e}} = \frac{1}{\tau_{COND,e}} + \frac{T^*}{\tau_{CONV}}$$

The C factors permit consideration of different ion containment times for various ion species. The T^* factor is the fraction of the average ion energy assumed to be carried out of the system by particles which diffuse from the containment volume and are lost to a surface.

2.1.4 Alpha Energy Deposition

The fraction of the alpha energy produced which is delivered to the ions, $U_{\alpha I}/U_0$, is taken to be $T_e(\text{keV})/50(\text{keV})$. The fraction delivered to the electrons is $1 - U_{\alpha I}/U_0$. These expressions were derived from curves developed by Sigmar and Joyce⁶ and are accurate for $T_e \leq 20$ keV. It is important to know the relative fraction, since different T_e and T_i may result, especially at low density and high temperature.

2.1.5 Impurity Radiation Effects

This topic is covered in detail in Sect. 3.

2.2 Point Modeling — Static Code

A static code called POINT has been written to solve the detailed energy balance equations for a tokamak system at equilibrium. The physics included is the same as that described in Sect. 2.1 and in Ref. 1. The POINT code is used interactively to obtain sensitivity plots for selected variables as a function of system parameters. The calculation is fast, which permits the large field of variables to be surveyed, and the results are good approximations to the time-dependent equilibria. The following simplifying assumptions are made: $T_e = T_i$; the equilibrium alpha density is zero; and ohmic heating is neglected. Input parameters which must be specified include the aspect ratio, toroidal field, safety factor $q(a)$, effective charge Z_{eff} , fraction of ion density composed of deuterium charge of low- and high-Z impurities, ratio of low-Z density to electron density, thermal conversion and neutral beam efficiencies, and N_τ . If two of the following three parameters are given, then the third can be computed: electron density N_e , temperature T , and the ratio of electron pressure to poloidal field pressure, $\beta_{\theta e}$. A graphical display program allows cross-plotting of virtually any two variables.

With these parameters, the program computes: the fusion reaction rate, stored energy, bremsstrahlung, line radiation losses, and the diffusion loss power. The deuterium-tritium ratio which optimizes the specific power produced, the injected power required, the beam-plasma fusion power, and the background fusion power produced are also found. The density and temperature are assumed to be uniform. A strong peaking of these quantities could lead to a significant change in background fusion power.

The STATIC code has been used extensively to examine the sensitivity of assumed reference cases to the variation of critical parameters. The results are usually displayed graphically as shown in Figs. 2.1-2.3.

Figure 2.1 displays information over the N_τ region of interest. Since N_τ is used as the independent variable, this figure has absolutely no inherent implicit scaling assumptions, and the information is valid for tokamaks of arbitrary size. Typical values have been chosen for the parameters $Z_{eff} = 2$, aspect ratio = 3, $q = 2.5$, and $B_t = 4$ T. Figure 2.1 is for operation at a temperature of $T = 8$ keV. The quantity P_b/V is the beam power per unit volume necessary to provide equilibrium operation in those cases where N_τ is less than that required for ignition. Note that P_b/V is sensitive to the assumed value of $\beta_{\theta e}$. The quantity P_α^M/P_α is the ratio of the Maxwellian or background-produced fusion alpha power to the total alpha power produced (which also includes beam-plasma fusion events). This ratio is independent of $\beta_{\theta e}$. If $P_\alpha^M/P_\alpha = 1$, then no beam-injected energy is required to maintain equilibrium. This corresponds to ignition, and it occurs in the present case at about $N_\tau = 10^{15} \text{ cm}^{-3}$ -s and $T = 8$ keV.

Figure 2.2 covers the same N_τ interval as Fig. 2.1. Again N_τ is taken to be the independent variable but assumptions are made relative to confinement time scaling laws in order to compute the plasma size necessary to maintain $T = 8$ keV for the assumed reference case parameters. The minor radius required to attain various N_τ values is shown for the scaling assumptions given in Sect. 2.1 and also for cases where the trapped-particle confinement times factored into the overall energy balance are improved by factors of $\sqrt{10}$ and 10. The latter may be considered as an effort to reflect optimism, an indication that trapped-particle modes are profile dependent and are predicted using a "worst case" analysis, or simply as a sensitivity calculation. This choice is left to the reader.

ORNL-DWG 76-10846

Fig. 2.1. Beam power per unit volume and the ratio of Maxwellian alpha power to total alpha power as a function of $N\tau$ for several $\beta_{\theta e}$ at $T = 8 \text{ keV}$.

ORNL-DWG 75-11170

TRAPPED PARTICLE LOSS VARIATION

 $T = 8 \text{ Kev}$

Fig. 2.2. Plasma radius required to attain temperature of 8 keV for different trapped particle loss coefficients.

Fig. 2.3. Beta poloidal and plasma radius as a function of $N\tau$ for different toroidal magnetic field strengths.

Note that an uncertainty of a factor of 10 in the trapped-particle confinement time at $N\tau = 10^{15} \text{ cm}^{-3} \text{ -s}$ results in an uncertainty of a factor of 2 in the plasma size required to attain that $N\tau$. Similarly, for $a = 2 \text{ m}$, the range shown for τ_{TP} results in a range of achievable $N\tau$ which includes $8 \times 10^{13} - 6 \times 10^{14} \text{ cm}^{-3} \text{ -s}$.

Figure 2.3 is similar to Fig. 2.2, except that the total poloidal beta is displayed, including the contribution from the fast-injected beam and alpha particles. This computation involves an integration over the energy range of the particles as they slow down. In the case shown, the sensitivity of size to changes in magnetic field strength on axis is displayed. Note that β_p is not sensitive to B_T , but that the size required to attain a particular $N\tau$ is strongly dependent on field strength when $N\tau \geq 10^{14} \text{ cm}^{-3} \text{ -s}$.

Figure 2.4 is an example of a useful technique used to determine the size of a tokamak fusion reactor. The minor radius, $a(\text{m})$, is plotted versus toroidal magnetic field on axis, B_t , for various values of B_{max} , the peak toroidal field at the toroidal windings. Values of $A = 3$ and $t_{bs} = 2 \text{ m}$ have been assumed for the aspect ratio and blanket thickness, respectively. In addition, ignition curves are shown for various ignition and operating temperatures.

For each value of T and B_t , the $N\tau$ value required for ignition is computed. The selected scaling law is then inverted to find the smallest minor radius which is consistent with operation at that $N\tau$. For clarity, only three different temperatures are displayed on Fig. 2.4, but any point on the plot above the 10 keV curve is a solution for two different temperatures. This is a result of the fact that $N\tau$ increases with temperature in the second trapped electron regime, but decreases with temperature once the trapped-ion regime dominates transport behavior. Therefore, there are two temperatures at which plasma heating, due to the energy input from fusion alpha particles, exactly balances plasma energy losses. This is illustrated graphically in Fig. 2.5. At point 1 the $N\tau$ achieved in a device of radius a_2 is just sufficient for ignition. A machine of radius a_1 will not ignite at any temperature. A machine of radius a_3 is ignited for temperatures between points 2 and 3, and reaches a stable equilibrium at point 3. Thus the system with plasma radius a_3 will ignite at point 2, will undergo a thermal excursion to point 3 since P_α is greater than P_{loss} following ignition, and will operate stably at point 3.

For example, in Fig. 2.4 with $B_{\text{max}} = 11 \text{ T}$, the minimum plasma radius for ignition is 2.3 m. Ignition would occur at $T \approx 10 \text{ keV}$, corresponding to operation at point 1 on Fig. 2.5.

2.3 Reactor Modeling — Time-Dependent SYSTEMS Code

The functional dependences of the transport laws which will govern reactor plasma dynamics are not known at present. The detailed energy balance which will dictate the parameters attained in the plasma is better understood than previously, but many uncertainties remain. All of these things will be clarified as the experimental and theoretical research programs progress. In parallel with these efforts, it is necessary to assess the implications of various possible outcomes of this research. The models previously described are used in this section to predict the size and operating characteristics of several reactor-scale devices. The discussion will specifically center around the ORNL Experimental Power Reactor Reference Design. It is assumed here that the EPR is a low- β ignition device that will produce net power during a single operating cycle. The ignition requirement immediately imposes constraints on the $N\tau$ which must be achieved by the system. Figure 2.6 shows the beam power per unit volume which is required in the model to maintain the energy balance as a function of $N\tau$. The temperature assumed is 8 keV; $B_T = 4 \text{ T}$, $\beta_{\theta e} = 1.0$, and $A = 3$. When the beam power required is zero, ignition has been achieved and the thermonuclear alpha particles produced maintain the energy

ORNL-DWG 76-10847

Fig. 2.4. Minimum minor radius required for ignition at given B_T and T ($1/10 \times$ T.P. losses).

Fig. 2.5. The density-confinement time product as a function of temperature for several plasma radii, and the ignition condition.

ORNL-DWG 76-4334

Fig. 2.6. Beam power per unit volume and the ratio of Maxwellian fusion power to total fusion power as a function of $N\tau$ for several Z_{EFF} at $T = 8 \text{ keV}$.

balance. The ignition point for various Z_{eff} values is more clearly shown on the ordinate labeled P_F^M/P_F . The total fusion power produced is a sum of that due to beam-plasma fusion events and that due to background or Maxwellian events. When the Maxwellian fusion power, P_F^M , is equal to the total fusion power, P_F , ignition has been achieved. Thus, in the present example, $N\tau = 5 \times 10^{14} \text{ s/cm}^3$ for ignition at $T = 8 \text{ keV}$ and $Z_{\text{eff}} = 1$. At $Z_{\text{eff}} = 2$, $N\tau = 10^{15} \text{ s/cm}^3$; at $Z_{\text{eff}} = 3$, ignition cannot be achieved. For the EPR a low-Z impurity (carbon) and $Z_{\text{eff}} = 2$ are assumed.

Figure 2.7 shows curves similar to those on Fig. 2.6 except that various operating temperatures are assumed. Operation at increased temperature reduces the $N\tau$ required for ignition but does not necessarily reduce the size of the device required to achieve the $N\tau$ itself. This again depends directly on the scaling laws and their functional dependences. Here the model is that described in Sect. 2.1 with the trapped-particle mode coefficients (convection and conduction) assumed to be one-tenth those estimated by linear stability analyses.⁵ Figure 2.2 shows the impact of varying the trapped-particle mode coefficients, and Fig. 2.3 shows the effect of varying the magnetic field on axis.

Performing the trade-offs indicated above results in the basic EPR machine parameters given in Table 2.1. It should be emphasized that this device is based upon the scaling described in Sect. 2.1, with stable plasma operation at $q = 2.5$, low- β equilibria, and a maximum magnetic field strength at the toroidal field coil superconducting winding of 11 T. It should also be recognized that radial profiles do not significantly alter the global power balance.

Table 2.1. EPR machine parameters

Major radius, R_0	6.75 m
Plasma radius, a	2.25 m
Field strength on axis, B_T	4.8 T
Field strength at winding, B_{max}	11.0 T
Wall area, S	600 m^2
Plasma volume, V	676 m^3
Safety factor, q	2.5
Plasma current, I	7.2 MA
Beam power, P_b	50 MW
Beam energy, E_b	200 keV

The above considerations established a size for EPR. An assessment of EPR performance involves many parameters, and the results are very sensitive to the assumptions employed. Based on the values in Table 2.1, an initial temperature of 100 eV ($T_e = T_i$), and an initially clean plasma, attainable steady-state operating points for the EPR were determined using the dynamic SYSTEMS code which solves the equations in Sect. 2.1 above. The key parameter difference in each case is initial density. Figure 2.8 shows the results for thermal power as a function of the operating density. For example, consider the curve labeled "Full losses, $P_b = 100 \text{ MW}$." These cases correspond to driven operation where 100-MW beams, together with the alpha power, balance the plasma losses; i.e., ignition is not achieved since full trapped-particle losses are assumed and the device was not sized for ignition under those circumstances.

Thus power outputs of between 100 and 440 MW are possible when the operating density is varied between $2 \times 10^{13} \text{ cm}^{-3}$ and $1 \times 10^{14} \text{ cm}^{-3}$, respectively. The corresponding β_{0e} and T for the various

ORNL-DWG 76-4333

Fig. 2.7. Beam power per unit volume and the ratio of Maxwellian fusion power to total fusion power as a function of $N\tau$ for several temperatures at $Z_{EFF} = 2$.

Fig. 2.8. EPR steady-state operating points, P_{thermal} vs N_e .

operating points are shown on Figs. 2.9 and 2.10. Note from Fig. 2.10 that in the driven cases T is approximately 9 keV over a wide range of operating densities. This is due to the fact that the dissipative trapped-ion mode confinement time is proportional to $\beta_{\theta e}^2/T^{11/2}$. The temperature dependence strongly dominates the scaling, so that the driven machine with $P_b = 100$ MW can only achieve $T \sim 9$ keV. Since $\beta_{\theta e} \propto N_e T_e$ and T_e is approximately constant, $\beta_{\theta e}$ varies linearly with density as shown on Fig. 2.9. The output power in the driven cases varies approximately as the square of the density (Fig. 2.8). This also follows from $P(th) \propto \beta^2 \propto N^2$. The driven cases are uninteresting because they are unlikely to yield net power or commercial power.

Figure 2.8 also shows operating points for cases where the trapped-particle mode losses are reduced by a factor of 10 relative to those typically assumed. In the lower density cases, the beams operate continuously and thermal powers between 100 MW and 400 MW are possible. In accordance with the reasoning above, the operating temperature again is about constant ($T \sim 14.5$ keV). See, for example, Fig. 2.10. The ignition curve is of interest here. At densities above approximately $6 \times 10^{13} \text{ cm}^{-3}$, ignition is achieved by injecting 50 MW of neutral deuterium atoms into the plasma for about 5 sec. After ignition the beam power is terminated and the plasma system attains the steady-state parameters shown in Figs. 2.8-2.10, again depending on the density of operation. Thermal power outputs between ~ 200 and ~ 800 MW, temperatures between ~ 10 and ~ 18 keV, and β_p between ~ 1 and ~ 3 are typical parameter ranges which could be achieved in the EPR.

The parameters for the chosen operating point for the EPR are shown in Table 2.2. In this case, it has been assumed that the plasma is initially free of impurities. Sputtering occurs according to the model described in Sect. 2. It is also assumed that the neutral density at the plasma boundary is $10^{15}/\text{m}^3$ and that the density decreases exponentially with a characteristic length of 5 cm. This permits charge-exchange sputtering to be included in the modeling. An impurity control scheme is assumed to be 99% effective with respect to charged particles and 90% effective with respect to neutral charge-exchange sputtering, i.e., the return impurity flux is assumed to be attenuated by 90%. These models and assumptions result in an increase in Z_{eff} to 1.34 after 50 sec of plasma burning. This is tolerable, and probably corresponds to the degree of impurity control required in the case of a stainless steel wall if the impurities entering the plasma mix throughout the reacting volume on a time scale short compared to a burn time. The effect of low-Z walls and a more general description of burn cycle limitations due to sputtering are given in the next section.

Table 2.2. Steady-state plasma parameters and power balance

Electron density, N_e	$7.4 \times 10^{19} \text{ m}^{-3}$
Electron temperature, T_e	13.0 keV
Ion temperature, T_i	12.2 keV
Energy confinement time, τ_E	3.7 sec
Particle confinement time, τ_p	9.6 sec
Poloidal beta, β_p	1.9
Thermal power, P_{th}	410 MW
Neutron wall loading, P_M/S	0.55 MW/m^2
Impurity level, Z_{eff}	1.34 ($t = 50$ sec)
Neutral density ($r = a$), N_0	10^{15} m^{-3}
Characteristic length, λ	5 cm

Fig. 2.9. EPR steady-state operating points, $\beta_{\theta e}$ vs N_e .

ORNL-DWG 76-10855

Fig. 2.10. EPR steady-state operating points, T_i vs N_e .

Alpha power, P_{α}	+82 MW
Transport power, P_{tr}	-66 MW
Bremsstrahlung power, P_{br}	-9 MW
Line radiation power, P_{ℓ}	-6 MW
Synchrotron power, P_{synch}	-1 MW
$\sum_k P_k = 0$ (ignition)	

REFERENCES

1. S. E. Attenberger, F. B. Marcus, and D. G. McAlees, *Dynamics of Tokamak Plasma Experiments and Reactors*, ORNL/TM-5509, Oak Ridge (to be published).
2. M. Roberts et al., *Oak Ridge Tokamak Experimental Power Reactor Study Reference Design*, ORNL/TM-5042, Oak Ridge (November 1975).
3. J. M. Dawson, H. P. Furth, and F. H. Tenney, "Production of Thermonuclear Power by Non-Maxwellian Ions in a Closed Magnetic Field Configuration," *Phys. Rev. Lett.* 26: 1156 (1971).
4. J. A. Rome et al., *Neutral Beam Injection into a Tokamak: Particle Orbit Loss-Regions in Tokamaks and Their Effects on Plasma Heating*, ORNL/TM-4855, Oak Ridge (April 1975).
5. S. O. Dean et al., *Status and Objectives of Tokamak Systems for Fusion Research*, WASH-1295, USAEC, Washington, D.C. (1974).
6. D. J. Sigmar and G. Joyce, "Plasma Heating by Energetic Particles," *Nucl. Fusion* 11: 447 (1971).

3. IMPURITY MODELING AND EFFECTS ON SYSTEM DYNAMICS

The plasma simulation model includes treatment of losses from line radiation, radiative recombination, and dielectronic recombination.¹ The model used is based on theoretical calculations by Merts et al.² The calculation is for iron and takes into account that, depending on plasma parameters, the iron may be partially ionized in some regions of the plasma. The results are dependent on plasma electron temperature. The impurity radiation loss is found to be approximately constant below a temperature T_e^L , and again approximately constant (with reduced magnitude) above a temperature T_e^V . A linear interpolation is used to connect the two regions. This radially dependent loss process is included in the model.

The electron temperature and electron density profiles are assumed to be centrally peaked parabolic functions of r (minor radius). The widths of these profiles are fixed by the initial conditions, namely the initial values for the average temperature and density and their assumed values at the plasma edge. The profiles are then determined. The peak values are re-computed as a function of time, since the code follows the evolution of all plasma parameter averages.

The high-Z impurity density is treated similarly, but the peak of the profile is allowed to occur at arbitrarily chosen radii. In particular, the impurity density may be peaked at the wall. This model allows impurities to build up at the plasma edge or at the plasma center, so that the implications of various impurity diffusion theories may be assessed.

With the above assumptions, the total radiation integral over the plasma volume is an eighth-order polynomial function of r^L and r^V , the radii at which the temperatures T_e^L and T_e^V occur. At each time step, these radii and the corresponding total radiation are evaluated.

Figures 3.1 and 3.2 show typical profiles for electron density, electron temperature, impurity density, and atomic radiation. The two cases differ in that the impurity profile is centered at the plasma edge in one case and at the plasma center in the other.

The high-temperature region, the transition region, and the low-temperature region can be distinguished between $r = 0$ and $r = a$. The remainder of the shape is due to the volume element variation with r and the density distributions used.

The equations used for the impurity radiation loss are:

$$P_{\ell}(\text{w/m}^3) = C(T_e) \times 10^{-32} N_e(\text{m}^{-3}) N_{Fe}(\text{m}^{-3}), \quad (1)$$

where

$$C(T_e) = 4.1 \quad \text{when} \quad T_e \leq T_e^L$$

and

$$C(T_e) = 0.32 \quad \text{when} \quad T_e \geq T_e^V;$$

$$C(T_e) = (4.1 - 3.78) \frac{T_e - T_e^L}{T_e^V - T_e^L} \quad \text{when} \quad T_e^L \leq T_e \leq T_e^V. \quad (2)$$

For iron, $T_e^L = 1.0$ keV and $T_e^V = 2.1$ keV.

The buildup of the carbon and iron impurity densities due to sputtering from the walls is considered. The sputtering processes included in the N_C and N_{Fe} equations are due to incident particles which result from diffusion, i.e., deuterons, tritons, alphas, carbon and iron ions; fusion neutrons and neutrons incident from the blanket; fast neutral atoms from charge-exchange; and chemical sputtering of carbon.

Fig. 3.1. Line and recombination radiation losses (integrated over θ and ϕ) vs r .

Fig. 3.2. Line and recombination radiation losses (integrated over θ and ϕ) vs r .

The ions that diffuse from the plasma are assumed to be monoenergetic. The dependence of the sputtering yield on incident particle energy is given by³

$$f(T) = \frac{4(T - E_C)^{3/2} E_M^{1/2}}{E_M^2 + 3T^2}, \quad E_M = 6 \text{ keV}, \quad E_C = 0.07 \text{ keV}. \quad (3)$$

The magnitudes of the sputtering coefficients depend on the incident species.

The neutron sputtering rate is proportional to the neutron production rate in the plasma due to background and beam-plasma fusion events. The neutron sputtering coefficient is taken from Ref. 4. The result is actually for a 0.5-cm niobium wall and a lithium-cooled blanket. It includes the 14.1-MeV source neutrons, and the random neutron flux at all energies which results from interactions in the blanket. The charge-exchange flux to the wall is computed from a spatial integration over the plasma. As in the radiation case above, radially dependent ion and neutral profiles are used.

The sputtering yield per unit volume is

$$\frac{1}{V} \int_V \gamma f(T(r)) \sigma_{CX} \frac{N_0}{4} (\gamma_D N_D v_D + \gamma_T N_T v_T) dV,$$

where $\gamma = \gamma_C$, or γ_{Fe} when appropriate. The neutral density is assumed to fall exponentially, as given by

$$N_0(r) = N_0(a) e^{-(a - r)/\lambda},$$

where λ is the characteristic length for neutral attenuation.

Chemical sputtering due to the interaction of deuterium and tritium with carbon on the wall is included. The sputtering coefficient is taken from data⁵ for carbon temperatures assumed to be greater than 1000°C. The sputtering coefficients, γ , are shown in Table 3.1.

Table 3.1. Peak sputtering yields

Carbon and Iron	
Deuterons	0.04
Tritons	0.06
Alphas	0.64
Carbon	13.0
Iron	13.0
n	3.0×10^{-5}
chemical	0.03 (carbon only)

Figure 3.3 summarizes the technique used to incorporate sputtering and impurity effects into the plasma simulation model. Spatially averaged time-dependent parameters are computed and advanced in time in the program. As part of the calculation, spatial profiles for the plasma parameters, impurity species and neutrals in the system are assumed. The profiles are chosen to have average values consistent with those in the code. Using the spatially dependent information, charge state distributions and line and recombination radiation losses are calculated. The results are integrated over the plasma volume and the integral results are used in the spatially independent energy balance computation. Similarly, the charge-exchange flux, which is dependent on space and energy, can be computed from the assumed profiles. This source of particle flux, together with the flux resulting from particle

ASSUMED PROFILES ARE NORMALIZED
TO AVERAGE VALUES

Fig. 3.3. Incorporation of sputtering and impurity effects in the plasma simulation model.

diffusion, is coupled to the wall by the sputtering model previously discussed. The diffusion flux is assumed to have a temperature of 100 eV. Sputtered impurities are assumed to reenter the plasma with some efficiency and on a profile which is selected *a priori*. It should be noted that major uncertainties exist in this area. For example, impurity transport and therefore the density profiles are not well understood; sputtering models (especially at low energies) remain to be verified; and detailed neutral behavior is an active area of concern. Using the impurity model, impurity control criteria and system responses to possible outcomes are described.

A reference case is established using the EPR machine parameters given in Table 2.1 of Sect. 2. The initial density is $.75 \times 10^{20} \text{ m}^{-3}$, the neutral density at $r = a$ is assumed to be 10^{15} m^{-3} , and the neutral profile decreases exponentially into the plasma with a 5-cm characteristic length. The impurity confinement time is assumed to be equal to that for the deuterons, and the wall is taken to be stainless steel (iron for calculational purposes). It should be noted that the reference case assumes an impurity level consistent with a "divertor efficiency" of 90%, even though there is no divertor in the EPR design. This efficiency is the fraction of the charged-particle flux diffusing toward the wall that does not sputter. Also, 90% of the impurities sputtered do not reenter the plasma. If a divertor is ultimately chosen as the means of impurity control, then this model requires that it must remove 90% of a charged particle flux which crosses the scrape-off zone. In this case, some means must be provided to re-inject a corresponding flux of fuel if the burn time is not to be limited. Alternatively, the efficiencies establish a target for other impurity control schemes. The reference case, curve 1 in Fig. 3.4, achieves a burn time of 100 sec.

Sensitivity calculations shown as curves 2-4 in Fig. 3.4 indicate results for (1) no flux attenuation, (2) 70% attenuation but impurity accumulation during the burn ($\tau_{\text{imp}} = \infty$), and (3) an increased neutral edge density which enhances sputtering due to charge exchange. In all cases, burn time is limited to less than 60 sec.

The reactor control model used here leads to the discontinuities which appear in the curves. After about 5 sec, ignition occurs and the injected beam power is discontinued where $T_i = 10 \text{ keV}$. The temperature continues to increase, and when β_0 approximates to A, the cold fueling rate which is used to maintain the plasma density is decreased. This is considered to be a realistic reactivity control. This leads to the abrupt temperature change near $t = 10 \text{ sec}$. When β_0 is less than A, fueling continues to maintain the fuel ion density. If the temperature decreases to 10 keV, the beams are turned on in an attempt to maintain the temperature within the range of EPR practicability. For example, curve 3 in Fig. 3.4 shows the effect of impurity accumulation. The fuel ion density decreases as a result of the buildup of impurities. The reactivity therefore decreases and the temperature drops. When a level of 10 keV is reached, the beams are used but continued fuel-ion decrease prevents the burn from continuing beyond $\sim 50 \text{ sec}$.

Curves 5 and 6 in Fig. 3.4 show the results when a carbon or low-Z wall is assumed. In curve 5, $\tau_{\text{imp}} = \tau_{\text{ion}}$; after initial transients, ignition is reestablished at $t = 50 \text{ sec}$ and a long burn is found to be possible. When $\tau_{\text{imp}} = \infty$, again the burn is limited to $\sim 50 \text{ sec}$ even in the low-Z case.

Figure 3.5 shows the model used for sputtering by deuterons. For cases where carbon is assumed, the chemical sputtering yield⁵ is taken to be constant. In the curve labeled "O.R.," the particle sputtering yields are energy-dependent³ and in the same ratios as those given by Behrisch⁴ for Nb. The absolute values are normalized to those for deuterons (0.04). A numerical formulation or fit to the recent sputtering data from Behrisch et al.⁶ was obtained from Conn⁷ and is shown for comparison.

In the modified EPR reference case (shown in the lower portion of Fig. 3.5), sensitivity to the sputtering model assumed is apparent. The diffusion flux is assumed to impinge on the wall with an energy of 100 eV. Sputtering due to this source is tolerable as indicated by the burn curve labeled

1. REFERENCE CASE: DIVERTOR EFF. = 90%, $\tau_{imp} = \tau_D$,
 $N_0(\text{edge}) = 1 \times 10^{15}/\text{m}^3$, STAINLESS STEEL WALL
2. NO DIVERTOR
3. DIVERTOR EFF. = 70%, $\tau_{imp} = \text{INFINITY}$
4. $N_0(\text{edge}) = 5 \times 10^{15}/\text{m}^3$
5. NO DIVERTOR, CARBON LINER
6. NO DIVERTOR, CARBON LINER, $\tau_{imp} = \text{INFINITY}$

Fig. 3.4. Time dependence of temperature for various impurity control conditions.

ORNL-DWG 76-10854

Fig. 3.5. Sputtering coefficient as a function of incident deuteron energy and time dependence of temperature for two sputtering models.

"Conn" which represents calculational results obtained using that sputtering model. The O.R. model limits the burn to less than 40 sec, which indicates that the dominant sputtering source results from charge exchange at energies greater than ~ 1.5 keV where the sputtering coefficient in the O.R. model exceeds that of Conn for the iron model.

Finally, Fig. 3.6 shows similar results for the same reference EPR except for the case where the alpha particles accumulate. The lower curves show the density behavior. Note that when $\tau_\alpha = \infty$ for burn times less than ~ 120 sec, fueling maintains the fuel ion density approximately constant. However, the alpha and electron densities are increasing. Losses also increase during this period and the temperature decreases, resulting in a reactivity decrease. After $t = 120$ sec, $T \sim 10$ keV has occurred and the beams are turned on to maintain a 10-keV temperature. The deuteron density is approximately constant for $t > 120$ sec due to injection. However, to maintain $\beta < A$, the triton density must be decreased to accommodate the increasing alpha number density. The reactor power history indicates that for $\tau_\alpha = \infty$, the burn time will be limited to ~ 100 sec. This assumes, of course, that $\beta_p < A$ is a meaningful limit on plasma pressure.

REFERENCES

1. E. C. Crume, Jr. (ORNL), private communication, December 1975.
2. A. L. Merts, R. D. Cowan, and N. H. Magee, Jr., *The Calculated Power Output from a Thin Iron-Seeded Plasma*, LA-6220-MS, Los Alamos, New Mexico (1976).
3. D. M. Meade et al., "The Effects of Impurities and Magnetic Divertors on High-Temperature Tokamaks," *Plasma Physics and Controlled Nuclear Fusion Research*, IAEA-CN-33/A15-4, Vol. I, p. 605, Tokyo; 1974.
4. R. Behrisch, *Nucl. Fusion* 12: 481 (1972).
5. M. I. Guseva and V. M. Gusev, *Conf. on Engineering Problems of Thermonuclear Reactors*, Leningrad (July 1974).
6. R. Behrisch et al., "Measurements of the Erosion of Stainless Steel, Carbon, and SiC by Hydrogen Bombardment in the Energy Range of 0.5-7.5 keV," to be published in *J. Nucl. Mater.*
7. R. W. Conn (University of Wisconsin), private communication, 1975.

ORNL-DWG 76-10852

Fig. 3.6. Time-dependent fusion thermal power and density for two assumptions on alpha-particle containment.

4. POLOIDAL FIELD SYSTEM DESIGN

The tokamak EPR has superconducting toroidal field coils, and its plasma is expected to reach high β_p due to thermonuclear ignition. It is required to have a burn time of around 100 sec for economic operation. Because superconducting TF coils are used, it is desirable to have a poloidal field system which minimizes the pulsed fields that intersect the TF coils, permitting attainment of the highest toroidal field with the least risk of superconductor quenching. High β_p values are more easily obtained by closely coupled vertical field coils that can respond quickly to large increases of plasma β . Also, the attainment of long pulse time is aided by a reduction in the stored poloidal field energy and the number of volt-seconds required to drive the plasma current.

A magnetic field shielding system for the ORNL EPR reference design has been proposed to fulfill these requirements.¹ Its concept, implementation, and implications are the central theme of this section. As distinct from the conventional poloidal field system, such as in TFTR, where the OH (ohmic heating primary) and VF (vertical field) windings are wrapped around the outside of the TF coils, the proposed shielding system introduces a set of S-VF (shielding-vertical field) windings inside the TF coils and possibly a set of decoupling windings outside and a set of T-VF (trimming-vertical field) windings at the locations of the conventional VF windings. A comparative study of the two systems for a circular plasma is presented in Sects. 4.4 and 4.5.

The MHD equilibrium and stability requirements are discussed briefly in Sect. 4.1, followed by a description of the field shielding concept in Sect. 4.2. Its implementation is presented in Sect. 4.3. Its implications in terms of pulsed poloidal fields and volt-seconds and power supply requirements are presented in Sects. 4.4 and 4.5, respectively.

4.1 Equilibrium and Stability Requirements

The vertical field on the midplane ($R, Z = 0$) required to hold a circular plasma can be estimated by the formula^{2,3}

$$B_V(R) = \frac{I_p}{cR} \left(\ln \frac{8R}{a} + \beta_p + \frac{\ell_i}{2} - \frac{3}{2} \right), \quad (1)$$

where I_p is the toroidal current, a is the minor radius, $\beta_p \propto \bar{P}/I_p^2$, and ℓ_i is the plasma internal inductivity (inductance per unit length in electromagnetic units). However, this expression is obtained while assuming $A \gg 1$ and $\beta_p < 1$. Since the EPR plasma has $A = 3$ and $\beta_p = 1.9$, it is important to establish the equilibrium field requirements for free-boundary, $A = 3$, and $\beta_p \sim A$ plasmas.

4.1.1 Equilibrium Requirements

The equation for axisymmetric MHD equilibria may be written in terms of the poloidal flux ψ as⁴

$$\frac{\partial \psi}{\partial R^2} - \frac{1}{R} \frac{\partial \psi}{\partial R} + \frac{\partial^2 \psi}{\partial Z^2} = -4\pi R^2 P' - FF' = 4\pi R J_T, \quad (2)$$

where P and $F(\psi)$ are pressure and toroidal flux, respectively, J_T is the toroidal current density, and primes denote density with respect to ψ . The solutions and boundary values of ψ are obtained by an iterative procedure.⁵ The functions $P(\psi)$ and $F(\psi)$ are presented in such a fashion that P and J_T are zero beyond the innermost flux surface that intersects the limiter. For fixed plasma current and with sufficient vertical field, convergent solutions are always obtained over a wide variety of

pressure and current profiles.

Typical solutions with various β_p values in a vacuum field of decay index $n (= -\partial \ln B_v / \partial \ln R) = 0.5$ are shown in Fig. 4.1. It is seen that, as β_p approaches A, the plasma cross section becomes progressively noncircular with horizontal elongation. This will result in an inefficient use of the available volume in the plasma chamber. It is therefore of interest to determine the optimum vertical field decay index to ensure a circular plasma cross section at large β_p .

The dependence of the plasma cross section on the decay index n for $\beta_p = 2.6$ and nearly fixed values of B_v at the geometric center is shown in Fig. 4.2. It is clear that a decay index around 0.2 is appropriate for circular plasmas at β_p approaching A.

4.1.2 Stability Requirements

Plasma stability in the presence of rigid motions and ideal local modes is considered. With respect to column motions of a circular plasma column, it is sufficient to require $1.5 > n > 0$ for stability against vertical and horizontal displacements.⁶ Circular plasmas at high β_p with $n = 0.2$ are therefore stable in this aspect. However, the decay index varies over the plasma volume, so n at the plasma center should be somewhat greater than zero.

With respect to ideal local modes, the Mercier^{7,8} stability criterion is used. It can be symbolized as

$$\frac{1}{4} S^2 + F - G \geq 0, \quad (3)$$

where the three terms denote stabilization due to shear, magnetic well, and finite pressure. This criterion is numerically evaluated for a series of free-boundary equilibria that have reasonable current profiles and are centered by vertical fields with decay index $n = 0.5$.⁹ The critical β , the ratio of plasma to toroidal field pressure (averaged over volume), and the q (safety factor) values allowed by Eq. (3), are plotted as functions of β_p in Fig. 4.3. The maximum value of critical β , approximately 4%, is obtained in the region of $q_a = 3-5$, at a β_p of around 2. This constitutes one of the criteria used to specify the 3% design value of β for EPR plasma at $\beta_p \approx 1.9$. The q_a value of 3.5 at this β_p is higher than the design value for q , so flatter current profiles will not be necessary judging from the local stability considerations.

It is of interest to note that the $q_0 \geq 1$ requirement used here for stability agrees with the well-known Kruskal-Shafranov condition. Therefore, the EPR plasma is also stable against the $m = 1$, $n = 1$ kink mode.

4.2 Shielding Concept and Computing Method

The underlying concept of the magnetic shield system can be illustrated by Fig. 4.4., where two perfectly-conducting toroidal shells divide the space into three regions. In order that Region II is free from poloidal field, surface currents J_{OH} and J_S on C_{OH} and C_S , respectively, must be set up to satisfy

$$J_{OH} = |B_{OH}|/\mu_0, \quad J_S = |B_P|/\mu_0, \quad (4)$$

in a static and lossless state.

Fig. 4.1. Free-boundary equilibria with fixed I_p and $n = 0.5$. B_0 is the vertical field required to hold a plasma with $\beta_p = 0$, and B_z is the applied vertical field.

ORNL-DWG 76-10843

Fig. 4.2. Dependence of the plasma shape on the decay index n , with $\beta_p = 2.6$, $A = 3$, and $B_v(R_0) = 2.0 B_0$, where B_0 is the vertical field required to hold a plasma with $\beta_p = 0$ (see Fig. 4.1).

ORNL-DWG 76-10845

Fig. 4.3. The critical β , q_a , and q_0 as a function of β_p for stability of ideal local modes.

4.2.1 Channeling the OH (Ohmic Heating) Flux

In Region I of Fig. 4.4, it can be asserted that, for a given total current

$$I_{OH} = \oint_{C_{OH}} \mathbf{s} d\mathbf{l} \cdot \mathbf{J}_{OH},$$

there are unique static solutions of $B_{OH}(R, Z)$ and $J_{OH}(\mathbf{l})$. This is because $B_{OH} (= i_R B_R + i_Z B_Z)$ satisfies the equation,

$$\left(\frac{\partial^2}{\partial R^2} - \frac{1}{R} \frac{\partial}{\partial R} + \frac{\partial^2}{\partial Z^2} \right) \psi = 0,$$

$$B_R = -\frac{1}{R} \frac{\partial \psi}{\partial Z}, \quad B_Z = +\frac{1}{R} \frac{\partial \psi}{\partial R}, \quad (5)$$

with $\psi = 0$ at $R = 0$ and at infinity, and $\psi = \text{constant}$ on the shell C_{OH} . Thus, if J_{OH} is driven through a large number of filaments on C_{OH} , the magnetic field produced will be confined to Region I. Because J_{OH} can shield the TF coils and the plasma from the OH flux, and thus completely link the OH flux with the plasma loop, J_{OH} becomes the ideal OH current distribution in tokamaks.

4.2.2 Shielding Against the Plasma Flux

In the presence of a plasma current I_p , the eddy current J_S induced on C_S also represents the ideal current distribution that confines the plasma flux to Region III. This current has the following properties:

1. Its sum is equal and opposite to the plasma current,

$$I_S = \oint_{C_S} \mathbf{s} d\mathbf{l} \cdot \mathbf{J}_S = -I_p. \quad (6)$$

2. It produces the equilibrium vertical field, B_{VF} , that holds and shapes the plasma column.
3. The current $J_S(R, Z)$ can adjust itself to changes in I_p and its internal distribution to maintain the complete shielding.

Our goal is to determine the locations of a set of shielding coils that essentially reproduce these properties.

4.2.3 Tailoring the Vertical Field

The currents J_{OH} and J_S are such that a specific form of B field is produced in a certain region. The field shielding concept can thus be easily generalized to a field tailoring concept.

For a given plasma equilibrium enclosed in a conducting shell, a unique vertical field B_{VF} in Region III is generated by J_S . In Region II, however, J_S provides exactly the negative of the plasma field, cancelling the plasma flux external to C_S . It is therefore clear that we can break down J_S into two parts: J_{OH} and J_{VF} , where $\oint_{C_S} \mathbf{s} d\mathbf{l} \cdot \mathbf{J}_{OH} = -I_p$ and $\oint_{C_S} \mathbf{s} d\mathbf{l} \cdot \mathbf{J}_{VF} = 0$, with J_{OH} producing zero field and J_{VF} producing B_{VF} in Region III. The current J_{OH} is similar to J_{OH} on C_{OH} and is uniquely determined for a given I_p . As a result, J_{VF} becomes uniquely determined for a given B_{VF} .

This conclusion greatly simplifies the problem of calculating the shielding current J_S , in view of the fact that the approximate forms of B_{VF} are well-known. It is then relatively straightforward to determine the J_{VF} that tailors the B field in Region III to B_{VF} , and the J_{OH} that produces the most nearly zero B field in Region III. Typical forms of J_S , J_{OH} , and J_{VF} are shown in Fig. 4.5. It is seen that for a given I_p , there is a critical B_{VF} , beyond which some portion of J_S on the inside part of C_S becomes reversed.

Fig. 4.5. Shielding (eddy) current J_S ($=J'_\text{OH} + J'_\text{VF}$) on a toroidal, circular, conducting shell. The major and minor radii of the shell inner surface are 6.75 and 3.18 m, respectively. The maximum vertical field B_z at the plasma center before J_S can change sign is 0.34 T.

4.2.4 Method of Poloidal Field Coil Design

An optimal coil design aims at best reproducing the effects of J_{OH} and J_S with practical current-carrying coils that allow for adequate access to plasma and feasibility in engineering. There have been many efforts in this area using various approaches.¹⁰ We have chosen a relatively simple approach that puts emphasis on reproducing the proper vacuum fields in Regions II and III of Fig. 4.4. This approach involves two steps. First, the surface current distribution that best reproduces a prescribed field in a region is calculated. This current is then approximated with discrete coils that carry currents in small multiples of a single current amplitude.

4.2.5 Optimal Current Distribution

An example is the OH current, J_{OH} . What are the currents I_i ($1 \leq i \leq N$) in N filaments equally spaced on C_{OH} , so that the $\underline{B}$ field is minimized in Region I and II? This I_i can be determined by minimizing

$$W = \sum_{j=1}^M w_j \left[\sum_{i=1}^N I_i b_z(r_j, z_j; R_i, Z_i) - B_0/r_j, z_j \right]^2, \quad (7)$$

subject to the constraint,

$$\sum_{i=1}^N I_i = I_0/2 \quad (8)$$

In Eq. (7), $b_z(r_j, z_j; R_i, Z_i)$ is the B_z field at (r_j, z_j) due to the i^{th} pair of filaments at $(R_i, \pm Z_i)$ with unit current. The field locations (r_j, z_j) can be uniformly distributed in Regions II and III, and w allows us the freedom of emphasizing a portion of the field locations. To obtain J_{OH} , we set $B_0 = 0$ and I_0 equal to a given constant. Only the B_z field needs to be considered because B_R is small due to the up-down symmetry in geometry.

When the derivatives of W and I_N are taken with respect to I_i ($1 \leq i \leq N-1$), Eqs. (7) and (8) are changed to

$$\begin{aligned} \sum_{i=1}^{N-1} A_{ki} I_i &= B_k \quad (i = 1, \dots, N-1), \\ A_{ki} &= \sum_{j=1}^M w_j [b_z(r_j, z_j; R_k, Z_k) - b_z(r_j, z_j; R_N, Z_N)] \\ &\quad \cdot [b_z(r_j, z_j; R_i, Z_i) - b_z(r_j, z_j; R_N, Z_N)], \\ B_k &= \sum_{j=1}^M w_j [b_z(r_j, z_j; R_k, Z_k) - b_z(r_j, z_j; R_N, Z_N)] \\ &\quad \cdot [B_0(r_j, z_j) - b_z(r_j, z_j; R_N, Z_N)]. \end{aligned} \quad (9)$$

I_i is then obtained by solving Eq. (9). We found that it is adequate to use $N = 8$ to 14 for consistent results. To calculate J_{VF} , we should set $I_0 = 0$ and $B_0 = [B_{VF}(R, Z)]_z$.

4.2.6 Approximation with Discrete Coils

We can now fit $I_i(\ell_i)$ with

$$J(\ell) = \sum_{n=0}^K a_n \cos(n\pi\ell/L),$$

$$a_n = \frac{1}{N} \sum_{i=1}^N I_i \cos(n\pi\ell_i/L), \ell_i = (i - \frac{1}{2}) \frac{L}{N}, \quad (10)$$

where ℓ is the arc length along C_{OH} that has a total length $2L$, and ℓ_i is the arc length to the i^{th} equally-spaced filament (pairs). We can then determine the locations ℓ_i' of N_c pairs of coils carrying equal current $I_0/2N_c$, that approximate $J(\ell)$. Their locations satisfy

$$I_{OH}(\ell_i') = (i - \frac{1}{2}) I_0/2N_c \quad (i = 1, \dots, N_c),$$

$$I_{OH}(\ell) = \int_0^\ell d\ell \cdot J(\ell) = \frac{1}{\pi} \sum_{n=0}^K \frac{a_n}{n} \sin(n\pi\ell/L). \quad (11)$$

In the case of J_{VF} , whose total current is zero, N_c must be even. We then have

$$I_{VF}(\ell_i') = (i - \frac{1}{2}) \max |I_{VF}(\ell)| / N_c \quad (i = 1, \dots, N_c/2), \quad (12)$$

and there are two solutions of ℓ_i' for each i .

This approach is flexible in the value of N_c used; the i^{th} coil so located replaces a segment of $J(\ell)$ for $\ell_1 < \ell < \ell_2$. If ℓ_i' is an undesirable location, $2N_c$ can be used in place of N_c in Eqs. (11) and (12). Now two half-sized coils at $(\ell_i')_1$ and $(\ell_i')_2$ replace the original ℓ_i -coil with $\ell_1 < (\ell_i')_1 < (\ell_i')_2 < \ell_2$. A few such iterations will usually yield a satisfactory coil configuration.

4.3 Coil Locations and Circuits .

The approach just outlined has been employed in determining the poloidal field coil locations and currents. In an analogy to the model of Fig. 4.4, they involve the OH (primary), the S-VF (shielding-vertical field), and the T-VF (primary-vertical field) coils. As shown in Fig. 4.6a, the OH and T-VF coils are located on contours that conform to the TF coil shape. A set of decoupling coils in the vicinity of the OH coil is also included to be connected in series with the S-VF coils, and to nearly eliminate the reactive power supply to S-VF coils. The functions of the decoupling coils will be discussed in Sect. 4.5. A poloidal system with only the OH and VF coils external to the TF coils is shown in Fig. 4.6b for comparison purposes. In the actual EPR design, any multiple of the specified number of coils may be used to obtain optimal current, voltage, and power supply conditions, so long as the flux pattern is maintained.

4.3.1 OH Coils

The OH coils are composed of a central solenoid with a length of 6.4 m and 192 turns plus 8 pairs of bunched coils. (A later version replaces this solenoid by a 4.6 m solenoid plus some bunched coils, without reducing the total number of turns.) The OH coils have a total of 320 turns and carry a peak of 61-MA turns to provide a peak-to-peak flux swing of 130 Wb. The maximum air core field is 7.2 T. The locations and the currents of these coils are shown in Table 4.1.

4.3.2 S-VF Coils

The S-VF coils are imbedded on the nuclear shield surface toward the blanket, resulting in a space limitation. For this reason, we have considered coils that are connected in series and carry current in one direction. The current and the optimized locations of an 8-pair system are listed in Table 4.2.

Fig. 4.6. Poloidal field system (a) with magnetic shielding and (b) without.

Table 4.1. Locations and currents of the OH coils (0.19 MA/turn)

<u>No.</u>	<u>R(m)</u>	<u>Z(m)</u>	<u>Turns</u>
1	1.5 - 1.9	0 - 3.2	96
2	2.16	3.48	16
3	2.44	4.26	16
4	3.02	5.26	16
5	4.18	6.18	8
6	5.77	6.43	4
7	7.89	5.80	2
8	10.05	4.25	1
9	11.60	1.99	1

Table 4.2. Locations and currents of the S-VF coils (0.45 MA/turn)

<u>No.</u>	<u>R(m)</u>	<u>Z(m)</u>	<u>Turns</u>
1	4.24	1.90	1
2	5.44	2.87	1
3	6.54	3.14	1
4	7.54	3.05	1
5	8.42	2.67	1
6	9.12	2.07	1
7	9.62	1.31	1
8	9.87	0.45	1

They provide a vertical field of 0.34 T at the plasma center, the field having a decay index $n = - (R/B_z) \cdot (dB_z/dR)$ between 0.15 and 0.25 in the plasma region. To provide a field higher than this requires that some of the S-VF coils carry reversed current (see Fig. 4.5). So, these coils constitute a field-shielding system that produces the maximum vertical field with the least amount of conductor.

According to Eq. (1), we find that 0.34 T is sufficient to hold a plasma with $R_p = 6.75$ m, $a = 2.25$ m, $I_p = 7.2$ MA, $\alpha_i = 1$, and $\beta_p = 1$. However, the value of β_p that can be held by this field is accurately determined by calculating the free-boundary MHD equilibrium, as indicated in Sect. 4.1. Detailed results of this will be discussed in Sect. 4.4.

4.3.3 T-VF Coils

Since β_p is expected to increase beyond 1 by auxiliary heating mechanisms or thermonuclear ignition, it is necessary to introduce some trimming-VF coils external to the TF coils. The optimized locations and current are listed in Table 4.3, producing an additional vertical field of 0.1 T (with $0.0 \leq n \leq 0.55$ in plasma). According to Eq. (1), the total field is then sufficient for a plasma with $\beta_p = 2$. This additional vertical field intersects the TF coils. Since the plasma temperature is expected to build up slowly from $\beta_p = 1$ to 2, only a small addition to the eddy current heat load on the TF coils is expected. The magnitude of this added pulsed field will be discussed in the next section.

Table 4.3. Locations and peak currents of the T-VF coils (0.21 MA/turn)

<u>No.</u>	<u>R(m)</u>	<u>Z(m)</u>	<u>Turns</u>
1	1.70	3.50	- 2
2	2.30	3.90	- 2
3	4.50	6.30	1
4	6.65	6.28	1
5	9.54	4.72	1
6	11.69	1.75	1

4.3.4 Decoupling Coils

Eight pairs of decoupling coils are introduced to reduce the power supply requirements for the S-VF coils. This can be done by simply connecting them in series with the S-VF coils to cancel the voltage induced in the S-VF coils by the OH flux swing. Since the S-VF coils are more closely coupled with the plasma than the decoupling coils, the S-VF coils are thus driven primarily by the plasma induction. As a result, the resistive loss in the S-VF coils becomes the principal power requirement. Considering the field shielding and flux swing effects, it is obvious that the decoupling coils should be arranged in a fashion similar to that of the OH coils. Their locations and currents are given in Table 4.4, and their utilities will be discussed in detail in Sect. 4.6.

Table 4.4. Locations and peak currents of the decoupling coils (0.45 MA/turn)

<u>No.</u>	<u>R(m)</u>	<u>Z(m)</u>	<u>Turns</u>
1	1.93 - 2.13	0 - 3.20	5
2	2.63	3.90	1
3	3.12	4.93	1
4	5.14	6.26	1

4.4 Poloidal Field Shielding

The poloidal fields that intersect the TF coils can be evaluated, using the plasma current profile obtained during the discharge. The plasma current profile in turn depends on the free-boundary MHD equilibria consistent with the currents in various sets of coils.

4.4.1 Time Behavior of Coil Currents

Typically assumed time behavior of the coil currents is shown in Fig. 4.7, together with the plasma current and β_j . Where β_p approximately equals β_j , β_j is a parameter used in the right hand side of Eq. (12a). Constant slopes of J_{OH} in time are assumed during the linear rise start-up and the constant phase of I_p . The currents I_{S-VF} and I_D (in decoupling coils) are assumed to track with I_p , and β_j is assumed to rise quickly during start-up and slowly afterwards. The actual behavior will depend upon supplementary heating and plasma energy loss rates.

Note that the S-VF coils produce a vertical field sufficient to hold a plasma with $\beta_p \geq 1$. To make the plasma fill the chamber during the entire discharge requires the use of I_{T-VF} to adjust the net vacuum field as β_p changes. Figure 4.7b shows the required behavior of I_{T-VF} as a result of the

ORNL-DWG 76-7593

Fig. 4.7. The assumed behavior of the amp-turns in the OH, S-VF, D, T-VF, and VF coils shown in Fig. 4.2, together with the plasma current I_P and β_J . Five "snapshots" (I, II, III, IV, V) are calculated.

free-boundary equilibrium calculations to be shown here. This particular discharge model is a convenient choice for the assumed S-VF coil arrangement. Other models and S-VF coil arrangements will be discussed in Sect. 4.6.

The efficiency of the shielding scheme may be characterized by the shielding ratio, which is the ratio of the averaged poloidal field over the TF coils in similar tokamaks with and without the shielding coils. Here, it is expedient to use the instantaneous field, rather than the time rate of the field, as an early design study. Two otherwise identical poloidal configurations are shown in Fig. 4.6. Without the S-VF and the decoupling coils, the T-VF coils are called the VF coils. The sum of their current magnitude, I_{VF} , that centers the plasma is also given in Fig. 4.7b.

4.4.2 Poloidal Fields

The MHD equilibria which characterize the plasma use the rounded current and pressure profiles shown in Fig. 4.8. This may be a poor assumption for the plasma during the early stage of start-up. However, it is likely to be acceptable with respect to the poloidal fields at the TF coils.

Five "snapshots" of plasma equilibria and the accompanying poloidal fields during the discharge of Fig. 4.7 are calculated.

The results of the poloidal fields in the EPR tokamak and the conventional tokamak of Fig. 4.6 are shown in terms of the flux plots in Figs. 4.9a-i. Note that in each of the flux plots, the field strength is proportional to the density of the flux lines divided by the major radius R . It is seen from Fig. 4.4a that the OH coils channel the OH flux away from the TF coils. When the plasma current rises, Fig. 4.9 shows that the S-VF coils effectively shield the plasma flux from the TF coils, while the conventional coil system allows poloidal fields comparable to that at the plasma edge to reach the TF coils.

The poloidal field over the TF coils is also calculated. The averaged field strength is plotted in Fig. 4.10. It shows a net shielding ratio (ratio of average field without shielding to that with shielding in similar tokamaks) of between 2.0 and 6.1.

The highest shielding ratio, 6.1, is obtained at $\beta_j = 1.33$ when $I_{T-VF} = 0$ and I_{S-VF} supplies the total equilibrium field. In other words, I_{S-VF} and the S-VF coils closely represent the eddy current on a conducting shell encircling the given plasma. When $\beta_j \neq 1.33$, the S-VF coil locations are no longer optimal for the best field shielding effect even in a well-centered plasma. This is evident for the peak field, which in the shielded case is located at $(R, Z) = (3.08 \text{ m}, 0)$; and equals 0.53 T at $\beta_j = 0.4$, and 0.28 T at $\beta_j = 0.8$; and equals roughly 3 times the average instantaneous field.

These results suggest that the S-VF coils be relocated to provide the vertical field required for equilibrium with a lower β_j , typically only half the value of β_{j0} reached at the end of the current rise. In view of present-day experiments, β_{j0} should be between 0.5 ~ 1.0. According to Fig. 4.5, the S-VF coils will have a more nearly uniform distribution than that shown in Fig. 4.6a. It should be noted that the field at the TF coils due to I_{T-VF} averages as ~ 0.05 T, and is not important in the cases with shielding coils. This field is comparable to the averaged stray field (0.03 T) from the OH coils. However, I_{OH} produces a peak stray field of 0.15 T at the outer edge of the TF coils near the OH solenoid. This field is exceptionally large as compared with its average, 0.03 T. Since this maximum field contributes directly to the maximum pulsed field experienced by the TF coils, future designs of the OH coils must aim at reducing this field by a factor of roughly 5. Possible improvements include reducing the solenoid length and using more and smaller OH coil bundles near the solenoid.

The locations of the S-VF coils for a circular plasma are rigid in that they can be optimal for shielding and vertical field for only one value of β_p . As a result, the field shielding ratio of the

ORNL-DWG 76-10849

Fig. 4.8. The current and pressure profiles used in the equilibria at $\beta_J = .4$ and 1.3.

Fig. 4.9. "Snapshots" of field due to currents and corresponding to those of Fig. 4.3. Figs. 4.9(a), (b), (d), (f), and (h) are for the case with shielding coils, and Figs. 4.9(a), (c), (e), (g), and (i) are for the case without shielding coils.

Fig. 4.9 (d), (e).

Fig. 4.9 (h), (i).

ORNL-DWG 76-7589

Fig. 4.10. Average poloidal field over the TF coils as functions of t in tokamaks with and without shielding coils for the assumed current behavior shown in Fig. 4.3. Improvement for the case with shielding from (—) to (-----) is possible.

case considered is only 2.0 at the beginning of the start-up, although later it can be better than 6.0 when the "design" β_j is reached. This may be understood in terms of Fig. 4.5. It shows that as the vertical field increases to provide equilibrium at higher β_p , the distribution of J_S goes through dramatic changes.

There are two types of engineering remedies for this difficulty. One is individual control of the current in each S-VF coil which leads to a large number of connections. The other, as suggested by Fig. 4.5, is the replacement of J_S by J_{OH} and J_{VF} ; J_{OH} has a sum equal and opposite to the plasma current and produces zero field in the plasma region; J_{VF} has a sum of zero and produces the equilibrium vertical field. Then two independent sets of coils can be used to replace J_{OH} and J_{VF} separately. This arrangement will eliminate the need for T-VF coils. A shielding ratio of above 6.0 throughout the discharge is likely to be obtained (see Fig. 4.10). This is equivalent to a reduction of the eddy-current heating rate of the TF coils by a factor of over 36. However, more space and more amp turns would be required for this arrangement.

If these remedies are not used, there are still alternatives to the type of current rises shown in Fig. 4.7. These include:

1. Set $I_{T-VF} = 0$ until the completion of the current rise, but maintain $I_{S-VF} = -I_p$. This will push the plasma against the inside limiter and hence lower the shielding ratio further than that shown in Fig. 4.10.
2. Set $I_{T-VF} = 0$ and maintain $I_{S-VF} < I_p$ to center the plasma, or even let the plasma lean against the outside limiter. It would be interesting to obtain the shielding ratio for this case.

4.5 Volt Seconds and Power Supply

In a design that successfully eliminates field over a sizable volume, substantial savings in the stored field energy are realized. This feature is made obvious by Fig. 4.9a-i and by the field calculations in Sect. 4.4. A reduction in the field energy means reductions in volt-second and power supply requirements. A comparison of these requirements in poloidal field systems with and without shielding (Figs. 4.6a and b) follows.

With the decoupling coils (D) and the S-VF coils (S) connected in series ($I_D = -I_S$), the circuit equations for the OH coils (O), plasma (P), the S-VF coils, and the T-VF coils (T) can be written as

$$V_O = L_O \dot{I}_O + M_{OP} \dot{I}_P + (M_{OS} - M_{OD}) \dot{I}_S + M_{OT} \dot{I}_T, \quad (13)$$

$$0 = V_P = L_P \dot{I}_P + M_{OP} \dot{I}_O + (M_{PS} - M_{PD}) \dot{I}_S + M_{PT} \dot{I}_T + R_P I_P, \quad (14)$$

$$V_S = (L_S + L_P - 2M_{SD}) \dot{I}_S + (M_{OS} - M_{OD}) \dot{I}_O + (M_{PS} - M_{PD}) \dot{I}_P \\ + (M_{ST} - M_{DT}) \dot{I}_T + R_S I_S, \quad (15)$$

$$V_T = L_T \dot{I}_T + M_{OT} \dot{I}_T - M_{PT} \dot{I}_P + (M_{ST} - M_{DT}) \dot{I}_S. \quad (16)$$

In the above, L_O represents the single-turn averaged self-inductance of the OH coils, and M_{OP} represents the single-turn mutual inductance between the coils and the plasma loop, etc. Equations for the system of Fig. 4.6b can be obtained by omitting terms with subscript S or D.

For the coil systems of Fig. 4.6, the per (turn)² inductances have been calculated and listed in Table 4.5. In this case, the plasma is assumed to have a uniform current density. In an ideal system, one should have

$$L_{OO} = M_{OP} = M_{OS} = M_{OD}, \quad (17)$$

$$M_{PS} \gg M_{PD}, \quad (18)$$

$$M_{OT} = M_{DT} = 0. \quad (19)$$

Table 4.5. Per (turn)² inductances (μH) of the poloidal systems shown in Fig. 4.6
(The plasma is assumed to have a flat current profile.)

	<u>OH</u>	<u>Plasma</u>	<u>S-VF</u>	<u>Decoupling</u>	<u>T-VF</u>
OH	1.03				
Plasma	1.11	12.04			
S-VF	1.12	9.70	12.30		
Decoupling	1.09	1.35	1.30	1.59	
T-VF	0.11	2.25	3.06	0.16	3.02

Equation (17) implies no OH leakage flux with respect to the plasma loop, and complete decoupling of the S-VF coils from the OH coils. Equation (18) indicates that the S-VF coils are driven primarily by plasma induction. Equation (19) specifies a completely decoupled T-VF (or VF) coils system. It is seen in Table 4.5 that Eq. (17) is satisfied to within 9%; M_{PS} is more than 7 times M_{PD} ; and M_{OT} and M_{DT} are around 10% of L_{OP} . Improvements of these conditions are possible by refinement of the computations in Sect. 4.3.

In the case of the shielding coil, $I_S = -I_P$, $I_T \approx 0$, and by the conditions of Eqs. (17) - (19), Eqs. (13) - (15) can be reduced to

$$-V_0 = \dot{I}_P(L_P - L_0 - M_{PS} + M_{PP}) + R_P I_P. \quad (20)$$

$$V_S = \dot{I}_S(L_{ST} + L_D - 2M_{SD} + M_{PS} + M_{PD}) + R_S I_S. \quad (21)$$

It is now required to model the plasma resistivity term during the start-up phase. Approximate calculations by Bodin et al.¹¹ indicate that an amount of OH volt-seconds equivalent to $2 L_{Pj} I_P$ is expended during start-up. Since the plasma internal inductivity, L_j (emu), is of the order of 1 in present-day experiments (corresponding to parabolic current profiles), this translates to a loss of roughly 61 Wb for start-up. This means a time-averaged resistive plasma voltage drop of 30 V for a 2-sec current rise. Substituting the values of inductances from Table 4.5, we find that the amount of inductive volt-seconds required to set up the plasma current is less than 13 Wb for a flat current profile. So the flux swing necessary for start-up is 74 Wb.

The power supply requirements can be estimated as follows. Assuming that 30 Wb are needed after the start-up, the total flux swing amounts to 104 Wb. With the mean air core solenoid radius set at 1.7 m, a field evaluation at the air core shows that the peak OH current is 59-MA turns (this value is roughly 80% of the design value). The peak power supply to the OH coils, with $V_0 = 36$ V, is then 2.1 GW. The power supply to the S-VF coils can be estimated by Eq. (21) to be 76 MW reactive (at the end of current rise). After the current rise, the resistive losses will depend upon conductor material and operation temperature.

Similar calculations for the poloidal system without shielding begin with the reduced equations:

$$-V_0 = (L_P - L_0) \dot{I}_P + M_{PT} \dot{I}_T + R_P I_P, \quad (22)$$

$$V_T = -M_{PT} \dot{I}_P + L_T \dot{I}_T, \quad (23)$$

where we have set $I_S = 0$ in Eqs. (13) - (16) and used the conditions of Eqs. (17) - (19). For this case, the time dependence of I_T has been shown in Fig. 4.7. The time-averaged value for $\dot{I}_T$ during

start-up is 11.8 MA/sec. The inductive flux swing needed for start-up from the OH coils is then 26 Wb. The total flux swing needed from the OH coils for start-up is then 87 Wb, and that for the entire discharge is 117 Wb. This corresponds to a peak of 66 MA turns for the OH coils with a peak field of 6.4 T.

In this case the T-VF coils supply a large portion of the flux swing needed to drive the plasma current. This can be seen by estimating the power supply to the T-VF coils. At the end of the current rise, the reactive power according to Eq. (23) is 0.65 GW. The power supply to the OH coils at the beginning of the start-up is 2.9 GW with $V_0 = 44$ V.

The comparisons between the poloidal field systems, with and without shielding, with equal I_p and B_T , are summarized in Table 4.6. It is seen that the magnitude of power supply and stored energy can be reduced to 59% and 66% respectively by using the shielding system.

Table 4.6. Comparisons between magnetic shielding and conventional poloidal field systems

Magnetic Shielding (ORNL EPR)		Conventional (VFC Ext. to TFC)			
Pulse Field at TF coils (time average for start-up)					
Average $\dot{B}$.04 T/sec		0.21 T/sec			
Plasma Induction					
(1) I_P strongly coupled to I_S $M_{PS} - M_{PD} = 8.4 \mu H$		(1) I_P weakly coupled to I_T $M_{PT} = 2.3 \mu H$			
(2) I_S decoupled from I_0		(2) L_T decoupled from I_0			
OH V-S needed for start-up					
Inductive	13		26		
Start-up loss	61		61		
Steady-state loss	<u>30</u>		<u>30</u>		
	104		117		
Peak power supply and stored energy					
OH	2.1 GW	1.8 GJ	OH	2.9 GW	2.2 GJ
S-VF	76 MW (reactive)		VF	650 MW	
	50 MW (resistive)				
	$\int V_S dI_S = 76$ MJ (reactive)			$\int V_T dI_T = 650$ MJ	

4.6 Other Considerations

The advantages of the magnetic shielding system for an EPR include reduction of the pulsed field at the TF coils to roughly one-fifth and the power supply requirements to 59% of a similar tokamak without shielding. It is important to reduce the power requirements of the VF coils, which require a rapid response time to maintain plasma equilibrium. It is probably more important to reduce the peak air core field and stored energy, in order to minimize the technological difficulties involved

in pulsed high-field superconducting coils. It is seen that the magnetic shielding system is eminently more suitable for these purposes than conventional systems. In addition, in designs of large superconducting TF coils, it is reasonable to have comparable heating loads from conduction and eddy current. This implies that a substantial reduction of the pulse fields can have a far-reaching influence on the design and manufacture of TF coils. Considerable savings in cost, cooling power supplies, and structural strength can then be realized for the superconducting coils.

However, these advantages will have to be balanced by the engineering complications introduced by placing coils within the TF coils. Nevertheless, it is well known that tokamaks with noncircular cross sections,^{12,13} poloidal field divertors,¹⁴ magnetic multipole start-up,¹⁴ and adiabatic compression¹⁵ all require internal coils to be practical.

4.6.1 Divertor Configuration

A divertor configuration for EPR can be obtained by imposing a strong quadrupole field above and below the plasma column. S-VF coils number 2, 3, and 4 are used for this purpose. The necessary divertor current superimposed on coil number 3 is found to be $0.3 I_p$, so this coil now carries a current of 1.7 MA in the same direction as the plasma current. A current of $-0.15 I_p$ is also superimposed on coils number 2 and 4 to result in $-1.5 I_p$. A typical free-boundary equilibrium obtained with these currents is shown in Fig. 4.11, where the plasma pressure and current are confined within the separatrix location. It is seen that if the added total divertor current of 8.4 MA and the resulting nuclear shield requirements near the divertor zone can be provided, this divertor configuration could become an interesting addition to the EPR design.

4.6.2 Noncircular Cross Section

There are substantial advantages to be gained by using noncircular cross sections because of the associated MHD properties.¹⁶ An investigation in terms of the stability requirements of the ideal local modes has shown that the highest β can be obtained in a D-shaped plasma with an elongation of 1.6, a modest triangularity, and $\beta_p \approx 2.4$.⁹ An immediate question to be answered is whether such a D-shaped plasma can be readily maintained at high β_p values with peaked plasma current profiles.

A typical result of a free-boundary D-shaped plasma is shown in Fig. 4.12. In this calculation, 8 pairs of S-VF coils are placed at a distance of 56 cm from the plasma boundary to produce a vacuum field with negative decay index (Fig. 4.12a). The resulting plasma equilibrium at $\beta_p \approx 2.4$ with peaked current profile is seen to retain an elongation of 1.66 (Fig. 4.12b). It is well known that a circular plasma becomes unstable in its vertical position when the decay index is negative.⁶ However, the stability of a D-shaped plasma against axisymmetric motions will require future examination in detail.¹⁷

Acknowledgment

Computations of poloidal field coil locations have been obtained with the aid of J. R. Moore.

REFERENCES

1. M. Roberts and E. S. Bettis (eds.), *Oak Ridge Tokamak Experimental Power Reactor Study Reference Design*, ORNL/TM-5042, Oak Ridge (November 1975).
2. V. D. Shafranov, *Review of Plasma Physics*, Vol. 2, p. 103, Consultants Bureau, New York, 1966.
3. J. M. Greene, J. L. Johnson, and K. Weiner, "Tokamak Equilibrium," *Phys. Fluids* 14: 671 (1971).

ORNL-DWG 76-6828

1	0.1735606E	02
2	0.1486427E	02
3	0.1237248E	02
4	0.9880686E	01
5	0.7388893E	01
6	0.4897100E	01
7	0.2405308E	01
8	0.0	
9	0.1735606E	00
10	0.3471212E	00
11	0.8647156E	-01
12	0.2578278E	01
13	0.5070068E	01
14	0.7561859E	01
15	0.1005365E	02
16	0.1254544E	02
17	0.1503723E	02
18	0.1752902E	02
19	0.2002081E	02
20	0.2251262E	02
21	0.2500441E	02

Fig. 4.11. Divertor, magnetic limiter configuration. The total divertor amp-turns are equal to 1.2 times the plasma current.

ORNL-DWG 76-10844

Fig. 4.12. (a) Typical free-boundary, D-shaped plasma flux plot. (b) The vacuum equilibrium field.

4. J. D. Callen and R. A. Dory, "Magnetohydrodynamic Equilibria in Sharply Curved Axisymmetric Devices," *Phys. Fluids* 15: 1523 (1972).
5. *Thermonuclear Division Annual Progress Report for Period Ending December 31, 1974*, ORNL-5053, Oak Ridge (July 1975).
6. S. Yoskikawa, "Application of the Virial Theorem to Equilibria of Toroidal Plasmas," *Phys. Fluids* 7: 278 (1964).
7. C. Mercier, "Critère de Stabilité d'un Système Toroidal Hydromagnétique en Pression Scalaire," *Nucl. Fusion Suppl. Pt. 2*, 801 (1962).
8. J. M. Greene and J. L. Johnson, "Stability Criterion for Arbitrary Hydromagnetic Equilibria," *Phys. Fluids* 5: 510 (1962).
9. Y-K. M. Peng et al., *MHD Equilibria and Local Stability in Axisymmetric Tokamaks*, ORNL/TM-5267, Oak Ridge (to be published).
10. K. Toi and T. Takeda, *Numerical Method to Obtain the Optimum Configuration of External Magnetic Field Coils in a Tokamak Device*, JAERI-M-6018, Japan Atomic Energy Research Institute (February 1975); W. M. Stacey et al., *Tokamak EPR Studies*, ANL/CTR-75-2, Argonne, Illinois (June 1975).
11. H.A.B. Bodin, "Start-Up," Fusion Reactor Design Problems, paper presented at IAEA Culham Workshop, Abingdon, Oxfordshire, England, January 29 - February 15, 1974.
12. T. Ohkawa, "Multipole Configurations with Plasma Current," *Bull. Am. Phys. Soc.* 13(11): 1527 (1968).
13. JET (Joint European Torus), *Culham Laboratory 12th Annual Report*, UKAEA Research Group, Abingdon, Oxfordshire, England, 1974.
14. K. E. Weimer et al., "Studies of MHD Equilibria in a Divertor Tokamak," *Bull. Am. Phys. Soc.* 20: 1279 (1975).
15. Christensen et al., "A Machine for Adiabatic Toroidal Compression," *IEEE Trans. on Nucl. Sci.* NS-18 (4): 49 (1971).
16. H. P. Furth, "Tokamak Research," *Nucl. Fusion* 15: 487 (1972).
17. E. Rebhan, "Stability Boundaries of Tokamaks with Respect to Rigid Displacements," *Nucl. Fusion* 15: 277 (1975).

5. TOROIDAL FIELD RIPPLE

5.1 Introduction

In the design of tokamak devices, the ripple produced by the finite number of toroidal field coils destroys the ideal axisymmetry of the configuration and is responsible for additional particle trapping, loss regions, and plasma transport. In the region of low collision frequencies, the magnetic field ripple might be the main factor determining the transport of particles and energy.¹⁻⁴ In the background plasma, collisions are usually strong enough to scatter the particles out of the ripple-trapping region in velocity space before they drift out of the plasma. Thus, this effect leads only to an enhancement of the transport coefficients.¹⁻⁴ For particles trapped in the toroidal field, the average location of their banana orbits can experience a new radial drift due to the presence of ripples, which results in an additional diffusion term.^{2,5}

At the higher ion energies typical of neutral beam particles injected into large devices (≥ 100 keV), or alpha particles produced from DT fusion (3.5 MeV), collisional effects are much weaker and the fast ions trapped in the ripples can drift out of the plasma before collisional detrapping occurs.

In this section we investigate the effects of toroidal field ripple on: (1) the plasma transport coefficients, (2) the loss of fast ions trapped in the toroidal field ripples, and (3) the relaxation of toroidal flows that might be induced by any unidirectional component of neutral beam injection⁶ or alpha particle losses.⁷

Considering a simple model of tokamak confinement geometry, the resulting magnetic field can be written as¹⁻⁴

$$\underline{B} = \frac{B_0}{h} \left(\frac{B_r}{B_\phi}, \theta, 1 - \delta(r, \theta) \cos N\phi \right), \quad (1)$$

where r and θ are polar coordinates in the minor cross section of the torus, and ϕ is the angular coordinate along the magnetic axis; $h = 1 + (r/R_0) \cos \theta = 1 + \epsilon \cos \theta$; $\delta(r, \theta) = (B_{\max} - B_{\min}) / (B_{\max} + B_{\min}) \ll r/R_0$ is the depth of the ripples; N is the number of toroidal field coils; and $\theta = \epsilon/q$ where q is the safety factor. The radial component of $\underline{B}$ necessary to satisfy $\nabla \cdot \underline{B} = 0$ can be found from

$$\frac{\partial}{\partial r} (rhB_r) + \frac{r}{R_0} \frac{\partial}{\partial \phi} (B_\phi) = 0, \quad (2)$$

which is on the order of $B_r \sim \delta B_0 \sin N\phi$. Hence, the field strength is approximately

$$B \approx B_\phi = B_0 [1 - \epsilon \cos \theta - \delta(r, \theta) \cos N\phi]. \quad (3)$$

The actual relationship of δ to r and θ depends substantially on the coil design and on the size and shape of the gap between the coils. In a typical case, to find the function $\delta(r, \theta)$, we must carry out computer calculations for the particular magnet design. However, a great number of calculations have shown that it is possible to write an approximate expression for the ripple $\delta(r, \theta)$ which is suitable for the purpose of our calculations, as²

$$\begin{aligned} \delta(r, \theta) &= \delta(r) g(\theta) \\ &= \delta_a \left[\frac{\delta_0}{\delta_a} + (r/a)^n \right] \exp(-\beta\theta^2) \end{aligned} \quad (4)$$

where a is the plasma radius; δ_0 and δ_a are the values of δ at $r = 0$ and $r = a$, respectively, with $\delta_0/\delta_a \sim 10^{-3} - 10^{-4}$; and β is a parameter which can be determined from the calculations and closely depends on the shape of the coils and plasma locations, usually, $0.5 \leq \beta \leq 1$. In the ORNL EPR design, $\delta_a = 2.2\%$, $\delta_0 = 0.012\%$, $n = 4$ and $\beta = (2.8/\pi)^2$.

The first point to consider is the threshold energy above which particles execute more than one bounce motion in a ripple before being scattered out of the loss region ($v < \delta^{1/2} v_{\perp}$) associated with the ripples. This energy for deuterium and tritium ions is given by ¹⁻²

$$E_{th} \approx \delta^{-3/4} (R_0/N)^{1/2} [8 \times 10^{-13} n_e Z_{eff}]^{1/2} \text{ (eV)} . \quad (5)$$

The threshold energy is a factor of $(m_i/m_e)^{1/4}$ larger for electrons. For the present (future) tokamaks, the ion threshold energy is typically on the order of .3 keV (3 keV) at the plasma edge and a factor of about 10^2 (50) larger at the plasma center (see Table 5.1).

Table 5.1. Typical parameters

	ORMAK Upgrade	ORNL EPR
Major radius, cm	90	675
Minor radius, cm	30	225
Number of coils	24	20
δ ($r = a$)	4.2%	2.2%
δ ($r = 0$)	0.009%	0.012%
n_0 , cm ⁻³	7×10^{13}	8×10^{13}
T_i , keV	2 - 3	9 - 10
B_{T0} , kG	45	48
$q(a)$	3 - 4	2.5
Z_{eff}	4	2
$E_{th}(r = a)$, keV	0.3	3
$E_{th}(r = 0)$, keV	30	150
$E_*(r = a)$, keV	10	48
$E_*(r = 0)$, keV	100	380

The second point to consider is that when a particle has been trapped in a ripple, its guiding center drifts more or less along a contour of constant B . Since such contours are not closed within the torus for weak ripples, the particles are not confined unless they are scattered out of the ripple loss region before they can reach the wall. This consideration defines a critical energy, above which a particle will reach the wall if it is trapped in the ripple, given by ¹⁻²

$$E_* \approx [10^{-14} \frac{BR_0 a}{\delta} n_e Z_{eff}]^{2/5} \text{ (eV)} . \quad (6)$$

The critical energy is a factor of $(m_i/m_e)^{1/5}$ larger for electrons. Again considering typical present (future) tokamak base cases, we can find the critical ion energy to be on the order of 10 keV (50 keV) at the plasma edge and a factor of about 10 (10) greater near the plasma center (see Table 5.1).

With these critical energies for relevant ripple-related phenomena in mind, we now proceed to discuss the effects of toroidal field ripple on the plasma.

5.2 Transport Modifications

The dominant transport modifications come from the enhanced ion heat conduction of the background plasma resulting from the ripple-trapping of ions with energies exceeding the threshold energy and from ripple-induced banana drift. (The effect on electron heat conduction is much smaller because the threshold energy is much higher.) From the estimates of E_{th} and E_* given in Sect. 5.1, we can see plasma ion temperatures anticipated in present and future tokamaks to be in the range of $E_{th} < T_i < E_*$, the relevant ion thermal conductivity coefficient χ_i is in the middle range of collisionality.^{1,2}

The expression for the ion heat conductivity coefficient in the above-mentioned regime for ripple-trapping is²⁻⁴

$$\chi_i^{\delta} = 46.5 \frac{\delta(r)^{3/2} G(\alpha, \beta)}{v_{ei}} (m_i/m_e)^{1/2} v_{drift}^2, \quad (7)$$

and for banana drift is²

$$\chi_i^{BD} = 9.23 \epsilon^{3/2} \delta^2(r) \frac{N^2}{(n+2)^2} \frac{v_{drift}^2}{v_{ei}} (m_i/m_e)^{1/2} q^2 Q^2, \quad (8)$$

where $v_{ei} = 1.6 \times 10^4 \langle Z \rangle (n_0/10^{13})/T^{3/2}(\text{keV})$, $q = \epsilon B_T/B_\theta$, $v_{drift} = cE/eBR$, $Q = d\ln q/d\ln r$, and $G(\alpha, \beta)$ is a function² which depends on the magnetic configuration. The computed value of $G(\alpha, \beta)$ for different β values is plotted against $\alpha(r) = \epsilon/Nq\delta(r)$ in Fig. 5.1.²

The ion heat conduction coefficients for these two mechanisms, ripple-trapped (δ) and banana-drift (BD), have been calculated for the parameters typical of the ORNL EPR. The coefficients have then been compared with the total heat conduction loss due to the neoclassical (NC), trapped-particle modes [trapped-ion modes (TIM) and trapped-electron mode (TEM) 1 or 2], and pseudoclassical (PS) coefficients currently used in the plasma transport simulation codes. The expressions used in the calculations are given below for $T_e = T_i$.⁸

$$\chi_i^{NC} = \epsilon^{3/2} q^2 \rho_e^2 v_{ei} (m_i/m_e)^{1/2}$$

$$\chi_e^{TIM} = \chi_i^{TIM} = \frac{\epsilon^{5/2}}{4} \frac{v_{drift}^2}{v_{ei}} (R/r_n)^2$$

$$\chi_e^{TEM1} = 3 \epsilon^{3/2} \frac{v_{drift}^2}{v_{ei}} (R^2/r_n r_T)$$

$$\chi_e^{TEM2} = C_2 (q^2/\epsilon^2) \rho_e^2 v_{ei}$$

$$\chi_e^{PS} = C_0 (q^2/\epsilon^2) \rho_e^2 v_{ei}$$

$$\chi_i = \chi_i^{NC} + \chi_i^{TIM}, \quad \chi_e = \chi_e^{TIM} + \chi_e^{PS} + \min \chi_e^{TIM1 \text{ or } 2}$$

Fig. 5.1. $G(\alpha, \beta)$ versus α .

where $r_n = [1/n \, dn/dr]^{-1}$, $r_T = [1/T \, dT/dr]^{-1}$, $C_0 \sim 4 - 10$, $C_2 = 0.06 \, \epsilon^{1/2} (d \ln T / d \ln n) (m_i/m_e) (B_\theta / \theta B_T)$, $\theta = r_n / L_s$, $L_s^{-1} = (r/R_0 q^2) (\partial q / \partial r)$, and ρ_e is the electron Larmor radius in the toroidal field.

In the calculations, a parabolic temperature profile, $T = T_0 (1 - r^2/a^2)$, has been assumed with $T \sim n^2$, $n = n_0 (1 - r^2/a^2)^{1/2}$. The safety factor q is calculated assuming a current density profile that varies like $T^{3/2}$. The results are shown in Fig. 5.2. Although the heat conduction coefficients associated with the ripple-trapping and banana drift are comparable to those calculated from the ion neoclassical and electron pseudoclassical coefficients, the trapped-ion and trapped-electron mode coefficients are about 20 times larger at their peaks, if poloidal angle effects are ignored ($\beta = 0$). The inclusion of poloidal angle effects reduces the losses associated with ripple trapping by about a factor of 10 to 15, but does not significantly change the banana-drift losses. The overall enhanced ion heat loss due to ripples is found to have a negligible effect on plasma confinement for the ORNL EPR design. However, if one assumes better confinement — for example, if ignition is to be achieved at a higher toroidal field (say 70-85 kG) — energy losses must be no larger than one-tenth of the trapped particle losses. Even under this condition, the enhanced heat losses due to the 2.2% field ripple at the edge are less than 10% of those due to trapped-particle instabilities.

5.3 Fast Ion Losses

It is well-known that deviations from axisymmetry in toroidal devices lead to the formation of loss regions because particles trapped in the ripple mirrors tend to be localized and are therefore subject to an uncompensated ∇B drift motion. A high-energy particle trapped in a ripple well drifts vertically to the wall in a time much shorter than that taken to slow down or to scatter in pitch angle and become untrapped. Consequently, the trapped high-energy particles in the bumpy field will be lost directly from the plasma without contributing to plasma heating and will be deposited on discrete surface regions of the vacuum wall. The loss of fast ions produced by neutral beam injection due to ripple-trapping is similar to the finite banana width effects that have been found to be important in present neutral injection experiments.⁶ However, since δ varies as (r/a) ,⁴ the critical energy for direct loss is greater than the injection energy of 100 (or 200) keV for $r/a \lesssim 0.65$ (0.4). In the parallel injection case, the fast ions must scatter from their essentially parallel injection velocity to a nearly perpendicular velocity to be lost by this process. Since the fast ions lose some fraction of their energy (~ 10 -50% depending on Z_{eff}) during this scattering process, the critical energy must be somewhat lower than the injection energy for significant losses; hence, the region of plasma free from such losses is even greater than the $r/a \lesssim 0.65$ (0.4) estimate given above. The fraction of beam particles (with parallel injection) diffusing into the ripple loss region during this slowing-down process will be neglected.

For an isotropic source of alpha particles, the fraction of particles, $f_{\alpha 1} \approx (2\delta)^{1/2}$, trapped in a ripple well will be lost directly. Those alpha particles that are not lost immediately after their production can diffuse to the ripple loss region during the slowing-down process by pitch-angle scattering. To evaluate the latter process we need to determine the α -particle distribution function during the slowing-down process. The α -particle distribution function f satisfies⁶

$$\tau_s \frac{\partial f}{\partial t} = \frac{1}{v^2} \frac{\partial}{\partial v} (v^3 + v_c^3) f + A (v_c^3/v^3) \mathcal{L} f + S \tau_s, \quad (9)$$

where $\tau_s = 0.12 (T_e/1 \text{ keV})^{3/2} (m_\alpha/m_H) / Z_\alpha^2 (n/10^{13} \text{ cm}^{-3})$ is the slowing-down time of α -particles, $A = 1/4 \, m_i \langle Z \rangle / m_\alpha [Z]$, $[Z] = \sum n_j Z_j^2 (m_i/m_j) / \sum n_j Z_j$, $\langle Z \rangle = \sum n_j Z_j^2 / \sum n_j Z_j$, $\mathcal{L}$ is the Lorentz scattering operator, $E_c = m_\alpha v_c^2/2 = 14.8 T_e$, $(m_\alpha/m_H)^{1/3} (m_\alpha [Z] / m_i)^{2/3}$ is the critical energy during the slowing-down process at which energy is being transferred equally to plasma ions and electrons, $S = \dot{n}_\alpha \delta(v - v_0) / 4\pi v_0^2$ is an assumed isotropic source of α -particles, and $\dot{n}_\alpha$ is the rate of α -particles being produced.

Fig. 5.2. Heat conduction parameters versus plasma radius for the ORNL EPR design parameters.

We assume the inverse aspect ratio ϵ is much larger than the ripple size δ , so that there is a class of α -particles that is deeply trapped in the toroidal field but untrapped in the field ripple. Since the classical Coulomb interaction is predominantly small-angle scattering, this class of α -particles is most easily lost to the ripple-trapped region by pitch-angle scattering. For convenience we define the pitch-angle variable

$$\chi^{-2} = 2\mu B_0 / [E - \frac{e\phi}{m\alpha} - B_0(1 - \epsilon \cos \theta + \epsilon^2 \cos^2 \theta + \delta)], \quad (10)$$

where μ and E are the magnetic moment and energy per unit mass of the α -particle. We can perform the bounce average in the ripple field periods by neglecting the slow poloidal angle variation (which is on the order of $2\pi/Nq$) and obtain the bounce average

$$\mathcal{L} \approx \frac{1}{\delta} \frac{\partial^2}{\partial \chi^2}. \quad (11)$$

The slowing-down equation can then be solved easily. We find that the fraction of α -particles diffusing into the ripple loss region by collisions is

$$f_{\alpha 2} \cong \left[\frac{m_i \langle Z \rangle}{6m_\alpha [Z]} \ln \left(\frac{v_*^3 + v_c^3}{v_o^3 + v_c^3} \frac{v_o^3}{v_*^3} \right) \right]^{1/2}, \quad (12)$$

where $1/2 m_\alpha v_*^2 = E_*$. It is clear that most of the α -particles can be lost by this process; however, loss occurs primarily on the order of or below E_c , so it appears that they deposit their energy before they escape from the system. Therefore the effect of their loss on the power balance might not be significant.

5.4 Relaxation of Toroidal Flow by Ripple Viscosity

In general, magnetic field ripple has the beneficial effect of producing a viscous drag on toroidal rotations or flows in the plasma, such as may be introduced by neutral beam injection and/or α -particle losses. Although there are several possible mechanisms by which toroidal momentum can be lost in a torus, most of the damping mechanisms are not very effective in large size plasmas. The non-ambipolar momentum loss due to the rippled toroidal field is always present, and is usually the dominant toroidal flow relaxation process. The equation governing the toroidal momentum is

$$m_i n_i \frac{\partial U}{\partial t} = \frac{\partial P_{\text{BEAM}}}{\partial t} - \frac{m_i n_i}{\tau_\delta} (U - U_0), \quad (13)$$

in which τ_δ is the characteristic ripple-induced decay time, and U_0 is the residual velocity. In the ripple collisionality regime of interest, $U \gg U_0 \approx 0$, and τ_δ is given by^{1,9}

$$(\tau_\delta)^{-1} = 9.28 \delta^{3/2} G(\alpha, \beta) (r/R)^2 \left[\frac{\Theta v_i}{r v_{ii}} \right]^2 v_{ii}, \quad v_{ii} < N v_i \delta^{3/2} / R \quad (14)$$

where $\Theta = B_\theta / B_T$ and $v_i = \sqrt{2T_i/m_i}$ is the ion thermal velocity. The beam source term, $\partial P_{\text{BEAM}} / \partial t$, is given by

$$\frac{\partial P_{\text{BEAM}}}{\partial t} = \dot{n}_f m_f v_o, \quad (15)$$

in which $\dot{n}_f = [I_{\text{BEAM}}/e] H(r)/\text{Volume}$ is the rate of fast-ion density input from the beam and $m_f v_o$ is the fast-ion momentum. Available models were used to calculate the radial shape factor of the fast ions, $H(r)$, deposited by neutral beam injection for 200 and 150 keV deuterons.¹⁰

The above equation may be solved for the steady state rotational velocity ($\partial U/\partial t = 0$) to yield

$$U = \frac{1}{m_i n_i} (\partial P_{\text{BEAM}}/\partial t) \tau^{\delta} \quad (16)$$

The quantity that indicates whether the induced rotation is large enough to modify the plasma equilibrium and stability or not is the toroidal Mach number defined in terms of the ion acoustic sound speed (V_s)

$$M = U/V_s .$$

The values for U and M were calculated for the ORNL EPR base case, assuming unidirectional tangential injections of 40 MW and 100 MW (in the beam-driven mode) of deuteron beams. The results are shown in Fig. 5.3. Here we see that ripple viscosity (for an edge ripple of 2.2%) is the dominant damping mechanism for the EPR. The flow velocity is much smaller than the sound speed, except near the center of the plasma where internal disruptions or other effects may dominate the energy transport. Since the apparent toroidal flow velocity for unidirectional injection can apparently be quite high near the plasma center, it may be very important to balance co- and counter-injection beams so as to cancel the momentum input there as much as possible. For this purpose a higher magnetic field ripple, particularly near the plasma center, would be desirable.

Acknowledgments

These calculations were performed in conjunction with J. D. Callen and K. T. Tsang, with computational aid from J. R. Moore and T. Uckan.

REFERENCES

1. N. A. Uckan, K. T. Tsang, and J. D. Callen, "Toroidal Field Ripple Effects in Large Tokamaks," *Proc. 6th Symposium on Engineering Problems of Fusion Research*, p. 1105 (1976).
2. N. A. Uckan, K. T. Tsang, and J. D. Callen, *Effects of the Poloidal Variation of the Magnetic Field Ripple on Enhanced Heat Transport in Tokamaks*, ORNL/TM-5438, Oak Ridge (to be published).
3. T. E. Stringer, "Effect of the Magnetic Field Ripple on Diffusion in Tokamaks," *Nucl. Fusion* 12: 689 (1972).
4. J. W. Connor and R. J. Hastie, "Neoclassical Diffusion Arising from Magnetic-Field Ripples in Tokamaks," *Nucl. Fusion* 13: 221 (1973).
5. J. N. Davidson, "The Effect of Toroidal Field Ripple on Trapped Particle Transport in Tokamaks," *Bull. Am. Phys. Soc.* 20: 1240 (1975).
6. J. D. Callen et al., "Neutral Beam Injection into Tokamaks," *Proc. 5th Conf. Plasma Physics and Controlled Nuclear Fusion Research*, Vol. 1, p. 645 (1975).
7. D. G. McAlees, *Alpha Particle Energetics and Neutral Beam Heating in Tokamak Plasmas*, ORNL/TM-4661, Oak Ridge (November 1974).
8. S. O. Dean et al., *Status and Objectives of Tokamak Systems for Fusion Research*, WASH-1295, USAFRC, Washington, D. C. (1974).

ORNL-DWG 75-15768

Fig. 5.3. Toroidal flow velocity versus plasma radius for the ORNL EPR design parameters.

9. K. T. Tsang and E. A. Frieman, *Toroidal Plasma Rotation in Axisymmetric and Slightly Non-Axisymmetric Systems*, ORNL/TM-4948, Oak Ridge (1975).
10. J. A. Rome, J. D. Callen, and J. F. Clarke, "Fast Ion Spatial Distribution for Tangential Injection," *Nucl. Fusion* 14: 141 (1975).

6. START-UP: BREAKDOWN, PRE-IONIZATION EFFECTS, RUNAWAY PRODUCTION, COUPLED RATE EQUATIONS

6.1 Introduction

In a tokamak, gas breakdown generally results from the induced electric field. The principal heating mechanism is the plasma current. The establishment of the full value of this current may entail runaway production, increased volt-second requirements on the transformer until ionization is complete, and large resistive losses which result in intense wall bombardment. The auxiliary method of pre-ionization by microwaves at the electron cyclotron frequency may alleviate some of these problems. A set of coupled equations including plasma and molecular physics is used to model start-up, and is compared with analytic predictions.

6.2 Gas Discharge Parameters

For the purposes of comparison, ORMAK¹ is considered to be representative of present tokamaks. The ORNL EPR reference design² has the large-tokamak attribute of a moderate plasma driving voltage due to a superconducting air core transformer. The parameters of these tokamaks are listed in Table 6.1.

Table 6.1. Tokamak parameters

	<u>ORMAK</u>	<u>ORNL EPR</u>
Major radius, R (cm)	80	675
Minor radius, a (cm)	23	225
Voltage/turn (V)	40	50
Electric field, E (V/cm)	.0796	.0118
Toroidal field, B _T (G)	1.8×10^4	4.8×10^4
Pressure H ₂ , p (torr)	4.4×10^{-4}	(1-4) $\times (4.4 \times 10^{-4})$
E/p (V/cm x torr)	1.81×10^2	26.8-6.70

Ionization rates, drift velocities, and electron temperatures have been measured in linear gas discharge experiments^{3,4,5,6} and applied to the analysis of breakdown in toroidal discharges.^{7,8} The parameter α is the ionization rate per centimeter per electron, and p is pressure. The reduced parameter α/p is a function of E/p (see pp. 180-182 of Ref. 3). An analytic expression for hydrogen in the region of interest is

$$\frac{\alpha}{p} = 5.4 \exp \left(\frac{-139}{(E/p)} \right), \text{ ionizations/cm x torr} \quad (1)$$

for

$$20 < E/p < 1000$$

The value of α/p decreases rapidly for $E/p \leq 100$, so that efficient ionization requires $E/p \geq 100$; α/p asymptotes to a maximum above this level. The drift velocity of the electron ensemble in the gas, due to elastic collisions, is approximated by ⁷

$$V_D = 3.5 \times 10^5 E/p \quad (\text{cm/sec}), \quad (2)$$

from a linear extrapolation^{3,4} of data with $E/p \leq 20$. The ionization rate per electron per second is $\nu = \alpha V_D$, and per unit volume is $n_e \nu$. The electron density is n_e . The drift velocity is added to

the higher random velocities of an electron cloud. The electron temperature T_e in eV is approximated⁷ by

$$\begin{aligned} T_e &= 0.1 (E/p) \text{ eV} & \text{for } E/p < 100 \\ T_e &\approx 10 \text{ eV} & \text{for } E/p > 100 \\ & & \text{(due to ionizations)} \end{aligned} \quad (3)$$

based upon Ref. 4. A more accurate expression may be obtained from plots of the random velocity/drift velocity.³

6.3 Plasma Buildup Rate

In the absence of losses, or for low losses, the density buildup equation is

$$\frac{d}{dt} n_e = \nu n_e = \alpha v_D n_e, \quad (4)$$

where $\nu (= \alpha v_D)$ is the ionization rate. If E/p is constant, then the solution is

$$n_e = n_{e0} \exp(\alpha v_D t). \quad (5)$$

If η is defined as n_e/n_{e0} , then the buildup time, τ , is

$$\tau = \ln \eta / \alpha v_D. \quad (6)$$

Devices like ORMAK have large values of E/p and p , which result in large values of α and rapid ionization rates. As examples, $\eta = 10^2$ is taken for good pre-ionization, and $\eta = 10^{12}$ for no pre-ionization. Table 6.2 shows the machine parameters, breakdown electric field, typical values of E/p , and buildup time for these values of η . For ORMAK and EPR, $p = 4.4 \times 10^{-4}$ torr is assumed.

Table 6.2. Ionization time

	<u>V/Turn</u>	<u>Pressure (torr)</u>	<u>η</u>	<u>τ(sec)</u>	<u>$V\tau$(V-sec)</u>
ORMAK	40	4.4×10^{-4}	10^2	6.6×10^{-5}	2.6×10^{-3}
ORMAK	40	4.4×10^{-4}	10^{12}	4.0×10^{-4}	1.6×10^{-2}
EPR	50	4.4×10^{-4}	10^2	3.5×10^{-2}	1.75
EPR	50	4.4×10^{-4}	10^{12}	2.1×10^{-1}	10.5

In EPR, at this pressure, the very low pre-ionization level results in a consumption of 10% of the available volt-seconds during ionization. An efficient pre-ionization mechanism is therefore desirable, especially at higher filling pressures.

6.4 Choice of Pressure in EPR

From the above considerations efficient breakdown in EPR requires $E/p \geq 100$, in which case the electrons would have a temperature of approximately 10 eV due to ohmic heating during the very early breakdown state. For $\vec{E}_{EPR} = 1.2 \times 10^{-2}$ V, this implies a maximum filling pressure of 1.2×10^{-4} torr at $E/p = 100$. This low pressure, if ionized uniformly in the chamber, results in a low-density plasma. However, many large tokamak start-up proposals include small radius start-up by moving material limiter,^{9,10} vertical field control,¹¹ or magnetic limiters.¹² In these cases, a small plasma will

rapidly ionize the gas in the chamber, increasing the density by the ratio of the plasma volume to chamber volume.

High filling pressures imply a long pre-ionization time, requiring efficient particle containment. The question of confinement time in a torus and the associated equilibrium state is extremely complex in the regime of zero or very low current, and will not be considered in detail here. Since mechanisms such as field errors or electric fields could cause rapid losses, a high ionization rate seems prudent. However, too high a value of E/p can lead to runaway electrons.

6.5 Runaway Electrons

Runaway electrons result from free-fall in the toroidal electric field. Experiments in several tokamaks, including ORMAK,^{13,14} show that the runaway flux depends on conditions of density and temperature which are only imprecisely known. Theoretical estimates can be made in the presence of plasma or neutral gas.

The critical field¹⁵ in a very weakly ionized plasma (neutral gas) is

$$E_{cn} = 4\pi e^3 n_0 (Z/2.72) \bar{\epsilon} \approx 7 \times 10^{-15} n_0 \quad (\text{V/cm}) , \quad (7)$$

where n_0 is the neutral gas density. Hence, the critical field is $E_{cn}/p = 248$. Operating a tokamak at $E/p > 100$ gives a high probability of a significant runaway production; in fact, even values of E/p lower than this by an order of magnitude can lead to runaways.

The critical field in a plasma¹⁶ is

$$E_{ci} = \frac{4\pi^2 n_e}{T_e} \ln \frac{mv_e^2 D}{e^2} \approx 2.6 \times 10^{-13} \frac{n_e}{T_e} \ln \frac{4 \times 10^9 (T_e)^{1/2} mv_e^2}{n_e^{1/2}} \quad (8)$$

where v_c is the critical velocity for runaways, and T_e and mv_c^2 are in electron-volts. The logarithmic term has a value of ~ 10 for typical conditions, and it is relatively insensitive to parameters. Where $E/p \approx 100$, $T_e \approx 10$ eV, so that

$$E_{ci}/p \approx 9.2 \times 10^3 (n_e/n_0) . \quad (9)$$

To avoid runaways requires $E_{ci}/p \gg E/p \approx 100$, which implies that $n_e/n_0 \gg 10^{-2}$. At a typical filling pressure of 4×10^{-4} torr, $n_e \gg 1.5 \times 10^{11} \text{ cm}^{-3}$.

It should be noted that in the absence of a movable or magnetic limiter, it will probably be necessary to pre-ionize the entire volume to prevent runaways, since an electric field is induced over the entire plasma volume until the current builds up. Electrons in the un-ionized region can run away if $E/p \approx E_{cn}/p$.

6.6 Pre-ionization by Microwave ECRH

A small-scale microwave breakdown experiment¹⁶ has been interpreted from a theoretical paper¹⁷ on maximum electron temperatures. The experiment was conducted in a metal torus of $R_0 = 22$ cm, $a = 4.4$ cm, with a toroidal field on axis of 3×10^3 G, and no plasma current. The electric field of the μ -waves E_0 determines a maximum electron energy¹⁷

$$W_1 = 2m_e c^2 \left(\frac{eE_0}{m_e \omega_0 c} \right)^{2/3} \quad (10)$$

where E_0 is in statvolts/cm, B_0 is the resonant field in gauss, and ω_0 is the electron gyrofrequency in radians/sec. This limit occurs because of relativistic detuning of the resonant frequency.

In a resonant cavity, of specified ω , the electric field is

$$E_0 (\text{statvolts/cm}) = \left(\frac{8\pi mcQ}{eB_0 (\text{gauss})} \frac{\text{Power (ergs/sec)}}{\text{Volume (cm}^3)} \right)^{1/2} \quad (11)$$

If this model is applicable, then a maximum electron energy may be calculated as shown in Table 6.3.

Table 6.3. Maximum electron energies

	<u>ORMAK</u>	<u>EPR</u>
Q	10^4	10^4
B	1.8×10^4 G	4×10^4 G
λ_μ	.60 cm	.27 cm
ω_0	3.16×10^{11}	7.03×10^{11}
P	10^{12} ergs/sec	10^{13} ergs/sec (1 MW)
V	8.35×10^5 cm ³	6.75×10^8 cm ³
E_0	.976 statvolts/cm	7.3×10^{-2} statvolts/cm
W_L	1.46×10^3 eV	152 eV

In fact, runaway electrons due to microwave heating are experimentally observed, so the validity of the model is in question. If a maximum temperature does exist, then the power input is limited by the ionization rate at which additional electrons are created.

6.7 Heating Volume

The efficiency of microwave pre-ionization depends upon the effective heating volume. The minimum pressure for breakdown is calculated¹⁶ by setting the ionization rate $N_n \bar{\sigma} v_i$ equal to the electron $vB \times B$ drift loss rate. The value N_n is the neutral density; $\bar{\sigma} v_i$ is the ionization rate¹⁸ equal to 7×10^{-9} cm³ sec⁻¹ at $T_e = 10$ eV. At 100 eV, $\bar{\sigma} v_i$ equals 4×10^{-8} cm³ sec⁻¹ and decreases slowly at higher T_e to 3×10^{-8} cm³ sec⁻¹ at 1000 eV. The curvature drift is vertical, implying that both ionization and loss occur in a thin annulus at the resonant surface in the torus. Only a small volume would be ionized. Normally the curvature drift would result in an outwards polarization drift. However, if the drift current can flow through a metal chamber wall, an equilibrium may be obtained.

In order to ionize the complete volume, some mechanism must be postulated by which electrons drift into the resonant zone, are heated, drift out, and ionize the remaining volume. Several effects, such as electric fields, density gradients, force balance drifts, etc., could occur. Also, if the pre-ionization microwave pulse were applied at the same time as the current driving field, a small current would suffice to cause particle drifts on flux surfaces. The subsequent pre-ionization over the entire volume would facilitate breakdown.

6.8 Microwave Pre-ionization Rate Equations

The density and temperature rate equations for negligible losses are

$$\frac{dn_e}{dt} = n_e n_0 (\sigma v)_i \quad (12)$$

$$\frac{d}{dt} \left(\frac{3}{2} n_e T_e \right) = - E_i n_e n_0 (\sigma v)_i + P_\mu \quad (13)$$

where $(\sigma v)_i$ is the ionization rate of H_2 by electrons, E_i is the energy required to produce an electron-ion pair (~ 93 eV) at low electron temperatures,³ and P_μ is the absorbed microwave power per unit volume. The solutions to these equations, for constant $(\sigma v)_i$ and n_0 , are

$$n_e = n_{e0} \exp(t n_0 (\sigma v)_i) = n_{e0} \exp(t/\tau_b) \quad (14)$$

$$\frac{dT_e}{dt} = -n_0 (\sigma v)_i (T_e + \frac{2}{3} E_i) + \frac{\frac{2}{3} P_\mu}{n_{e0} \exp(t n_0 (\sigma v)_i)} \quad (15)$$

The parameter τ_b is the exponential-folding time for the density increase. Initially, the second term in the T_e equation dominates, due to low density, and T_e increases rapidly to a maximum value due to relativistic detuning of the resonance. Eventually $dT_e/dt = 0$, and then

$$P_\mu (\text{initial}) = \frac{3}{2} n_{e0} \exp(t n_0 (\sigma v)_i) \times n_0 (\sigma v)_i (T_e + \frac{2}{3} E_i) \quad (16)$$

If the temperature saturation model is applicable, we can deduce the qualitative behavior of the system by an approximate analytic solution of the rate equations. As the density rises, eventually the full power will become available. In the early states, T_e is typically between 100 eV and 1 keV, so that $(\sigma v)_i$ is roughly constant. Thus, the characteristic density level at which the exponential buildup changes to a slower buildup is

$$n_{eb} = \frac{2}{3} \tau_b P_\mu / (T_e + \frac{2}{3} E_i) \quad (17)$$

At this level, T_e begins to drop, and below 100 eV the $(\sigma v)_i$ rate drops rapidly. If the temperature change is slow compared to the density change, then

$$\frac{d}{dt} n_e = P_\mu / (E_i + \frac{3}{2} T_e) \quad (18)$$

If the pre-ionization time, τ_p , is several times τ_b , then the final density at the end of the pre-ionization is approximately

$$n_{ep} = P_\mu \tau_p / (E_i + \frac{3}{2} T_e) \quad (19)$$

and the required microwave power is calculated from the desired density. If T_e has dropped far below 100 eV, then it can be neglected in the power calculation. If T_e is large, then E_i is reduced, so $E_i + 1.5 T_e$ is roughly constant below $T_e = 100$ eV.

The most efficient pre-ionization will occur if the pre-ionization time is less than the particle containment time. The scaling for turbulent Bohm diffusion is similar to the single-particle drift time τ_D , which is

$$\tau_D = \frac{aRB}{10^8 (2T_e)} \text{ (cm, G, eV, sec)} \quad (20)$$

In EPR, at $T_e = 100$ eV, τ_D equals .36 sec, and increases at lower T_e . The rate is therefore determined by the desired pre-ionization time. To obtain a low flux swing, we choose a time of .021 sec (1.5 V-sec) and an electron density N_e/N_n high enough to avoid runaways ($N_e/N_n \sim 0.1$), so that the density is $1.56 \times 10^{12} \text{ cm}^{-3}$. With $(E_i + 1.5 T_e) \approx 100$ eV, the specific power is $7.5 \times 10^{15} \text{ eV/cm}^3 \text{ sec}$, and the total power is 162 kW. This assumes 100% absorption.

6.9 Coupled Rate Equations

In this section we consider the problem of the complete ionization and heating of the gas in a toroidal discharge. Initially, the torus will be filled with molecular hydrogen (or its isotopes D and T) with some degree of pre-ionization. A complete description of the breakdown process requires density and temperature equations for H_1^0 , H_2^0 , H_1^+ , H_2^+ , and electrons, plus an equation for the current, including inductive and resistive contributions. The reaction rate for these processes is derived from data in Refs. 3, 8, and 18-27. These equations are similar to those of Ref. 28, and are an extension of those in Ref. 29, which includes impurities, but includes only H_1^0 of the hydrogen species. The volumes are:

$$V_0: \text{Volume available to } e, H_1^+, H_1^0 = 2\pi R_0 \pi r_0^2$$

$$V_1: \text{Volume of torus chamber} = 2\pi R_0 \pi r_1^2$$

$$V_2: \text{Volume available to } e, H_1^+, H_1^0, H_2^+, H_2^0 = 2\pi R_0 \pi (r_a^2 - r_2^2)$$

The plasma radius r_0 is less than the chamber radius r , in the case of a material or magnetic limiter. When the plasma becomes so dense or hot that H_2^0 molecules cannot penetrate to the center of the plasma, r_2 is non-zero. However, during the breakdown stage, it will be shown that the electron temperature is so low, 1-2 eV, that the H_2^0 can easily fill the whole volume. Also, the energy balance is dominated more by molecular physics than by transport (under certain assumptions), so that a zero-spatial-dimension model provides an extremely good description. The H_1^+ , H_2^+ , H_1^0 , and H_2^0 densities and temperatures are given in order as n_{1+} , n_{2+} , n_{10} , n_{20} , and T_{1+} , T_{2+} , T_{10} , T_{20} . The volume-averaged electron densities in the V_0 and V_2 volumes are n_{e0} and n_e , which are equal if $V_0 = V_2$. The reactions rates S_1 through S_9 and S_{2x} , S_{3x} , S_{4x} are listed in Table 6.4. These include processes for ionization, excitation, dissociation, and charge exchange. In the particle loss rate terms, which also cause an energy loss, the confinement time is Bohm-like, which is similar to the single-particle drift rate. The protons diffusing to the wall recombine to return as H_1^0 with probability ψ_1 , and as H_2^0 with probability ϕ_1 . Similarly, the H_2^+ ions diffusing to the wall recombine as H_1^0 with probability ψ_2 , and as H_2^0 with probability ϕ_2 .

Table 6.4. Rate coefficients $S(\text{m}^3 \text{ sec}^{-1})$ and processes

Rate	Process
S_1	$e + H_1^0 \rightarrow e + H_1^+ + e$
S_2	$e + H_2^0 \rightarrow e + H_2^+ + e$
S_3	$e + H_2^0 \rightarrow e + H_1^+ + H_1^0 + e$
S_4	$e + H_2^0 \rightarrow e + H_1^0 + H_1^0$
S_5	$e + H_2^0 \rightarrow e + H_1^+ + H_1^+ + e + e$
S_6	$e + H_2^+ \rightarrow e + H_1^+ + H_1^0$
S_7	$H_1^+ + H_1^0 \rightarrow H_1^0 + H_1^+$
S_8	$H_1^+ + H_2^0 \rightarrow H_1^0 + H_2^+$
S_9	$H_2^+ + H_2^0 \rightarrow H_2^0 + H_2^+$
S_{2x}	$e + H_1^0 (n = 1) \rightarrow e + H_1^{0*} (n = 2)$

6.10 Neutral Temperature Equations

The equations for the neutral temperatures T_{10} and T_{20} assume that the neutrals are reflected off the walls with zero loss of energy. Alternatively these equations can be "turned off" and the neutral temperature can be fixed at the wall temperature, approximately 10^{-1} eV. These two assumptions can lead to very different results. In one case, the ion temperatures and neutral temperatures rise together. In the other, the ion temperature falls toward the wall temperature because of charge exchange. The detailed behavior depends upon particle and energy wall reflection coefficients.

A further set of processes not yet considered in detail is that of the elastic collisions of ions with neutrals and neutrals with neutrals. The $H_2 - H_2$ collision cross section is given in Ref. 30 (Fig. 16.75, p. 1458) and the $H_1^0 - H_2^0$ cross section is given in Ref. 30 (Fig. 16.71, p. 1452). At energies of 1 eV and below, the approximate value of σ is $5 \times 10^{-15} \text{ cm}^2$. The cross sections for collisions of H_1^+ , H_2^+ , and H_3^+ on H_2^0 are derived from Ref. 8 (Fig. 2.55, p. 43). The values are all in excess of 10^{-15} cm^2 for energies less than 9 eV, and they increase rapidly at low energies. Thus at high neutral densities, the mean free path for these collisions may be much less than the machine radius, so that the neutral temperature could rise above the wall temperature. A high value of the wall reflection coefficient would enhance this effect. The data from Ref. 31 (Fig. 2) show that the reflection coefficients of neutral atoms increase rapidly below 10 eV, and are above 0.7. Reference 31 (Fig. 3) also shows that the reflection coefficients increase as the angle of incidence changes from normal incidence. Therefore, high neutral temperatures may result.

6.11 Analytic Model of Minimum E/p for Complete Ionization

An analytic model may be developed to predict the minimum E/p required to achieve complete ionization. The model considers a system of electrons, H_2^+ ions, and H_2^0 neutrals. The electrons are heated by ohmic heating, ionize the H_2^0 to H_2^+ , and heat the H_2^+ by Coulomb transfer. The warm H_2^+ ions lose their energy to the walls by rapid charge exchange with H_2^0 . The analysis is the same for H_1^+ on H_1^0 . The ohmic power ($\text{keV/m}^3 \text{ sec}$) is

$$P_{OH} = \frac{E^2}{n(1.6 \times 10^{-16})} = 3.79 \times 10^{24} \frac{E^2 T_e^{3/2}}{Z_{eff}^2 n \Lambda} \quad (21)$$

If $T_{2+} \approx T_{20} \approx 0$, then the electron-ion Coulomb power is

$$P_{coul} = \frac{n_e^2 \Lambda n \Lambda (.77 \times 10^{-19})}{\sqrt{T_e}} \quad (22)$$

At low density, the temperature would increase rapidly, except that ionization and excitation reactions limit it to a few eV. The reaction rates are extremely strong functions of temperature in this range. The value of $\Lambda n \Lambda$ is approximately 10, and $Z_{eff} = 1$. If the plasma is to be completely ionized, the ohmic power must always exceed the Coulomb power, giving

$$\frac{ET_e}{n_e} > 1.425 \times 10^{-21} \frac{\text{V} \cdot \text{m}^2}{\text{kev}} \quad (23)$$

The electron temperature is roughly approximated from low ionization gas discharge physics (see Ref. 6, p. 3):

$$T_e(\text{keV}) = 10^{-6} E(\text{V/m})/p(\text{torr}) . \quad (24)$$

If the final electron density is approximated by the filling pressure $p = n_e/3.54 \times 10^{22}$, we obtain

$$\frac{E}{p} > 7102 \frac{\text{V/m}}{\text{torr}} = 71 \frac{\text{V/cm}}{\text{torr}} . \quad (25)$$

This is only approximate, because the H_2^+ temperature will rise above zero, and the H_2^0 will disappear as full ionization is approached. Even so, this indicates a threshold for full ionization. Similarly, if E/p is below this level, the maximum obtainable density is:

$$\frac{n_e}{n_0} = \frac{E/p}{7102} , \quad E/p < 7102 . \quad (26)$$

These predictions will be checked in the exact numerical simulations.

6.12 Numerical Results: Application to EPR

The numerical results from the start-up code for EPR are shown in Table 6.5 for three illustrative cases. In each case, the initial conditions are given. The second set of parameters is either approaching equilibrium or at the point of complete ionization. The neutral temperatures are fixed at the wall temperature, and the recycling parameters cause H_1^+ to return entirely as H_1^0 , and H_2^+ as H_2^0 . The voltage is constant at 50 V, the maximum available from the superconducting air core. The operating density of EPR is $0.8 \times 10^{20} \text{ m}^{-3}$, so a filling pressure of nH_2^0 of $0.4 \times 10^{20} \text{ m}^{-3}$ might be desirable. The E/p value is 10.4, much lower than the minimum predicted by the analytic theory. To test this theory shown in Case I in Table 6.5, the filling pressure is set so as to give $E/p = 70$, just below the predicted minimum. After .030 sec, the ion density has saturated at about 20% of the initial neutral density. In this case, the main energy loss is by charge exchange from H_1^+ rather than H_2^+ . The electron temperature and current are slowly falling, while the density is rising with a time constant of 11 sec. Thus we tentatively conclude that there is a threshold for ionization at approximately the level of the analytic prediction. In Case II, the filling pressure was lowered by a factor of 6, raising E/p to 417. In this case ionization proceeds rapidly, half the plasma being ionized after .020 sec. The electron temperature is rising at a rate of .007 keV/sec and the current at 2.15 MA/sec. As the temperature rises, the plasma voltage drop will be due to inductive rather than resistive effects. At full ionization, the density would be only $0.2 \times 10^{19} \text{ m}^{-3}$, a factor of 40 below the desired level.

One method of overcoming this difficulty is to employ a small radius start-up. In Case III, the plasma volume is 1/4 of the chamber volume. The ionization proceeds more rapidly, the final density is $0.8 \times 10^{19} \text{ m}^{-3}$, and both the temperature and current are rising rapidly. The problem is then to expand the plasma, while feeding gas to provide the density. This will be the subject of future analysis.

The time-dependent behavior of these cases is shown in Figs. 6.1, 6.2, and 6.3.

Fig. 6.1. EPR start-up: $T_{10,20} = 0.1$ eV, $n_{20} = 0.6 \times 10^{19} m^{-3}$.

Fig. 6.2. EPR start-up: $T_{10,20} = 0.1$ eV, $n_{20} = 0.1 \times 10^{19} \text{ m}^{-3}$.

Fig. 6.3. EPR start-up: $T_{10,20} = 0.1$ eV, $n_{20} = 0.1 \times 10^{19} \text{ m}^{-3}$, $v_1/v_0 = 4$.

Table 6.5. EPR start-up parameters

$R_0 = 6.75 \text{ m}$	$T_{\text{well}} = .0001 \text{ keV}$	$\psi_2 = 0$
$a_0 = 2.25 \text{ m}$	$V = 50$	$\phi_1 = 0$
$B_T = 4.8 \text{ T}$	$\psi_1 = 1$	$\phi_2 = 1$
$T_{10} = T_{20} = \text{constant}$		
$T_e(0) = .002 \text{ keV}$	$I(0) = .1 \times 10^5 \text{ A}$	
$T_{1+}(0) = T_{2+}(0) = T_{10}(0) = T_{20}(0) = .0001 \text{ keV}$		
$n_{1+}(0) = n_{2+}(0) = n_{20}(0) = .1 \times 10^{17} \text{ m}^{-3}$		

Parameter	Case I	Case II	Case III
$n_{10}, \text{ m}^{-3}$	$.6 \times 10^{19}$	$.1 \times 10^{19}$	$.1 \times 10^{19}$
Time, sec	.030	.020	.018
V_1/V_0	1	1	4
$p_{20}, \text{ torr}$	1.69×10^{-4}	2.82×10^{-5}	2.82×10^{-5}
$E/p_{20}, \text{ V/cm-torr}$	70	417	417
$T_e, \text{ keV}$	$.159 \times 10^{-2}$	$.270 \times 10^{-2}$	$.105 \times 10^{-1}$
$T_{1+}, \text{ keV}$	$.867 \times 10^{-3}$	$.146 \times 10^{-2}$	$.102 \times 10^{-1}$
$T_{2+}, \text{ keV}$	$.878 \times 10^{-3}$	$.148 \times 10^{-2}$	$.671 \times 10^{-2}$
$n_e, \text{ m}^{-3}$	$.126 \times 10^{19}$	$.119 \times 10^{19}$	$.807 \times 10^{19}$
$n_{1+}, \text{ m}^{-3}$	$.924 \times 10^{18}$	$.117 \times 10^{19}$	$.807 \times 10^{19}$
$n_{2+}, \text{ m}^{-3}$	$.335 \times 10^{18}$	$.213 \times 10^{17}$	0
$n_{10}, \text{ m}^{-3}$	$.270 \times 10^{19}$	$.825 \times 10^{18}$	$.900 \times 10^{14}$
$n_{20}, \text{ m}^{-3}$	$.387 \times 10^{19}$	$.112 \times 10^{16}$	0
$\dot{T}_e, \text{ keV/sec}$	$-.241 \times 10^{-1}$	$.721 \times 10^{-2}$	$.310 \times 10^1$
$\dot{n}_h, \text{ m}^{-3} \text{ sec}^{-1}$	$.102 \times 10^{20}$	$.636 \times 10^{20}$	$.203 \times 10^{18}$
$\dot{I}, \text{ A/sec}$	$-.500 \times 10^6$	$.215 \times 10^7$	$.632 \times 10^7$
$I, \text{ A}$	$.622 \times 10^5$	$.894 \times 10^5$	$.100 \times 10^6$

A significant change occurs in the results when the neutral temperatures are allowed to vary. In Figs. 6.4-6.7, the neutral temperature equations were used with the wall-collision energy-loss model. The plasma volume equalled the chamber volume, and the four filling pressures were 0.1, 0.3, 0.6, and $1.2 \times 10^{19} \text{ m}^{-3}$. In Fig. 6.4, the plasma ionized more rapidly than in Fig. 6.3. Similarly, Fig. 6.5 shows rapid ionization at a higher pressure. At higher pressures, in Figs. 6.6 and 6.7, ionization is proceeding at a much slower rate. The main power loss is now due to excitation radiation from the S_{2x} term. The temperatures have all equilibrated. If the wall-collision model is correct, then ionization at higher pressures is possible.

Even these higher pressures are insufficient for operation at high plasma density. This work is still in early stages, but it seems that low density or small radius start-up will be necessary, unless much higher voltages can be applied during the plasma breakdown stage.

Acknowledgment

A. C. England and J. T. Hogan made valuable suggestions in the microwave breakdown and coupled rate-equations section. T. Uckan aided in small radius start-up studies.

Fig. 6.4. EPR start-up: $T_{10,20}$ variable, $n_{20} = 0.1 \times 10^{19} \text{ m}^{-3}$.

 Fig. 6.5. EPR start-up: $T_{10,20}$ variable, $n_{20} = 0.3 \times 10^{19} m^{-3}$.

Fig. 6.6. EPR start-up: $T_{10,20}$ variable, $n_{20} = 0.6 \times 10^{19} \text{ m}^{-3}$.

Fig. 6.7. EPR start-up: $T_{10,20}$ variable, $n_{20} = 0.12 \times 10^{20} \text{m}^{-3}$.

REFERENCES

1. J. D. Callen, J. F. Lyon, and M. Murakami (eds.), *ORMAK Plasma Confinement and Neutral Beam Injection Heating (Experiment and Theory)*, ORNL/TM-4790, Oak Ridge (January 1975).
2. M. Roberts and E. S. Bettis (eds.), *Oak Ridge Tokamak Experimental Power Reactor Study Reference Design*, ORNL/TM-5042, Oak Ridge (November 1975).
3. A. von Engel, *Ionized Gases*, pp. 180-182, Clarendon Press, Oxford, England, 1965.
4. S. C. Brown, *Basic Data of Plasma Physics*, Technology Press, Cambridge, Mass., 1965.
5. S. C. Haydon (ed.), *Discharge and Plasma Physics*, University of New England, Armidale, Australia, 1964.
6. R. Papoular, *Electrical Phenomena in Gases* (English translation), Iliffe Books Ltd., London, England, 1965.
7. R. Papoular, *The Genesis of Toroidal Discharges*, EUR-CEA-FC-769, Fontenay-aux-Roses, France (April 1975).
8. D. C. Schram and J. Rem, *Eddy Currents in the ALCATOR Tokamak*, Rijnhuizen Report 75-88, Appendix C, Jutphaas, Netherlands (March 1975).
9. S. Itoh et al., "Dynamic Limiter Experiments in JFT-2 Tokamak," *Nucl. Fusion Lett.* 14: 581 (1974).
10. J. P. Girard, M. Khelladi, and D. A. Marty, "Low Density Tokamak Reactor Ignition by Fast-Neutral-Atom Injection and Buildup at Higher Density by Use of a Moving Limiter," *Nucl. Fusion* 13: 686 (1973).
11. F. B. Marcus et al., "A Magnetic Shielding System for the Tokamak Experimental Power Reactor," *Proc. 6th Symposium on Engineering Problems of Fusion Research*, p. 1110 (1976).
12. K. W. Weiner et al., "Studies of MHD Equilibria in a Divertor Tokamak," *Bull. Am. Phys. Soc.*, 20: 1279 (October 1975).
13. D. A. Spong, J. F. Clarke, and J. A. Rome, *Relativistic Electron Production in the ORMAK Device*, ORNL/TM-4120, Oak Ridge (June 1973).
14. D. A. Spong and J. F. Clarke, *Runaway Electron Dynamics*, ORNL/TM-4432, Oak Ridge (December 1973).
15. A. V. Gurevich, "On the Theory of Runaway Electrons," *Sov. Phys.-JETP* 15(5): 094-912 (May 1961).
16. A. I. Anisimov, N. I. Vinogradov, and B. P. Poloskin, *Resonance Microwave Breakdown in a Toroidal Magnetic Field* (trans.), ORNL-tr-4078, Oak Ridge (1975).

17. A. Hakkenberg and M. Pitt. Weenink, "The Solution of the Relativistic Equation of Motion of a Charged Particle in a Homogeneous Magnetic and a Rotating Electric Field Perpendicular to Each Other," *Physica* 30: 2147-2152 (1964).
18. R. L. Freeman and E. M. Jones, *Atomic Collision Processes in Plasma Physics Experiments*, CLM-R137, UKAEA, Culham Laboratory, Abingdon, Oxfordshire, England (May 1974).
19. E. Hinnov, A. S. Bishop, and F. W. Hofmann, *Interpretation of Atomic Hydrogen Light in the C Stellerator*, MATT-270, Princeton Plasma Physics Laboratory, New Jersey (May 1964).
20. R. G. Alsmiller, Jr., *Dissociation and Ionization of H_2^+ by Electrons and Protons*, ORNL-3232, Oak Ridge (January 1962).
21. C. F. Barnett, J. A. Ray, and J. C. Thompsen, *Atomic and Molecular Cross Sections of Interest in Controlled Thermonuclear Research*, ORNL-3113, Oak Ridge (revised, August 1964).
22. S. C. Brown, *Basic Data of Plasma Physics*, Technology Press, Cambridge, Mass., 1965.
23. B. Peart and K. T. Dolder, "The Production of De-Excited H_3^+ Ions and Measurements of the Energies of Two Electronically-Excited States," *J. Phys. B: Atom. Molec. Phys.* 7(12): 1567-1573 (1974).
24. D. W. Vance and T. L. Bailery, "Inelastic Collisions of H_2^+ and N_2^+ Ions With Hydrogen Molecules," *J. Chem. Phys.*, 44(2): 486-493 (January 1966).
25. B. Peart and K. T. Dolder, "Measurements of Cross Sections for the Production of Protons by Collisions Between Electrons and Vibrationally De-Excited H_3^+ Ions," *J. Phys. B: Atom. Molec. Phys.* 8(8): L143-L145 (1975).
26. B. Peart and K. T. Dolder, "Measurement of the Dissociative Recombination of H_3^+ Ions," *J. Phys. B: Atom. Molec. Phys.* 7(14): 1948-1952 (1974).
27. H. W. Drawin, "Collision and Transport Cross Sections," EUR-CEA-FC-383 (revised), Fontenay-aux-Roses, France (January 1967).
28. E. Hinnov, A. S. Bishop, and H. Fallon, Jr., "Computer Description of the C Stellerator Ohmic Heating Discharges in Hydrogen," *Plasma Phys.* 10: 291-307 (1968).
29. R. J. Hawryluk and J. A. Schmidt, "Start-up Phase of a Large Tokamak, Paper 3B12, *Bull. Am. Phys. Soc.*, Series II, 20(10) (October 1975).
30. H. S. W. Massey, *Electronic and Ionic Impact Phenomena, Volume III: Slow Collisions of Heavy Particles*, Clarendon Press, Oxford, England, 1971.
31. O. S. Oen and M. T. Robinson, "Computer Studies of the Reflection of Light Ions from Solids," *Nucl. Instrum. Methods*, 132: 647 (1976).

7. FUELING

Fueling is considered a necessary procedure for the operation of a fusion reactor in that it balances particle diffusion losses at the intended operating conditions, i.e., it maintains the density of the plasma, and to some extent controls the reactivity of the device. Moreover, a method of fueling must be sought which avoids excessive neutral fluxes at the chamber wall due to charge-exchange processes.

7.1 Plasma Recycling

Plasma-wall recycling may be inadequate since the ionization of wall-neutralized refluxing ions and reflected charge-exchange neutral atoms prevents their deep penetration back into the core region. Only the hottest of ions in the plasma, once charge-exchanged into neutrals, have mean free paths against ionization and further charge-exchange as large as the plasma. In an EPR-sized plasma, these very fast neutrals are most likely to be buried in the wall rather than reflected.^{1,2,3} A large diffusive mixing of edge-refluxing particles into the plasma core is necessary for this method to succeed.

7.2 Neutral Gas Blanket

Neutral atom fluxes created by the dissociation of gas molecules, adsorbed to the chamber walls or injected into the chamber, are likewise attenuated by ionization and the combined effects of charge exchange.

To prevent damaging neutral fluxes at the chamber wall, a dense cold plasma must be interposed between the near-wall gas region and the hot plasma core to cool wall-directed neutrals. This, however, increases the attenuation of core-bound neutrals. The charge-exchange process serves in effect to increase the path length of neutral atoms between the wall and the core, thus increasing the attenuation of these atoms due to ionization. The existence of the cool plasma layer is determined by the relative rates at which neutrals ionize in this region and at which the ionized products diffuse out of this region. Some question exists as to the pertinent diffusion rate, though it seems safe to believe that longer particle-confinement times will be accompanied by the longer persistence of such layers.

One adverse effect of fueling the plasma center by ionizing Franck-Condon neutrals which are generated at the plasma edge is a fast neutral flux at the chamber wall which may increase intolerably with the fueling rate.

7.3 Pellets

There is some indication that fueling by the ballistic injection of fuel pellets may to some degree avoid the problem of excessive wall loading by neutrals. A high-density gas forms around the pellet adjacent to the pellet surface and expands radially into the plasma. The neutral gas conserves the momentum from the pellet motion; as it expands, it travels with the pellet. The electrons from the plasma drive the evaporation responsible for the gas which they also heat and accelerate. The evaporation takes place in the presence of the neutral gas shielding layer. Inelastic collisions with gas molecules attenuate the electron power flux at the surface of the pellet from its values in the gas shield. The gas is dissociated and ionized by reactor electrons in a region some few pellet radii from the pellet surface, forming a dense warm plasma (a temperature of a few electron-volts) which is heated by the reactor ions. The ions slow down rapidly in the peripheral regions of the expanding high-density ablatant. Warm secondary electrons, generated profusely in this ionization

region by the inelastic collisions of reactor electrons with gas molecules and atoms, have a high mobility in the ablatant and effectively short out the radial electric fields generated by the inward-moving electrons. A negative potential arises near the pellet surface where ionization does not occur. Recombination of cold ions and electrons can increase this potential and affect the ablation process by decreasing the inward flux of electrons. The rate at which the newly-formed fuel plasma is "scraped off" by the forward motion of the pellet-ablatant ensemble through the magnetic field determines the generation rate of neutral charge-exchange atoms.

One can argue that if the fuel plasma is deposited locally and is highly ionized, the density of charge-exchange centers is highest in the region where the expanding ablatant ionizes. Since the density of the ablatant in this region is quite high, the charge-exchange neutrals are likely to be slow, due to the slowing-down of the ions or of the fast charge-exchange products. The products then are of little consequence to neutral wall loading, as their ionization is imminent.

If, on the other hand, the fuel plasma is carried some distance with the pellet by particle drift motion caused by the electric field generated by the motion of the plasmoid across field lines,^{4,5} then the fuel plasma itself serves as a shield, slowing down incident ions before they can charge-exchange with the less ionized ablatant nearer the pellet surface.

A detailed analysis of these processes is in progress, but experimental results will be necessary to confirm that with respect to neutral wall loading in an EPR device, fueling by pellets holds a distinct advantage over fueling by gas injection. Implicit in the remainder of Sect. 7 is the assumption that such a distinct advantage exists and justifies a thorough study of pellet injection and ablation as a means of fueling large reactor systems.

7.3.1 A Steady-State Formulation of the Pellet-Plasma Interaction

The surface of a solid-hydrogen pellet immersed in a tokamak plasma is subjected to charged and neutral particle fluxes which cause evaporation, dissociation, and ionization. The pellet surface is progressively eroded or ablated away exposing the solid beneath until the pellet ceases to exist. We are concerned with the lifetime of pellets insofar as they determine the velocities required to place the pellet in a region of the plasma surface where we wish to introduce a particle source.

In present-day tokamak plasmas, the energy flux at the surface of a pellet is principally caused by electrons from the background plasma which impinge upon it.⁶ This is due to two facts. First, the electrons in a plasma where $T_e \sim T_i$ are moving faster than the ions, or neutral species; and second, the slowing-down mean free path of the electrons in the dense gas and plasma which emanates from the region around the pellet is much longer than that of the ions.

Quantitatively, the energy flux of electrons in the outermost regions of the ablatant is higher than the ion flux by $v_e/v_i \sim \sqrt{m_i/m_e}$, and the energy flux of electrons at the surface is higher than the energy flux of ions there by $\sim \sqrt{m_i/m_e} e^{-\int ds/\lambda_{esd}} e^{-\int ds/\lambda_{isd}}$, where $\lambda_{esd}/\lambda_{isd} \sim (m_i/m_e)(KT_e^2/KT_i^2)$ and the integration is taken along the particle path through the ablatant. In a reactor, energetic alpha particles may play a role in the ablation process; their suprathermal energy tends to offset the effect of their larger mass in their slowing down. In the following, we will treat only the electron thermal component of the energy flux.

We model the pellet plasma interactions as the simultaneous propagation of two concentric deflagration waves: an inner evaporation wave at the solid gas interface and an outer ionization-dissociation wave at a gas-plasma interface.

The inner wave is a well-defined deflagration wave. Its thickness as given by the range of impinging electrons in the solid is very much smaller than the size of the pellet for modest-size pellets ($> 10\mu$). The properties of the material flow field in a frame moving with this wave change abruptly from solid to dense gas through a transition zone. A Chapman-Jouget expansion is assumed.⁸

The outer wave is a somewhat thicker deflagration wave which propagates through the expanding gas ionizing as it goes. Though the penetration depth of electrons in the external gas flow field is typically as large as the flow field itself, we are justified in considering the ionization wave as a deflagration process. This is, in a sense, similar to the inner wave if we consider the following argument:

The electrons incident on the dense gas surrounding the pellet undergo many small angle collisions as they stream toward the pellet surface. The degradation of their energy can be described by the loss function $L(E)$,^{7,8} which accounts for all energy loss processes between the electrons and the H_2 molecules in the gas flow. The efficiency for energy loss by processes which ionize or dissociate molecules, as well as the total loss function itself, becomes very small at electron energies below about 40 eV. The flowing gas shield set up by the pellet is thus rather transparent to low-energy electrons which cause vibrational excitations and rather more opaque to higher energy electrons which cause ionization.

Considering this fact, we define a diffuse boundary at r_g beyond which the gas flow is stipulated to be fully ionized by the electron flux. Within this radius, we describe the gas flow as an expanding free jet which is sonic at the pellet surface. We balance mass, momentum, and energy across the two boundaries. We use the properties of the flow to determine gas density profiles (or gas flux profiles) in the region intermediate between the two boundaries. This defines the rates at which the boundaries move in a manner self-consistent with the approximations made. This is done as a first iteration towards the determination of the gas flow including the effects of electron energy deposition in the gas flow itself. Calculation shows that in reactor-grade plasmas, the effect of the local deposition of electron energy is sufficient to drive the gas flow sonic throughout the gas flow field.

The approximations implicit in the pellet model are:

- (1) Ionization and dissociation losses are responsible for the inward motion of the outer gas-plasma boundary,
- (2) the total losses between the outer and inner boundary are responsible for the degradation of the electron energy flux impinging on the inner solid-gas boundary, and
- (3) all the energy flux across the inner boundary drives the evaporative deflagration wave.

The coupled nonlinear equations describing this model can then be written:

- (1) solid-gas boundary motion

$$\frac{d}{dt} r_p = - \frac{\tilde{\Gamma}_e n_{ev} \delta_g}{n_p (w_{ev} + \Delta w)} \quad (1)$$

- (2) gas-plasma boundary motion

$$\frac{d}{dt} r_g = - \frac{\tilde{\Gamma}_e n_{ion} + \text{diss}}{n_g w_{ion} + \text{diss}} + v_g \quad (2)$$

- (3) shielding function

$$\delta_g = 1 - \frac{1}{K T_e} \int_{r_g}^{r_p} L[E(r)] n(r) dr \quad (3)$$

- (4) pellet equation of motion

$$\frac{ds}{dt} = u_p \quad (4)$$

The definitions of the symbols are: n_p = solid pellet density; r_p = radius of evaporation deflagration wave; r_g = radius of gas-plasma boundary; $\tilde{\Gamma}_e = 1/\sqrt{2\pi} \tilde{n}_e K T_e \tilde{v}_{the}$ = electron energy flux. The symbol $\tilde{x}$ refers to quantities which are functions, $\tilde{x} = \tilde{x}[s(t), t]$, of time and the position of the pellet in an inhomogeneous plasma.

$n(E)$ = efficiencies for energy loss by the denoted processes

n_g = density of H_2 in the gas flow field at r_g , the gas-plasma boundary

$n(r)$ = density of H_2 in the gas flow field between r_p and r_g

w = average energy/atom required to effect the denoted processes

Δw = average thermal energy/atom convected away by the sonic flow at the pellet surface
approximately given by the transcendental relation (where $\Delta w \ll w_{eo}$)

$$\frac{\Gamma_e}{w_{eo}} = 4.97 \times 10^5 \frac{1}{\Gamma \Delta w} \exp \left[\frac{-0.0421}{\Delta w} + 6.075 + 24.46 \Delta w + 3542 (\Delta w)^2 \right]^9 \quad (5)$$

v_g = average radial velocity of H_2 in the gas flow field at r_g

δ_g = shielding function

$L(E)$ = total electron energy loss function

s = pellet position

u_p = pellet velocity

The initial conditions for this set of equations are:

$$1. \quad r_p(t=0) = r_g(t=0) = r_{po}$$

$$2. \quad \delta_g(t=0) = 1$$

$$3. \quad s(t=0) = 0$$

$$4. \quad ds/dt(t=0) = u_{po}$$

A method is now presented for the self-consistent calculation of the gas flow field and its shielding properties relative to the ablating pellet. The gas leaving the pellet surface is described by a shifted Maxwellian distribution. In keeping with the Chapman-Jouget condition, we take the flow to be locally sonic at the downstream side of the deflagration transition zone, in the frame of reference of this zone. Neglecting collisions, we solve the forceless Liouville equation

$$\left(\frac{\partial}{\partial t} + v \nabla \right) f = f_{MS} \frac{\delta(r - r_p)}{4\pi r_p^2} \delta(t - t_p) \quad , \quad (6)$$

where

$$f_{MS} = \frac{1}{\beta_0^{3/2} \pi^{3/2}} e^{-[v_{\parallel} - c_s]^2 + v_{\perp}^2 / \beta_0} \quad . \quad (7)$$

The coordinate system is shown in Fig. 7.1. Note that in the following, full advantage has been taken of the spherical symmetry in the problem.

The solution to Eq. (6) is given by

$$f = \frac{f_{MS}}{4\pi r_p^2} \theta(t - t_p) \delta(g) \quad , \quad (8)$$

where g is the relation

$$g = c(v_{\perp}^2 t^2 + (|v_{\parallel}| t + r_0)^2 - r^2) \quad , \quad (9)$$

and c is a constant we introduce for normalizing purposes.

We are interested in generating a kernel which describes conditions in the flow field at (r, t) , given conditions at (r_p, t_p) . First we calculate the normalized density by integrating Eq. (8) over the appropriate region in velocity space:

$$\bar{n}(r, t) = \frac{r^2}{r_p^2} \frac{n(r, t)}{n(r_p, t_p)} = \int \frac{dv_{\parallel}}{\beta_0^{3/2}} \frac{2\pi v_{\perp}}{\pi^{3/2}} \frac{dv_{\perp}}{\pi^{3/2}} e^{-(v_{\parallel} - c_s)^2 + v_{\perp}^2 / \beta_0} \delta(g) \quad (10)$$

For initially outward-moving particles, Eq. (10) becomes

ORNL-DWG 76-10858

Fig. 7.1. Pellet geometry.

$$\bar{n}_o(r,t) = \int_0^{r-r_p} \frac{2\pi dv_{||}}{\beta_o^{3/2} \pi^{3/2}} e^{-(v_{||} - c_s)^2/\beta_o} \int dv_{\perp} v_{\perp} e^{-v_{\perp}^2/\beta_o} \times \frac{\delta(v_{\perp} - v_{\perp}^o)}{|\frac{\partial g}{\partial v_{\perp}}|_{v_{\perp} = v_{\perp}^o}}, \quad (11)$$

where $v_{\perp}^o$ is a function denoting the appropriate zeros of g .

Likewise, for initially inward-moving particles, assuming specular reflection at the pellet surface,

$$\bar{n}_i(r,t) = \int_{\frac{r_p - r}{t - t_p}}^0 \frac{2\pi dv_{||}}{\beta_o^{3/2} \pi^{3/2}} e^{-(v_{||} - c_s)^2/\beta_o} \int dv_{\perp} v_{\perp} e^{-v_{\perp}^2/\beta_o} \times \frac{\delta(v_{\perp} - v_{\perp}^o)}{|\frac{\partial g}{\partial v_{\perp}}|_{v_{\perp} = v_{\perp}^o}}. \quad (12)$$

Having performed the integration, summed Eqs. (11) and (12), and normalized the resulting density kernel to unity at $t = t_p$ and $r = r_p$, we find

$$\bar{n}(r,t;r_p,t_p) = \frac{M}{TR^2 \pi^{1/2}} e^{\frac{M^2}{T^2} (R^2 - 1 + T^2)} \left\{ \frac{e^{2(T+1)(R-1)\frac{M^2}{T^2}} - 1}{(T+1)} + \alpha \frac{1 - e^{2(T-1)(R-1)\frac{M^2}{T^2}}}{(T-1)} \right\}, \quad (13)$$

where α is a reflection coefficient assumed for the present to be unity. The variables have been normalized as follows:

$$\begin{aligned} R &= \frac{r}{r_p} \\ T &= \frac{c_s(1 - \frac{1}{\epsilon})(t - t_p)}{r_p} \\ M^2 &= c_s^2(1 - \frac{1}{\epsilon})^2/\beta_o \end{aligned} \quad (14)$$

Here ϵ is the expansion ratio at r_p^+ , the downstream side of the inner deflagration wave transition zone. The expansion ratio is the ratio of the solid density to the density at r_p^+ .¹⁰

With Eq. (13), we may write the density in the gas flow field directly as the integral

$$N(r,t) = \int_{t-\tau}^t \frac{n_p}{\epsilon(t'_p)} dT K_n[r'_p(t'_p), t'_p; r, t], \quad (15)$$

where $t' = t - \tau$ is taken as the time at which the ablation process begins (so far as an observer at r, t is concerned), and K_n is related to $\bar{n}$ in Eq. (13) with the values defined in Eq. (14) replaced by functions to be determined by the ablation dynamical equations. The solid density of the pellet is represented by n_p .

The kernel in Eq. (15) is given by

$$\begin{aligned} K_n[r'_p(t'), t'_p; r, t] = \\ \frac{Me^{-\frac{M^2}{T^2}(R^2-1+T^2)}}{T\pi^{1/2}R^2} \left\{ \frac{e^{\frac{2(T+1)(R-1)M^2}{T^2}} - 1}{(T+1)} + \frac{1 - e^{\frac{2(T-1)(R-1)M^2}{T^2}}}{(T-1)} \right\}. \end{aligned} \quad (16)$$

In particular, we have the gas density at the gas-plasma interface as

$$n(r_g(t), t) = \int_{t-\tau}^t \frac{n_p}{\epsilon(t'_p)} dT K_n(r'_p(t'_p), t'_p; r_g(t), t) \quad (17)$$

Using an analogous technique we find the flux kernel, K_ϕ , which relates to the average particle flux in the gas flow field, to be

$$\begin{aligned} K_\phi(r'_p(t'_p), t'_p, r'_g(t'); r, t) = & \\ & \frac{-\frac{M^2}{T^2}(T^2+R^2-1)}{R^2 T \pi^{1/2}} \left[\frac{\left(v - c'_s \frac{T}{(T+1)2M^2} \right) e^{\frac{M^2}{T^2} 2(T+1) \frac{v}{c'_s}}}{(T+1)} + \frac{c'_s T}{2M^2(T+1)^2} \right. \\ & \left. + \alpha \left\{ \frac{\left(v + c'_s \frac{T}{(T-1)2M^2} \right) e^{-\frac{M^2}{T^2} 2(T-1) \frac{v}{c'_s}}}{(T-1)} - \frac{c'_s T}{2M^2(T-1)^2} \right\} \right], \quad (18) \end{aligned}$$

where

$$\frac{v}{c'_s} = (R-1)/T$$

and

$$c'_s = c_s \left(1 - \frac{1}{\epsilon(t'_p)} \right)$$

Dimensionless parameters are defined as before.

We may now rewrite the original ablation equations describing the model with the use of these relations. We have

$$\frac{d}{dt} r_p = - \frac{\tilde{r}_e}{n_p} \frac{n_{ev}}{w_{ev}} \delta_g \quad (19)$$

$$\frac{d}{dt} r_g = - \frac{\tilde{r}_e n_{ion+diss}}{n(r_g, t) w_{ion+diss}} + \bar{v}(r_g, t) \quad (20)$$

$$\delta_g[r_p(t), t] = 1 - \frac{1}{K \tilde{T}_e} \int_{r_g(t)}^{r_p(t)} L[E(r)] n(r, t) dr \quad (21)$$

$$\frac{ds}{dt} = u_p, \quad (22)$$

where $\tilde{x}$ refers to quantities which are functions of $s(t)$ and t because of the pellet's motion through an inhomogeneous plasma

r_p = pellet radius

r_g = gas-plasma interface radius

δ_g = shielding function

w = average energy involved in the atomic physics process denoted

$L(E)$ = electron energy degradation function in the continuous slowing-down approximation

s = pellet position

u_p = pellet velocity .

The density in the gas flow field appearing in Eq. (21) and Eq. (19) [through Eq. (21)] is given by

$$n(r,t) = \int_{t-\tau}^t \frac{n_p}{\epsilon(t'_p)} dT K_n[r'_p(t'_p), t'_p; r, t] , \quad (23)$$

where ϵ is the gas dynamic expansion ratio defined by

$$\epsilon = \frac{n_p T_s}{p_v(T_s)} \quad (24)$$

and

$$p_v(T_s) = \exp(a/T_s + b + cT_s + dT_s^2) \quad (25)$$

is the vapor pressure of hydrogen constrained to the critical point. The variable T_s is the temperature of the surface of the pellet given by the relation which balances power flux across the surface

$$3\phi_{eo} \tilde{T}_{eS}/\sqrt{2\pi} = \frac{n_p}{\epsilon} (W_{ev} + \frac{3}{2} (T_s - T_0) + \frac{\gamma}{2} T_s M_0^2) C_{so} M_0 . \quad (26)$$

The density in the gas flow field at the gas-plasma boundary appearing in Eq. (20) is given by

$$n(r_g, t) = \int_{t-\tau}^t \frac{n_p}{\epsilon(t'_p)} dT K_n(r'_p(t'_p), t'_p, r'_g(t'_p); r_g, t) . \quad (27)$$

Finally, the gas flow field velocity at the gas-plasma boundary appearing in Eq. (20) is given by

$$\bar{v}(r_g, t) = \frac{\int_{t-\tau}^t \frac{n_p}{\epsilon(t'_p)} dT K_\phi(r'_p(t'_p), t'_p; r_g, t)}{n(v_g, t)} , \quad (28)$$

with the initial conditions

$$\begin{aligned} s(0) &= 0 \\ r_p(0) &= r_{p0} \\ r_g(0) &= r_{p0} \\ K\tilde{T}_e &= K\tilde{T}_e[s(t), t] \\ \tilde{n}_e &= n_e[s(t), t] . \end{aligned} \quad (29)$$

The initial value problem is thus determined. Numerical solutions of the coupled equations (19), (20), (21), (22), (23), (26), (27) have been obtained.

The solutions of these equations indicate:

- (1) Our original assumption, that the gas flow field is unchanged by the local heat deposition in different plasma regimes, may be insufficient for reactor-grade plasma but is quite good for those plasmas in which experimental results have been obtained.
- (2) The time-dependence of the ablation process including the deposition of fuel and the spectral luminosity of the ablatant cloud can be unsteady.
- (3) The degree of ionization of the flow beyond the gas-plasma interface is incomplete. Statement (3) provides necessary information for the description of the effects of charge exchange in the transport of those pellet-produced neutrals which escape to the external plasma.

Implicit in the preceding analysis are some basic assumptions requiring explanation.

A. The electron energy flux which drives the evaporation, thus determining the shielding function, is calculated in such a way as to neglect elastic collisions of electrons with gas molecules, and self-consistent electric potential. The effect of this in the model is that electrons stream through the gas rather than diffusing, and the scalar flux of electrons is spatially constant. Inclusion of the elastic scattering effect enhances the flux of electrons at the pellet surface, but depresses the energy of the electrons there. This latter effect is the result of the longer effective path length traveled by electrons diffusing toward the pellet surface through the ablatant. The effect of the potential is to suppress the electron flux near the pellet surface. Efforts are now being made to find the importance of these counter-balancing effects in determining the electron power flux at the pellet surface.

B. The calculation made also assumes that the mean free path of molecules in the gas against self-collisions is of the order of the gradient scale length of the density, i.e., free-jet expansion. This is warranted for low ablation rates (rarefied ablatant). In EPR this assumption is at best questionable. Calculations show the difference to be relatively small if one uses a fluid approach that assumes the opposite, i.e., that the mean free path for molecular collisions is short relative to the density-gradient scale-length, provided the gas flow is not strongly heated. Strong heating tends to choke the gas flow and alter the profiles of molecular density. Qualitatively, the density is enhanced near the pellet and decreased at larger radii.¹¹

C. The charging of the pellet by the electron flux is a problem which has not yet been addressed in this research. One can argue somewhat convincingly, however, that secondary electrons, produced by the inelastic collisions of reactor electrons with electrons bound in gas molecules and atoms, can easily short out potentials which appear between the pellet and the plasma. The secondaries are the cool electrons which occur when hot electrons collide with bound electrons, giving up a large proportion of their energy to the bound electron, though retaining enough energy to resist capture. The net kinetic energy lost from the electrons is the binding energy of the target electron, which later undergoes inelastic collisions if its kinetic energy is large enough. Since the secondaries are typically too cool to undergo many inelastic collisions, they are quite mobile; the inward-bound hot electron flux is balanced by a cold outward flux of secondaries driven by a potential which is small compared with the reactor-electron thermal energy. This is not true for regions of the ablatant in which ionization does not take place. Thus a potential persists near the pellet surface relative to the external plasma which is large enough to turn around those incident electrons whose energies are smaller than that necessary to cause ionization.

There is a significant difference between the hydrogen pellet and a metal probe inserted into a plasma. Solid hydrogen is an insulator and has a negative surface work function. This means that while electrons incident on a metal probe are likely to find conduction bands in the metal lattice in which to reside, electrons incident on solid H_2 are less likely to stick. Those electrons which do stick to the pellet reside very near the surface and are subsequently liberated into the gas flow by the ablation process.¹² They quickly join the outward secondary flux and help to maintain the surface potential near the ionization threshold.

D. The effects of the magnetic field in which the pellet-ablatant ensemble is immersed have been largely neglected in the calculations. Three effects of such a magnetic field have been examined.

1. In the case that the pitch-angle scattering collision frequency for electrons is much smaller than the gyrofrequency, the effects of the magnetic field on the ablation rate become important, which occurs if the ablatant density of hydrogen molecules is $\leq 10^{16}/\text{cc}$. This condition is fulfilled for regions of the reactor plasma where the electron power flux is quite small, e.g., the plasma edge and

the periphery of the ablatant cloud. For all the calculations presented, however, these effects are neglected.

2. The motion of the pellet-ablatant ensemble across field lines causes an electric field to be seen in the plasma from the frame of reference of the pellets. The field is in the direction mutually perpendicular to both the direction of motion and the magnetic field in the ablatant.⁵ A current due to this electric field, which flows in the external medium, the plasma, causes a drag force on the ablatant in the peripheral regions of the pellet-ablation ensemble, tending to "scrape off" the ablatant and deposit it on the flux surface the pellet is crossing.¹³ The conductivity of the plasma in this peripheral region determines the magnitude of the effect. An essentially local deposition of ablatant beyond the gas-plasma boundary referred to in an earlier section has been assumed. Further study is needed of the changes in the conductivity of the plasma in this region due to ablatant-induced effects (cooling and magnetic field refraction).

3. The expansion of the ionized ablatant plasma into the reactor plasma, in the presence of a magnetic field, drives a diamagnetic current in the ablatant itself.^{4,15} For high- β ablatants this current can refract the external magnetic field around the region in which the pellet ablates. In all cases of EPR interest this is a weak effect. The power fluxes necessary for this effect to become large are beyond the reference design plasma parameters by a comfortable margin. The giant laser-pellet induced plasmas fit this category, and a wealth of interesting physics questions pertain to that problem.¹⁴

E. As mentioned earlier, the effects of alpha fusion products on ablation have been neglected. Though no justification can be made for this, analysis shows the power flux of suprathermal and 3.5-MeV alpha particles to be smaller than that of the electrons.¹⁶ The effects of runaway electrons have also been neglected.

The ablation behavior of a pellet (1) which is subjected to plasma conditions that change with time or (2) which is moving through an inhomogeneous plasma is determined by the relative rates at which the plasma conditions change and the ablatant shield changes. In the case of the pellet which moves slowly through gradual inhomogeneities in plasma conditions, the ablation shield adjusts itself so as to allow a somewhat slowly varying electron energy flux to impinge on the pellet surface. However, in the case of the pellet which encounters changes of electron energy flux occurring on a time scale that is faster than the time scale over which the shielding layer is set up, large excursions in ablation rate occur, followed by overshielding and shutdown of the ablation process.

In general, the geometric considerations are such that larger pellets are more likely to show the oscillatory behavior for faster changes in plasma conditions, since the time necessary to set up the shielding layer scales as pellet radius. Smaller pellets, on the other hand, have more tenuous shields (due to the $\sim 1/r^2$ shield density dependence) and tend to ablate faster and more steadily.

7.3.2 Results of the Pellet-Plasma Interaction Model

Figure 7.2 depicts the steady ablation of a 2-mm pellet encountering "gradual changes" in plasma conditions (electron density, temperature). The curves labeled $t_1 \dots t_9$ and the two graphs show the density (dotted curves) and electron temperature (solid curves) in the gas flow field at successive times. The large density near the pellet surface at t_1, t_2 is due to the initial transient behavior of the bare pellet as it enters the reactor plasma edge. At later times one sees the initial density "lump" relax and the entire shield density rise gradually, retaining its form. This is due to the rise in the reactor electron temperature and density. The ends of the curves (to the right) indicate the positions of the sharp gas-plasma boundary at which the power flux balance is struck (see Sect. 7.3.1) and the respective value of the reactor electron temperature and ablatant density there. The

Fig. 7.2. The steady ablation of a pellet encountering gradual changes in plasma conditions.

drop in incident electron temperature through the ablatant is due to the slowing down of primary electrons by inelastic collisions in the ablatant evaporate. The motion of the inner boundary is imperceptible on this scale. Approximately 10 μ sec into the interaction, t_0 occurs.

Figure 7.3a,b depicts the ablation of a pellet whose shield cannot keep up with the changes in plasma conditions. Here a finer time mesh is shown to indicate (Fig. 7.3a) the evolution of the shield density and (Fig. 7.3b) the electron temperature as the electrons are chilled by the ablatant. A full period in the oscillation is depicted by the figure. Notice that at t_1 the shield is at its lowest strength. That is, the number of molecules between the external plasma and the pellet surface is minimal. The result shortly thereafter is (1) the rapid increase in the temperature of incident electrons at the pellet surface and (2) an ablation excursion. The relaxation of the shield density which follows (approaching t_2) leaves the pellet undershielded to the reactor electron energy flux at t_2 , which has increased since t_1 due to the pellet motion through inhomogeneous plasma. A second excursion follows. The result is that the ablation rate is clearly above that of a pellet which is neither over- nor under-shielded throughout. On this time scale, the inner boundary motion is imperceptible.

On the basis of the results of calculations such as these, it is clear that the relationships between pellet size, injection velocity, and penetration depth are not simple linear or monotonic relationships. For example, having chosen a given size pellet for injection, one cannot expect the penetration depth to increase with velocity beyond the velocity which allows intermittent ablation to occur. Indeed, one can expect the penetration depth to reach a maximum for some optimal injection speed. On the other hand, having chosen a velocity one might expect several optimal sizes of pellets which will allow a maximal penetration depth.

With this insight into the pellet-plasma problem, we can be optimistic on three points:

1. We can expect the well-shielded pellet to shrink more rapidly as it becomes smaller, thus allowing larger specific depositions at the end of its lifetime.
2. We can choose optimal injection velocities well below the technological limits set by feasibility.
3. Within the range of reasonable technological development of future pellet-accelerating devices, we can choose optimally-sized pellets which remain intact in the acceleration process.

It seems indeed probable that 1-3-mm size pellets injected at $\sim 2-5 \times 10^3$ m/sec will be adequate for the EPR design.

The efforts presently under way will further refine our ideas regarding the ablation process. Altering the model from a sharp boundary outer deflagration-ionization process to a heated spherical expansion is now under way. The effects of alpha particles on the ablation process will be addressed.

7.3.3 Pellet Acceleration Techniques¹⁷

Mechanical Rotating Devices. The use of a rotating arm device resembling a high-speed centrifuge appears feasible and practical for mechanical acceleration of large pellets (on the order of 1 mm) to speeds of $1-6 \times 10^3$ m/sec. A device of this kind would have certain advantages, such as the ability to accelerate a variety of different pellet sizes through a wide range of velocities without the need to modify the physical characteristics of the device. Rotor-tip speeds of 1×10^3 m/sec are attained using high-strength aluminum as the material of construction. Velocities as high as 6×10^3 m/sec may be attainable by using high-performance composite materials in combination with a high-strength metallic hub. Although these speeds may not be adequate for EPR, they will allow us to develop scaling laws for pellet ablation rates and lifetimes.

NOTE: THE LARGE ELECTRON TEMPERATURES AT t_1 AND t_2 ARE FOLLOWED QUICKLY BY A LARGE DENSITY EXCURSION WHICH RELAXES IN $\sim 1 \mu\text{sec}$.

Fig. 7.3. The dynamic ablation of a pellet whose shield cannot keep up with the changes in plasma conditions.

Conceptually, an apparatus suitable for accelerating pellets would consist of a rotor with three equally-spaced arms designed so as to maximize rotor-tip speed. Channels machined into the rotor arms would serve to guide pellets as they move by centrifugal forces from the hub of the rotor to the tip of the arms, where they leave at a speed approximately 40% greater than the rotor-tip speed. Pellets would enter the guide channels after being clipped from a solid extruder by the rotating motion of the arms. If temporal changes in channel friction are not large and if they also enter the device with a time dispersion small compared to $1/\omega$ rotations, then all pellets would leave the assembly at a regular angular position with negligible dispersion. A device of this kind operating at 1000 Hz and having a rotor diameter ≈ 1 m will provide velocities above 10^3 m/sec. Pellets of the order of 1 mm diameter are envisioned, and the resulting acceleration will not exceed the estimated maximum allowable value of 4×10^7 m/sec² (see Sect. 7.3.4 of this report). High-speed centrifuges with characteristics similar to this are routinely operated in ORNL's Molecular Anatomy Program. These units are driven at high rotational speeds by commercial turbine drive systems powered with ordinary building supply air. The rotor is encased in protective armor plate and the rotor chamber is kept partially evacuated to reduce parasitic drag.

The solid extrusion apparatus required to feed the centrifuge at the needed rate also seems practicable, as described below in the section entitled "Solid Extrusion Devices."

Performance Evaluation of the Light Gas Gun. A light gas gun operates in the principle of a high-pressure gas exerting an accelerating force on one side of a projectile that is constrained laterally by the walls of a launch tube (gun barrel). The gas and projectile accelerate in the launch tube until the gas velocity reaches the "choked" or sonic value. To maximize performance, operating devices utilize gases of low molecular weight at elevated pressures where the sound speed exceeds the ideal gas value.

The departure from ideality is usually obtained by compressing a gas charge adiabatically with an explosively driven piston. Projectile velocities approaching 10^4 m/sec have been reported for such a device, which utilized gaseous hydrogen as the propellant (Ref. 2).

An evaluation of the performance of a light gas gun was made by approximating the properties of high-density gaseous hydrogen by a van der Waals fluid model. In reduced form, the van der Waals equation of state can be written as

$$P_r = \frac{8T_r}{3V_r - 1} - \frac{3}{V_r^2} \quad , \quad (30)$$

where the pressure, temperature, and specific volume are reduced by their respective values at the thermodynamic critical point. In terms of these reduced properties, the local speed of sound for a real gas is calculated from the following expression:

$$a_s = \sqrt{Z_c \gamma R T_c V_r^2 \left(\frac{\partial P_r}{\partial V_r} \right)_{T_r}} \quad , \text{ m/sec} \quad . \quad (31)$$

Here Z_c is the critical compressibility factor (3/8), R is the gas constant, and T_c the critical temperature (38 K). When values of the reduced volume approach 1/3 and/or the reduced temperature is high, the isothermal compressibility ($1/V_r \partial P_r / \partial V_r$) becomes large and results in high values of the sound speed. The specific heat γ is a complicated function of pressure and temperature, but its value varies little over a broad range in conditions ($\gamma \approx 1.4$). The gas constant R is inversely proportional to molecular weight and this is why low molecular weight gases are preferred.

For a reversible adiabatic compression process between state points 1 and 2, the resulting temperature ratio is given for this model by the simple result

$$\frac{T_{r2}}{T_{r1}} = \left(\frac{3V_{r1}-1}{3V_{r2}-1} \right)^{0.4} \quad (32)$$

A knowledge of the initial conditions V_{r1} , T_{r1} , and the device compression ratio, is sufficient information to calculate the state point T_{r2} , V_{r2} . The sound speed follows directly from Eqs. (30) and (31).

The gas gun described in Ref. 2 typically operates with an initial charge of hydrogen at ambient temperature and a pressure of 6.7 bars ($P_{r1} \approx 0.51$, $T_{r1} \approx 7.89$, $V_{r1} \approx 41.2$), and develops a compression ratio of 60 ($V_{r2} = 0.69$). The calculated values of the final chamber temperature and pressure are 2000 K and 5000 bars respectively. The resulting speed of sound computed from Eq. (31) is 6300 m/sec for H_2 . If deuterium or tritium were used, this value would be reduced due to the higher values of molecular weight.

Solid Extrusion Devices. Solid hydrogen extrusion devices are presently being considered as methods for production of pellets larger than 100 μm . A number of such devices have been fabricated elsewhere for the purpose of studying materials properties at low temperatures and of providing targets for laser plasma interaction studies. These devices share in common a piston and cylinder assembly and some kind of extrusion die. Basically, low-pressure hydrogen gas is condensed and solidified in the cylinder which is cooled by liquid helium. Pressure applied to the frozen hydrogen billet by the piston and drive assembly causes the hydrogen to flow plastically at the neck of the die, forming a continuous solid filament.

Inasmuch as the yield strength of solidified hydrogen is expected to be small, the pressure required to initiate the plastic deformation will not have to be particularly high. This can be calculated from a knowledge of the billet and die dimensions and estimated sample yield strength. If the billet and die radii are given respectively by R_0 and R_i , then the required extrusion pressure will be given by the expression

$$P = S_y \ln \left(\frac{R_0}{R_i} \right), \quad (33)$$

where S_y is the yield strength of the material to be extruded. For hydrogen the yield strength is expected to be of the order of 20 bars. Then to extrude a 1-mm wire from a 10-mm sample, a pressure of about 100 bars is required. This is easily attainable even with a small electric motor drive system. A cylinder with a length of 10 cm would have sufficient volume to produce a 10-m length of 1-mm wire. Extrusion rates of 1 m/sec should be attainable.

7.3.4 Mechanical Properties and Acceleration Limit of Solid Hydrogen

When the pellet experiences an acceleration, opposing inertial forces result which give rise to a longitudinal compression of the pellet. The compression is accompanied by a proportionate amount of lateral tensile strain in both principal transverse directions. The constant of proportionality, Poisson's ratio, typically has values in the range 0.25-0.30. As long as the amount of lateral strain does not exceed the elastic limit of the material, the solid structure will remain intact. Common materials can withstand at most a few tenths of a per cent of tensile strain before plastic deformation results. Inasmuch as solid hydrogen is relatively compressible, it will exceed its elastic limit under small compressive loads. An estimate of this yield strength can be made from a knowledge of the average elastic modulus E , which is related to the speed of sound c and mass density ρ by the relationship

$$E = \rho \left(\frac{\partial P}{\partial \rho} \right)_{\text{constant entropy}} = c^2 \rho \quad (34)$$

For hydrogen, $c \approx 2 \times 10^3$ m/sec, $\rho = 86$ Kg/m³, and hence $E = 3400$ bars. A similar result is obtained by considering the nature of the elastic restoring forces between crystal lattice sites.

By limiting the maximum elastic strain to the order of 1%, the maximum allowable stress that the solid can accommodate before yielding becomes

$$\Delta P_{\max} = E \frac{\Delta \rho}{\rho_{\max}} \approx 0.01 E = 34 \text{ bars} \quad (35)$$

Such a low yield strength is not surprising, since solid hydrogen is often used as the hydrostatic pressure-transmitting medium in measurements of low-temperature physical properties.

If the pressure given above were distributed uniformly over one side of a cube of dimension ℓ , the acceleration would be given by

$$A_{\max} = \frac{\Delta P_{\max}}{\ell \rho} \quad (36)$$

For a 1-mm pellet, we then expect maximum tolerable accelerations on the order of 4×10^7 m/sec². To attain a terminal velocity of 10^4 m/sec, a pellet of this size would require an accelerating path of 1 m. Bigger pellets and/or higher velocities would need longer acceleration paths. The same result applies if the tensile strength of hydrogen should prove to be lower.

A similar argument can be made when the pellet experiences larger impulsive loads which are applied over time periods shorter than the time required for the elastic stress wave to propagate through the pellet (10^{-6} sec for a 1-mm pellet). In this case, however, that amount of applied pressure in excess of the elastic limit (~ 34 bars) is propagated as a slowly-moving plastic wave that deforms the pellet as it propagates, dissipating energy as heat. The pellet can probably remain intact if such large waves do not propagate more than a fractional distance into the pellet bulk. For 1-mm pellets, this would require that large pressure excursions be limited to 10^{-6} sec or less in duration.

Acknowledgments

This work has benefited from discussions with C. A. Foster and S. Milora.

REFERENCES

1. M. T. Robinson, "The Reflection of Low Energy Hydrogen Atoms from Metal Surfaces," Third National Conference on Atomic Collisions With Solids, Kiev, U.S.S.R. (October 1974).
2. O. S. Oen and M. T. Robinson (ORNL), private communication, August 1975.
3. H. Verbeek and W. Eckstein, "Reflection of Hydrogen from Stainless Steel and Nb," Third International Conference on Surface Effects in Fusion Devices, San Francisco (February 1976).
4. D. J. Rose, "On the Fusion Injection Problem," Culham Technical Memo #82, UKAEA Culham Laboratory, Abingdon, England (1968).
5. L. Spitzer, Jr. et al., *Problems of the Stellarator as a Useful Power Source*, NYO-6047 (PM-S-14), Project Matterhorn, Princeton University (August 1, 1954).
6. S. L. Gralnick, "Solid Deuterium Evaporation in a Fusion Plasma," *Nucl. Fusion* 703, 1973.

7. W. T. Miles, R. Thompson, and A. E. S. Green, "Electron Impact Cross Sections and Energy Deposition in Molecular H_2 ," *J. Appl. Phys.* 43(2): 678 (1972).
8. M. G. Heaps and A. E. S. Green, "Energy Spreading and Angular Distribution of a Beam of Electrons in Molecular Hydrogen," *J. Appl. Phys.* 46(11): 4718 (1975).
9. R. Prydz, K. D. Timmerhaus, and R. B. Steward, "The Thermodynamic Properties of Deuterium," *Adv. in Cryogenic Eng.* Vol. 18 (1973).
10. S. L. Gralnick, *Solid Deuterium Evaporation in a Fusion Plasma*, Ph.D. Thesis, Columbia University (1972).
11. C. A. Foster, Ph.D. Thesis, University of Illinois (1976).
12. R. Clampitt, "Ablation of Solid Hydrogen by Electrons and Ions," CLM-P400, UKAEA Culham Laboratory, Abingdon, England (1974).
13. W. H. Bostick, *Proc. Symposium on Electromagnetic and Fluid Dynamics of Gaseous Plasma*, p. 211 (1961).
14. D. K. Bhadra, "Expansion of a Resistive Plasmoid in a Magnetic Field," *Phys. Fluids* 11: 234 (1967).
15. C. T. Chang, "The Magnetic Shielding Effect of a Refueling Pellet," *Nucl. Fusion* 15: 595 (1975).
16. D. Vasilow, private communication (to be published).
17. S. L. Milora, C. A. Foster, and L. D. Stewart, *Plasma Heating Quarterly Report*, ORNL (Jan. 6, 1976).

8. NEUTRAL BEAM INJECTION SYSTEM

8.1 Definition and Requirements

The function of the neutral beam injectors is to contribute to start-up heating of the plasma and, if ignition is not reached, to provide supplementary heating for achieving useful power levels in a beam-driven EPR. The injectors consist of: (1) a deuterium gas source which is supplied from the tritium process and vacuum pump system; (2) an ion accelerator including a deuterium plasma source, extraction electrodes, and accelerating electrodes; (3) a neutralizer; (4) an ion magnetic deflector; (5) a drift tube; and (6) controls. An ion energy recovery system is an optional feature for reducing power needs. The injectors are supplied with power by the injector power supply and are evacuated by the injector vacuum system.

Neutral injection system energy and power requirements are determined by properties of the plasma. On the basis of the analysis reported in Sect. 2, the design goal for neutral beams delivered to the EPR plasma was set at 50 MW at 200 keV in convenient modules. The duration of beam injection is governed by (a) the volt-second OH limit if the EPR is capable of purging helium and refueling, (b) the impurity-buildup limit, and (c) whether or not the plasma ignites.

8.2. Consideration of Possibilities

Figure 8.1 illustrates the basic injection system and identifies the various currents and efficiencies.

The ion accelerator involves two efficiencies: η_1 represents the fraction of the extracted beam which is the useful ion species [$D^+/(D^+ + D_2^+)$ for positive ions, $D^-/(D^- + e^-)$ for direct extraction of negative ions, or D^- production efficiency for an alkali-metal vapor charge-exchange of low-energy D^+ to D^- coupled with post-acceleration]; η_2 represents D^+ or D^- beam losses to accelerating grids and neutralizer as well as further geometric losses downstream.

The neutralizer is either a D^+ charge exchange neutralizing cell or a D^- stripping cell. The associated efficiency, η_3 , represents the fraction actually used of an equilibrium cell of conversion efficiency $F_{0\infty}$.

The deflection magnet removes the remaining charged particles from the beam that leaves the neutralizer. It may either dump them or focus them into the ion-energy recoverer.

The drift tube transports the D^0 beam from the deflector to the EPR plasma. The associated efficiency, η_4 , represents primarily ionization losses from collision with background gas in the drift tube.

The optional ion-energy recoverer is a direct conversion system with efficiency η_5 for converting the ion beam to dc power, and efficiency η_6 for dc-dc conversion and feedback into the ion accelerator.

Associated electric or equivalent neutral beam currents are:

I_A — the electric current into the ion accelerator first grid

I_0 — the neutral 200 keV D^0 current into the EPR plasma

I_R — the recoverable beam current

I_N — the net-electric acceleration current supplied externally.

These currents are given by:

$$I_0 = I_A \eta_1 \eta_2 \eta_3 F_{0\infty} \eta_4$$

$$I_R = I_N \eta_1 \eta_2 (1 - \eta_3 F_{0\infty}) \eta_5 \eta_6$$

$$I_N = I_A - I_R$$

The associated power is simply taken as $P = I \cdot V$; the ion production power is assumed to be negligible compared to the acceleration power.

Fig. 8.1. Model of neutral beam injector systems.

The possible neutral beam injection schemes are:

- (1) Positive ion acceleration, dumping of residual ion fraction after neutralization.
- (2) Positive ion acceleration, energy recovery of residual ion fraction after neutralization.
- (3) Negative ion acceleration (direct or indirect production), dumping of residual ion fraction after neutralization.
- (4) Negative ion acceleration (direct or indirect production), energy recovery of ion fraction after neutralization.

The preferred scheme depends on expected EPR performance and progress in the beam program. The simpler schemes require less beam development, but their higher net power requirement suggests their use for an igniting EPR with an energy storage system supplying their power. The more complex schemes, requiring more beam development, require lower net power, suggesting their use for a beam-driven EPR.

Designs representative of each of these cases are summarized in Table 8.1. Advantages and disadvantages are summarized in Table 8.2.

Table 8.1. Representative design values for four possible types of EPR neutral beam injectors

	Case 1	Case 2	Case 3	Case 4
Type	D^+	D^+ , recovery	D^-	D^- , recovery
I_0 , A	250	250	250	250
η_1	0.95	0.95	0.95	0.95
η_2	0.90	0.90	0.90	0.90
η_3	0.90	0.90	0.90	0.90
η_4	0.97	0.97	0.97	0.97
η_5	0.00	0.85	0.00	0.85
η_6	0.00	0.98	0.00	0.98
$F_{0\infty}$	0.20	0.20	0.67	0.67
I_A , A	1667	1667	500	500
P_A , MW	333	333	100	100
Number of sources	18	18	6	6
I_A /source	93	93	83	83
Beam lines	6	6	3	3
Sources/line	3	3	2	2
I_R , A	0	967	0	140
I_N , A	1667	700	500	360
P_N , MW	333	140	100	72
P_0 , MW	50	50	50	50
Efficiency, P_0/P_N	0.15	0.36	0.50	0.69

Table 8.2. Advantages and disadvantages of the four possible types of EPR neutral beam injectors

<u>Type</u>	<u>Advantages</u>	<u>Disadvantages</u>
D^+	Straightforward extension of TFTR technology.	Low electrical efficiency; large number of beam lines and injectors required (costs).
D^+ plus recovery	TFTR basic source and beam line technology; moderate electrical efficiency.	Large number of beam lines and injectors (costs); stored energy required until recovery effective; efficient high power recovery must still be developed.
D^-	Moderate electrical efficiency; fewer beam lines and injectors.	Existing experimental D^- sources have very low gas efficiency; high-power D^- sources must still be developed.
D^- plus recovery	High electrical efficiency; fewer beam lines and injectors.	Existing experimental D^- sources have very low gas efficiency; high-power D^- sources and efficient high-power recovery of D^- must still be developed.

8.3 Recommendations for Reference Design

The recommended reference design system is a negative ion source, with no recovery of charged-particle fraction energy (Case 3 in Table 8.1). This design, shown schematically in Fig. 8.2, couples the beam-handling experience gained from the 150-keV TFTR beams with the development of a suitable negative ion source to give an efficient overall system. The substantial reduction in the number of beam lines, injectors, and primary electric power of the negative ion system over the positive ion systems results in a significant reduction in EPR costs. The additional benefits gained from direct ion-energy recovery do not yet warrant the added complexity of the recovery system; however, if justified, energy recovery could be added later.

The recommended backup design system is a positive ion source, with no recovery of the energy in the charged-particle fraction (Case 1 in Table 8.1). That design is a direct extension of the 150-keV TFTR beams and injectors to slightly higher voltage and to twice the current per injector. TFTR positive ion injector experience is directly applicable here.

The more complex negative ion reference scheme requires the development of suitable negative ion sources, but requires lower electrical power input. That suggests its use for a steady-state, beam-driven EPR. The positive ion backup scheme is a direct extension of the TFTR injectors, requiring less development, but requiring more electrical power input. This suggests its use for an igniting EPR with an energy storage system supplying this start-up power.

8.4 Principal Development Requirements

All parts of the reference design beam system except the negative ion source are taken directly and scaled up in performance where needed from the ORNL beam system design for TFTR.¹ The beam transport efficiencies (η_1 , η_2 , and η_4) shown for the designs in Table 8.1, including the reference design, indicate a considerable improvement over present state-of-the-art injectors as represented by the ORMAK injectors. These efficiencies should be taken as realistic goals for the beam program to achieve by the time the injectors are needed for the EPR. The accumulated experience of the ORNL beam program, progressing in a direct line from ORMAK through the Princeton Large Torus (PLT), the Poloidal

Fig. 8.2. EPR reference design neutral beam injector system.

Divertor Experiment (PDX), ORMAK Upgrade, Tokamak Fusion Test Reactor (TFTR), and The Next Step (TNS) to EPR, will maximize the probability of achieving these goals.

The level of research and development should be substantially increased in the two schemes of direct ion energy recovery and the negative ion source to improve efficiency and reduce overall cost.

REFERENCE

1. *TFTR Neutral Beam Systems Conceptual Design*, ORNL/TM-5227, Oak Ridge (March 1976).

9. VACUUM SYSTEMS*

Large high-speed vacuum systems are required in tokamaks to provide the proper environment of low impurity density for the formation and ignition of a DT plasma. A design for a vacuum system for EPR was produced.^{1,2} This design was then used as the basis for extrapolating to vacuum system designs for two larger reactors, Demo and CPR.

Demo is a proposed 1900-MW(t) tokamak which would demonstrate all the technology required for electrical power generation by fusion, while CPR is a proposed 4300-MW(t) tokamak which would be a true prototype of a commercial fusion power reactor.

9.1 Vacuum Systems for EPR

EPR will require a main vacuum system to remove plasma exhaust, impurities, and outgassing from the plasma chamber, and a beam vacuum system to remove gas cell exhaust, ion source neutral leakage, and gases resulting from dumped unneutralized ions from the beam line. The main vacuum system will consist of a roughing system and cryosorption pumps or a combination of cryopumps and crypsorption pumps. Between-pulse pumpdown time and background gas pressure requirements should be fairly easy to meet. The plasma physics program must determine the requirements for pumping the plasma exhaust during a pulse. Estimates of the required speed during a pulse now range from the insignificant to the unachievable. The beam vacuum system will consist of separate roughing systems, cryopanel, and cryo-sorption pumps for each beam line. The required speeds are enormous, and the geometry must be worked out with great care to allow sufficient conductance so that the necessary pumping speeds may be achieved.

9.2 Main Vacuum System

The main vacuum system of the EPR will be connected to the plasma chamber. It will consist of a roughing system to bring the pressure down to the millitorr range, and a high-vacuum system to remove the plasma exhaust and reduce the pressure before pulses to an acceptable level. The criteria for the required overall speed for the main vacuum system have not yet been developed.

A criterion presently used for small short-pulsed tokamaks covers the base pressure due to plasma vessel outgassing and inleakage. Base pressures of 10^{-8} - 10^{-9} torr are required in these machines. With reasonable care in fabrication, and with provisions for baking out the plasma chamber surfaces, such pressure can be achieved in EPR with a fairly modest vacuum system.

A second criterion covers the speed with which the chamber is evacuated between pulses. It is believed the between-pulse evacuation will be necessary to remove impurities and fusion products from the plasma chamber, so that the subsequent pulse can be started with a virtually pure DT mixture of optimum proportions. A possible rough criterion would be a reduction in helium and impurity pressure by a factor of 10^3 in approximately one-tenth of a pulse duration.

Another possible criterion concerns the removal of particles leaving the plasma during a pulse. At this time it is not possible to determine if vacuum pumping during a pulse will be effective in reducing sputtered impurities, neutral buildup, and charge-exchange losses, or in preventing damaging plasma density gradients. Perhaps recycling of the plasma will make pumping during a pulse unnecessary and ineffectual. If there is a need for the vacuum system to maintain a low peripheral neutral density during the pulse, it may represent a most stringent criterion.

*This work was performed in conjunction with J. S. Watson, S. D. Clinton, and P. Fisher.

9.2.1 System Description

The roughing system for EPR is designed to allow for rapid pumpdown to minimize delays resulting from repairing and testing. The system will consist of three 5000-cfm Roots Blowers coupled to three 500-cfm mechanical pumps. Details of the arrangement have not yet been settled, but each of the three pump blower combinations will probably be connected to separate vacuum ports which will be described later. Bypass lines will be provided to allow the mechanical pumps to evacuate the plasma chamber to about 10 torr before the blowers are used. The three blowers are capable of reducing the pressure in the chamber to the millitorr range prior to the utilization of the high-vacuum system. The roughing system will be isolated from the chamber before the introduction of any material.

The exhaust from the blower mechanical pump combinations will be passed through a cold trap and then through a tritium removal system consisting of heater, catalyst beds, chillers, and dryers.

The high-vacuum system will consist of 40 cryosorption pumps each attached to the plasma chamber through a duct approximately 0.8 m² in cross-sectional area.

The ducts will have double bends and the appropriate configuration of voids and absorbers surrounding them in order to minimize neutron and gamma streaming. They will have an equivalent length of 5-6 m.

The vacuum pumps as originally conceived would be similar in design to the commercially available cryosorption pumps, although much larger. These pumps contain a molecular sieve-coated panel operated at 4.2 K at the center. An array of chevrons which is cooled to approximately 77 K is placed near the pump entrance. A second array of chevrons at 15 K is located between the panel and the 77-K array to limit radiation heating of the sieve coating. Information recently obtained indicates that bare metal cryopanel operating at 4.2 K can condense layers of hydrogen more than thick enough to meet the storage capacity requirements of EPR. In fact, recent experiments at ORNL³ indicate that cryosorption pumps will achieve high hydrogen storage capacities in the condensation rather than sorption mode with no decrease in pumping speed. Thus, there is the possibility of using pumps consisting of bare cryopanel cooled to 4.2 K surrounded by a single array of chevrons at 77 K to pump all but helium. This type of pump might be more compact and impose a smaller load on the refrigeration system than a cryosorption pump of comparable speed. A cryopump would require a cryosorption panel within it or attached to it in order to pump the helium which will be produced by fusion reactions.

Another alternative pump design might utilize sorption panels cooled to 15 K by helium vapor surrounded by 77-K chevrons to pump hydrogen and impurities, with a smaller cryosorption panel to pump helium. This design might also result in lower refrigeration requirements and a more compact design, but there are uncertainties concerning the storage capacity and pump characteristics for hydrogen isotopes on a molecular sieve at 15 K.

All pumps working in either cryosorption or cryo-condensation mode must be regenerated by allowing their temperature to rise. In order to achieve a balance between convenience and minimization of costly and dangerous tritium buildup on the pumps, each of the high-vacuum pumps will be regenerated once every 24 hours. It is expected that regeneration times can be kept to one hour or less, so that no more than two or three pumps will be out of service at any one time. Regeneration of 77-K arrays in the pumps to remove impurities will require more time, but will be necessary far less frequently.

Two valves will be located between each pump and its associated duct. One valve will be used for rapid closure to prevent pumpout of fuel prior to and perhaps during a pulse. The other will be used to isolate the pump from the chamber while the pump is being regenerated. The rapid closure valve will not have to seat tightly, but must reduce the conductance to the pump to insignificant levels. The isolation valves must seat very tightly since the pressures in the pump during regeneration may exceed 20 torr.

9.2.2 Performance Evaluation

The vacuum system may be judged on the basis of three criteria: maintenance of low outgassing base pressure, pumpdown between pulses, and handling of particle flux during pulses. The main vacuum system consists of 1-m diameter cryosorption pumps (or equivalent) at each of 40 vacuum ports.

The roughing system is designed to reduce the pressure in the torus to the millitorr range quickly enough so that the pump-down time does not become a cause of start-up delay. The plasma chamber and vacuum duct volume is 830 m³. The speed of the mechanical pumps is 42 m³/min. If these pumps alone are used for pump-down to 10 torr, it will take 86 min to reach that pressure. The time required to pump down from 10 torr to 10⁻³ torr after introducing blowers can be estimated from performance curves. These curves show that the efficiency of a Roots Blower, coupled to a mechanical pump of 1/10 its capacity, varies between 60 and 90% over the pressure range of interest and yields a pump-down time of 24 min.

Each vacuum pump is assumed to have a pumping speed of approximately 50,000 liters/sec for DT and 40,000 liters/sec for air at 300 K. Each duct is assumed to be equivalent to a cylinder 1 m in diameter and 5 m in length. Thus, the overall pumping speed of the system is 1,120,000 liters/sec for DT and 650,000 liters/sec for air, assuming an average temperature of 550 K in the duct and 300 K at the pump.

The volume of the main vacuum chamber, including ducts, is approximately 830,000 liters and the surface area is 1.2×10^7 cm². If the outgassing rate is 10⁻¹⁰ torr-liter/sec of air per cm², the base pressure of air will be 2×10^{-10} torr.

The time to reduce the DT pressure by a factor of 10³ is 5 sec, while to reduce air pressure by a similar factor, it takes 9 sec.

The calculated particle load leaving the plasma during a pulse is 2×10^{21} /sec. If one assumes that this particle flow is not recycled by the plasma but results in neutrals in an annular region between the plasma and the chamber wall, until pumped by the vacuum system, the neutral density in this region will be

$$\frac{2.1 \times 10^{21} \text{ particles/sec}}{1.12 \times 10^9 \text{ cm}^3/\text{sec}} \approx 2 \times 10^{12} \text{ particles/cm}^3.$$

9.3 Beam Vacuum System

The EPR will use up to 100 MW of 200 keV neutral beam injection to heat the plasma. The six beam lines can each contain two 8.5-MW beam sources. The beams will be of the negative ion type, utilizing sources which have yet to be developed. The negative ions from the source will be accelerated in a grid region to 200 keV. They will then pass through a stripping cell where the pressure will be maintained for optimum neutral beam production. Unneutralized negative ions as well as positive ions will be diverted by a magnetic field at the exit of the gas cell, while high-energy neutrals stream into a drift tube and hence into the plasma chamber.

A neutral beam colliding with gas leads to the production of positive ions. These ions will be deflected by the EPR's magnetic fields into the drift tube wall, representing loss of energy from the beam as well as a source of damage to the drift tube material; thus the particle density downstream of the gas cell must be minimized.

A vacuum system must be provided to pump neutral gases which flow from the ends of the stripping cell and the ion source, as well as the gas formed from the dumping of the unneutralized ions. At this time the ion source is undefined and no attempt was made to deal with the gas load it will represent. The criterion for the particle density downstream of the gas cell was based on limiting the beam loss and high-energy particle impact on drift tube walls to acceptable levels.

9.3.1 System Description

Enormous vacuum pumping speeds are required to maintain adequately low particle densities downstream of the stripping cell. The only feasible alternatives are cryogenic techniques. The most promising candidate is a large configuration of cryopanel operated at 4.2 K screened by 77-K chevron arrays, supplemented by a small cryosorption system to absorb helium issuing from the drift tube. The cryopanel and their associated chevrons will almost totally enclose the area from the stripping cell to the entrance of the drift tube.

Because significant quantities of tritium will enter this region from the drift tube, the beam vacuum system will be designed and enclosed in accordance with EPR tritium-handling criteria.

Regeneration of the cryopanel and cryosorption system will require an isolation valve between the beam line and the plasma vessel. The cryopanel will be exposed to neutron and gamma streaming and radiant heating from both the plasma and the beam line.

9.3.2 Performance Analysis

The beam vacuum system will handle gas loads from the ion source, the stripping cell, and the ion dump. The ion source design is not known at this time and the gas load from the source was not considered, although accommodating it may require significant changes in the vacuum system.

The stripping cell was assumed to be approximately 30 cm in diameter, which would be sufficient to handle 8.5 MW of neutral beam. Calculations were made to determine the optimum product for length and gas density in the cell. These calculations were based on a cross section of $6 \times 10^{-16} \text{ cm}^2/\text{module}$ for $D^- \rightarrow D^0$ and $1.5 \times 10^{-16} \text{ cm}^2/\text{module}$ for $D^0 \rightarrow D^+$ both at 200 keV. The cross sections for $D^0 \rightarrow D^-$ and $D^+ \rightarrow D^0$ were insignificant in comparison to the other two.

The calculations indicated that the optimum length and density product would be approximately $3.1 \times 10^{15} \text{ molecules/cm}^2$ which yields a net fraction of approximately 0.63 D^0 . If one chooses a length of 300 cm for the stripping cell, then one arrives at a pressure of $3 \times 10^{-4} \text{ torr}$ at 300 K.

Gas will flow from the stripping cell at either end, and if the gas is introduced to the cell at the middle of its length, a pressure gradient will exist outward from the middle toward either end. Since this pressure is in the molecular flow regime for a 30-cm-diameter tube, the pressure gradient will be linear. With a linear pressure gradient and the assumption that pressure just outside the stripping cell is much less than the pressure at the middle of the cell, it is possible to calculate the gas flow from the cell. With a linear gradient, the pressure at the middle of the cell should be twice the required average pressure. The conductance of the tube was calculated at $8.4 \times 10^3 \text{ liters/sec}$ from the middle to either end for D_2 at 600 K. This, combined with a particle density of 2.1×10^{13} at the middle of the tube, yields a flow of 1.8×10^{20} particles per second from each end. In addition, unstripped or ionized particles from the cell amount to 4.7 to 9.4×10^{19} particles per second, depending on the extent to which they form D_2 molecules.

Given the assumption that these particles inhabit a region between the stripping cell and the plasma 300 cm in length, and given the allowable ionization fraction of the neutral beam, the allowable particle density and thus the required pumping speed outside the stripping cell may be calculated. If one chooses 2% allowable ionization, the allowable average density over 300 cm is $4.4 \times 10^{11} \text{ molecules/cm}^3$ and the required pumping speed is $9 \times 10^5 \text{ liters/sec}$.

A conservative rating for a cryopanel assembly is 5 liters/sec per cm^2 for D_2 , and so approximately $2 \times 10^5 \text{ cm}^2$ of cryopanel are required for each 8.5-MW beam.

9.4 Extrapolation of EPR Design to Larger Reactors

As a first approximation, the main vacuum systems for the proposed Demo and CPR are assumed to vary roughly with the surface area of the plasma chamber, judging from the fact that the amount of space available for locating the ports is proportional to the surface area. It is assumed that the same size ducts and pumps will be used for the three plants, since ducts greater than 1 m in diameter will lead to even less tolerable neutron and gamma streaming problems. The hazard associated with the tritium in each pump, of course, increases with pump size.

The configurations and speeds for the main vacuum systems for the three reactors, in accordance with the above and the assumption that duct length remain constant at 5 m, are given in Table 9.1.

Table 9.1. EPR, Demo, CPR vacuum systems

	<u>EPR</u>	<u>Demo</u>	<u>CPR</u>
Major radius	6.75	10.5	13.5
Minor radius	2.25	3.5	4.5
Configuration	40 1-m ducts and pumps	80 1-m ducts and pumps	160 1-m ducts and pumps
Speed for DT (liters/sec)	1.12×10^6	2.24×10^6	4.48×10^6
Speed for air (liters/sec)	6.5×10^5	1.3×10^6	2.6×10^6

Note that in all cases each pump must be enclosed in a secondary enclosure, and regeneration equipment must be provided.

Using assumptions outlined in the preceding discussion of the EPR vacuum systems (i.e. outgassing of 10^{-11} torr- ℓ /sec-cm², pump-down of airlike impurities by a factor of 10^3 between pulses, and existence of an annular region where neutrals bounce between the plasma and walls until pumped) performance can be estimated from Table 9.2.

Table 9.2. Pump-down time and neutral density

	<u>EPR</u>	<u>Demo</u>	<u>CPR</u>
Base air pressure (torr)	2×10^{-10}	2×10^{-10}	2×10^{-10}
Pumpdown time between pulses (sec)	9	15	16
"Neutral Density" (cm ⁻³)	5×10^{12}	1.1×10^{13}	1.0×10^{13}

The beam vacuum system scaling can be based on the fact that an 8.5-MW beam requires 20 m² of cryopanel.

The roughing system for EPR consists of three 5000-cfm Roots Blowers each coupled to a 500-cfm mechanical pump. Demo should use about ten blower-pump combinations and CPR twenty, although fewer units of larger size might be used.

REFERENCES

1. J. S. Watson et al., *Scoping Studies of Tritium Handling in a Tokamak Experimental Power Reactor*, ORNL/TM-5080, Oak Ridge (January 1976).
2. M. Roberts et al., *Oak Ridge Tokamak Experimental Power Reactor Study Reference Design*, ORNL/TM-5042, Oak Ridge (November 1975).
3. C. D. Scott et al., *Experimental Engineering Section Semiannual Progress Report (Excluding Reactor Programs)*, Sept. 1, 1975 to Feb. 29, 1976, ORNL/TM (in preparation).
4. C. M. Van Atta, *Vacuum Science and Engineering*, p. 197, McGraw-Hill, New York, 1965.

10. ADVANCED FUEL PROSPECTS

10.1 Use of DD and D³He in Testing

The use of pure DT fuel in potential reactor situations may (and should eventually) lead to ignition of the nuclear fuel at about 6 keV with a consequent thermal runaway to 10 keV or higher, dependent on the transport loss modes operative. This could result in a large power load on the first wall with possible first-wall damage and severe outgassing of occluded DT fuel which may then feed the nuclear fire. Because of the presence in the reacting plasma of a positive density coefficient (i.e., an increase in the fuel density increases the output power level), the ignited plasma may undergo a second more severe power transient, with associated increased damage to the first wall, especially if the dissipative trapped ion mode is not present.

In view of these potential difficulties during ignition start-up, it is incumbent to perform preliminary studies on more controllable situations such as plasmas fueled by DD and D³He. The nuclear ignition temperature for these fuels is well above 6 keV, so that nuclear runaway cannot occur, and yet important plasma physics and reaction physics can be studied using the 2.5-MeV neutron from $d(d,n) {}^3\text{He} + 3.3 \text{ MeV}$ and the 14-MeV proton from $D({}^3\text{He}, p) \alpha + 18.4 \text{ MeV}$. Also, the 0.8-MeV ³He ion from the first reaction, and the 3.0-MeV proton and 1.0-MeV triton from $d(d,p) t + 4.0 \text{ MeV}$, will be of value to the pre-ignition studies. These particles will permit important diagnostic studies to be made leading to evaluation of charged-particle confinement, slowing down properties, *in situ* beam-plasma reaction properties, and ion temperatures. The power levels for DD and D³He plasmas will be lower by approximately two orders of magnitude than those of DT plasmas, so that alpha-particle sputtering effects, neutron radiation damage, and neutron activation will be minimal compared to the more reactive pure DT plasma operation. In addition, since the 14-MeV neutrons arise principally from the second-generation slowing down or beam-plasma reactions $d(t_{\text{fast}}, n)\alpha$, they can serve as selective monitors of plasma confinement and slowing down of the energetic tritons which originate at 1-MeV energy from $d(d,p) t$.

10.2 Thermal Runaway Without the Dissipative Trapped-Ion Mode Losses

In the presence of the dissipative trapped-ion mode losses, the thermal runaway of the plasma temperature following nuclear ignition of DT fuel will be limited to 10-20 keV (depending on size)¹ because of the strong dependence of $n\tau$ on $T_e^{-1/2}$.² Should this loss mode be absent (due to plasma and temperature profile effects or to non-applicability) the ion temperature can escalate to 200-400 keV ($T_e \sim 100 \text{ keV}$), since neoclassical confinement improves with temperature increase and the equilibrium is then determined principally by synchrotron radiation losses (see Fig. 10.1) and eventually by ash buildup (see Fig. 10.2).

10.3 DD Advanced Fusion Fuel Prospects

Such high-temperature DT runaways would automatically place the plasma operation in the Flux Conserving Tokamak (FCT) mode³, and fueling of DD would be not only possible but extremely appropriate⁴ (provided, of course, that pellet fueling or some other fueling technique works). This would permit programming a low-density ignited DT plasma to move into a high-density high-power DD burning regime, with the concomitant advantage of excess or "free" neutrons except for those required to replace the tritium burned in the ignition phase.

It should be noted that the excess tritium produced by DD reactions in the total burn period may obviate the need for any external tritium breeding blanket. One may thus be able to replace the

Fig. 10.1. Schema of burning characteristics of DT systems. The operating point has a desirable negative temperature characteristic (but an undesirable positive density characteristic).

Fig. 10.2. Radiation-dominated DT reactor (pulsed burn, closed system).

Fig. 10.3. Fractional burn-up in percent in steady-state catalyzed DD reactors.

molten lithium blanket heat-exchanger and thereby eliminate the potential chemical hazard due to a common-mode failure rupture between the water moderator region and the molten lithium region. Alternative blankets include (1) chemically fixed sodium (either in water or as a molten salt) with its 12.6-MeV energy release per neutron capture; (2) a sub-critical U blanket; (3) a thorium blanket to breed ^{233}U ; (4) a depleted uranium blanket to breed ^{239}Pu ; or (5) a burn-up region for actinide and fission product consumption.

The *in situ* burn-up of fuel in a dense quasi-steady state plasma is given by $\text{FBX} = 1/(1 + 1/n_F \tau \langle \sigma v \rangle_{XF})$ and characteristic values are given for deuterium-rich plasmas in Fig. 10.3. Burn-up of the radioactive tritium in plasmas with high ash content is extremely large (50-80%), provided only that the high temperatures and densities are maintained for at least some tens of seconds.

10.4 Noncircular Tokamaks

The noncircular tokamaks can in principle provide higher β 's [$\beta_{nc} \sim \beta_c (1 + \epsilon^2)/2$] than circular-shaped tokamak cross sections. Again in the absence of the dissipative trapped ion mode, the high temperatures along with the high β 's permit the transfer to the DD fueled regime with its many advantages. These advantages accrue independently of the FCT mode but would be further enhanced with FCT burning.

REFERENCES

1. J. Rand McNally, Jr. et al., "Fusion Dynamics of Lossy Reactor Systems and Power Balance Implications", *Plasma Physics and Controlled Nuclear Fusion Research*, Vol. 3, p. 511 (1975).
2. S. O. Dean et al., *Status and Objectives of Tokamak Systems for Fusion Research*, WASH-1295, AEC, Washington, D.C. (1974).
3. J. F. Clarke, *Bull. Am. Phys. Soc.* 20: 1228 (1975); J. D. Callen and R. A. Dory, *Phys. Fluids* 15: 1523 (1972).
4. J. Rand McNally, Jr., "Advanced Fuels for Nuclear Fusion Reactors," *Proc. Conf. Nuclear Cross Sections and Technology*, Vol. 2, p. 683 (1975); "Advanced Fusion Fuels," *Proc. 6th Symposium on Engineering Problems of Fusion Research*, p. 1012 (1976).

THIS PAGE
WAS INTENTIONALLY
LEFT BLANK

INTERNAL DISTRIBUTION

1. S. E. Attenberger	67. A. T. Mense
2. L. A. Berry	68. O. B. Morgan
3. E. S. Bettis	69. ORNL Patent Office
4. J. D. Callen	70. Y-K. M. Peng
5. J. F. Clarke	71. H. Postma
6. CTR Reports Office	72. M. Roberts
7. C. A. Flanagan	73. M. W. Rosenthal
8. G. G. Kelley	74. G. Schilling
9. G. D. Kerbel	75. Y-12 Document Reference Section
10. Laboratory Records Department	76. T. E. Shannon
11. Laboratory Records - ORNL RC	77. R. Sharp
12-13. Central Research Library	78. D. Steiner
14. J. W. Lue	79. L. D. Stewart
15. F. B. Marcus	80. D. Strickler
16. J. R. McNally	81-82. Thermonuclear Division Library
17-65. D. G. McAlees	83. J. S. Watson
66. H. C. McCurdy	84. N. A. Ukan

EXTERNAL DISTRIBUTION

85. C. C. Baker, General Atomic Company, San Diego, California 92138.
86. R. N. Cherdack, Burns & Roe, Inc., 283, Highway 17, Paramus, New Jersey 07652.
87. F. E. Coffman, Division of Magnetic Fusion Energy, Energy Research and Development Administration, Washington, D.C. 20545.
88. Robert W. Conn, Nuclear Engineering Department, University of Wisconsin, Madison, Wisconsin 53706.
89. N. Anne Davies, Mail Stop G-234, Division of Magnetic Fusion Energy, Energy Research and Development Administration, Washington, D.C. 20545.
90. S. O. Dean, Mail Stop G-234, Division of Magnetic Fusion Energy, Energy Research and Development Administration, Washington, D.C. 20545.
91. Dr. H. K. Forsen, Vice President, General Manager, Laser Enrichment Department, Exxon Nuclear Co., Inc., 777-106th Avenue NE, C-00777, Bellevue, Washington 98009.
92. T. K. Fowler, University of California, Lawrence Radiation Laboratory, P. O. Box 808, Livermore, California 94551.
93. H. P. Furth, Princeton Plasma Physics Laboratory, Princeton University, P. O. Box 451, Princeton, New Jersey 08540.
94. G. Gibson, Westinghouse Electric Corporation, Fusion Power Systems, Cheswick, Pennsylvania 15024.
95. M. B. Gottlieb, Plasma Physics Laboratory, Princeton University, Princeton, New Jersey 08540.
96. W. C. Gough, Electric Power Research Institute, Palo Alto, California 94304.

97. Roy W. Gould, Building 116-81, California Institute of Technology, Pasadena, California 91109.
98. J. Nelson Grace, Mail Stop G-234, Division of Magnetic Fusion Energy, Energy Research and Development Administration, Washington, D.C. 20545.
99. Harold Grad, Courant Institute, New York University, 251 Mercer St., New York, New York 10012.
100. R. Harder, General Atomic Company, San Diego, California 92138.
101. C. Henning, Mail Stop G-234, Division of Magnetic Fusion Energy, Energy Research and Development Administration, Washington, D.C. 20545.
102. R. L. Hirsch, Energy Research and Development Administration, Washington, D.C. 20545.
103. E. E. Kintner, Mail Stop G-234, Division of Magnetic Fusion Energy, Energy Research and Development Administration, Washington, D.C. 20545.
104. G. Kulcinski, Nuclear Engineering Department, University of Wisconsin, Madison, Wisconsin 53706.
105. Librarian, Culham Laboratory, United Kingdom Atomic Energy Authority, Abingdon, Berkshire, United Kingdom.
106. Plasma Physics Library, Plasma Physics Laboratory, Princeton University, P. O. Box 451, Princeton, New Jersey 08540.
107. L. M. Lidsky, Department of Nuclear Engineering, Massachusetts Institute of Technology, Cambridge, Massachusetts 02139.
108. R. Mills, Plasma Physics Laboratory, Princeton University, Princeton, New Jersey 08540.
109. M. Murphy, Mail Stop G-234, Division of Magnetic Fusion Energy, Energy Research and Development Administration, Washington, D.C. 20545.
110. J. O. Neff, Mail Stop G-234, Division of Magnetic Fusion Energy, Energy Research and Development Administration, Washington, D.C. 20545.
111. T. Ohkawa, General Atomic Company, San Diego, California 92112.
112. J. Powell, Brookhaven National Laboratory, Upton, Long Island, New York 11973.
113. Research and Technical Support Division, Energy Research and Development Administration, Oak Ridge, Tennessee 37830.
114. F. L. Ribe, Los Alamos Scientific Laboratory, P. O. Box 1663, Los Alamos, New Mexico 87544.
115. D. J. Rose, Department of Nuclear Engineering, Massachusetts Institute of Technology, Cambridge, Massachusetts 02139.
116. M. N. Rosenbluth, School of Natural Sciences, Princeton University, P. O. Box 451, Princeton, New Jersey 08540.
117. Z. M. Shapiro, Westinghouse Electric Corporation, Fusion Power Systems, Cheswick, Pennsylvania 15024.
118. Anthony Shu, Mail Stop G-234, Division of Magnetic Fusion Energy, Energy Research and Development Administration, Washington, D.C. 20545.
119. Weston M. Stacey, Jr., ANL-CTR Program, Building 208, Argonne National Laboratory, 9700 South Cass Avenue, Argonne, Illinois 60439.
120. C. Taylor, Controlled Thermonuclear Research, Mail Code L-382, Lawrence Livermore Laboratory, P. O. Box 808, Livermore, California 94550.

121. T. C. Varljen, Westinghouse Electric Corporation, Fusion Power Systems, Cheswick, Pennsylvania 15024.
122. J. M. Williams, Mail Stop G-234, Division of Magnetic Fusion Energy, Energy Research
123. Herbert H. Woodson, University of Texas, Austin, Texas 78712.
- 124-326. Given Distribution as shown in TID-4500 - Controlled Thermonuclear Processes and Plasma Physics.