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by

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Abstract

During FY 82, much effort was devoted to work in support of alternate confinement concepts, especially those involving field reversal. This work includes: 1) development of particle and energy confinement scaling for the Reversed Field Pinch (RFP) and the Field Reversed Theta Pinch (FRTP), and 2) analysis of start-up (heating and plasma build) for the spheromak and Field Reversed Mirror (FRM).

In addition, a block of projects were concerned with fusion product effects, including heating and ash build-up. These include, 1) a study of possible use of radial electric fields to control ash build-up in tokamaks, 2) effects of alpha-driven microinstabilities on heating in tokamaks, and 3) fusion product transport, including effects of large angle scattering on orbits, in EBT and FRM devices.

In a related study, the possibility of hot-ion mode operation (assuming strong transfer of fusion product energy to ions, e.g. via microinstabilities) was done with emphasis on calculation of ion-electron equilibration rates.

Among the major contributions from this work are a new integral transform method for fast calculation of 2-D magnetic field equilibria and the identification of conditions where alpha-driven instabilities are most strongly excited.

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I. Introduction

Effort in FY 1982 has focused on key areas in support of field-reversed, alternate concept devices, leading to a clearer understanding of scaling laws for use in reactor studies. Additionally, the critical role of fusion products in the plasma burn has required that their transport and energy transfer effects be examined in detail for the alternate concepts as well as main line devices.

In Section II.1 a theoretical model is developed to account for particle loss in the field-reversed theta pinch (F RTP) based on loss-cone scattering. Rotation, as observed in experiment, was then predicted based on our model of conservation of angular momentum. Reactor designs, using scaling laws derived in our studies, were found to be attractive for an FED-type device. In Section II.2 the development of a 1 1/2-D code for equilibrium-transport in an arbitrary aspect ratio spheromak is described. This features a fast, integral transform method for the equilibria calculations. Also, energy extraction by inductive coils and plasma stabilization by gyroscopic motion are discussed as means for improving spheromak reactor design. In Section II.3 the 6-D Fokker-Planck kinetic equation for the F RTP is solved and shown to agree well with experimental findings. Section II.4 describes the effects of electron currents in the start-up phase of a field reversed mirror (FRM). This analysis may explain why efforts to reverse the 2XIIB field have failed and possible remedies. In Section II.5 fundamental work on a transport code for the reversed field pinch (RFP) device is described based on a braided magnetic field model. Comparison with experimental results from ZT-40 suggests that

a "dynamo" model is also needed to give an accurate picture of the RFP discharge.

In Section III.1 the imposition of radial electric fields are found to improve overall tokamak performance by increasing the fuel confinement time. Methods to provide the fields, consistent with reactor designs, are suggested. Section III.2 discusses the transport of fusion products in the EBT and FRM devices. Prompt loss fractions in EBT are found to be larger than previously assumed. Nuclear elastic collisions have been studied as a mechanism for increasing fusion product energy deposition in the FRM. In Section III.3 fusion product driven microinstabilities are discovered to significantly enhance energy transfer to background ions via wave-particle interactions. However, when this model is included in a reactor code the overall effect is diminished. Finally, in Section III.4 several methods are discussed for decoupling the energy transfer between ions and electrons in fusion plasmas. These are shown to have a dramatic effect on the required ignition condition especially for advanced fuels.

II. ALTERNATE CONCEPTS

II.1 LOSS-CONE SCATTERING/ROTATION MODEL IN THE FIELD REVERSED THETA PINCH AND APPLICATIONS TO REACTORS

The Field-Reversed Theta Pinch (FTRP) represents a promising confinement approach to fusion power. However some experiments have encountered rapid particle losses and plasma disruption after a period of increasing rotation frequency. The objective of the present research was to develop theoretical models to explain the phenomena and simultaneously explore ways to reduce losses and spin-up. Further, reactor scenarios were studied as an extension of the particle and energy loss models.

II.1.1 Confinement and Rotation.

To account for the rapid particle losses, an ion transport theory was proposed based on the existence of a "loss-cone-like" region in H-P space (or a corresponding one in velocity space) for the FTRP (see last years progress report). This loss region was observed to have an effect on particle losses analogous to that of the conventional mirror loss cone, although now the loss region varies with radial position.

Particle confinement times predicted by this mechanism were found to fall within a factor of two compared with a wide variety of past and recent experiments on the FTRP (see Table I). The predicted confinement time and corresponding dynamics of plasma parameters also compare favorably with experimental observations from earlier FTRP devices and also the more recent FRX-A and B experiments at the Los Alamos National Laboratory, and the TRX-1 at Mathematical Science Northwest Inc.

This "loss-cone-like" transport is also shown to be consistent with plasma rotation measurements. It is proposed that the mechanisms associated with ion rotation are: 1) loss of angular momentum associated with particle

Table I. Projected Operating Parameters for FRX-C and Parameters for FRX-A and B for Comparison

Parameters	FRX-C		FRX-A	FRX-B
	Set #1	Set #2		
B_{ext} (T)	.9	.9	.68	.75
τ_B (μs)	200	200	150	150
P_o (mTorr)	2.4	24	7	9
n_m (10^{15}cm^{-3})	1.2	6.8	3.6	4.1
T_i (eV)	1140	200	185	240
T_e (eV)	480	85	120	100
r_c (cm)	45.0	45.0	25.0	25.0
R (cm)	6.7	8.9	4.0	4.0
a (cm)	3.4	3.4	1.4	1.4
L (cm)	54	54	35	36
ρ_i (cm)	.78	.33	.29	.31
K	.55	.75	.6	.6
N_L (10^{18}cm^{-1})	.2	2.0	2.7	3.2
Obs. τ_p (μs)	—	—	20 *	30 *
Pred. τ_p (μs)	340	105	51	63

* Observed τ_s .

leakage; and 2) "loss" of plasma moment of inertia due to changes of the density profile associated with particle and energy transport. This is shown graphically in Fig. 1. The growth time of the ion rotation frequency is found to be most strongly influenced by the particle and energy confinement times. This model gives numerical results which appear to be consistent with measured parameters within the experimental uncertainties. Typical results are shown in Fig. 2.

II.1.2. F RTP Prototype Power Reactors

The potential advantages of the F RTP reactors includes: 1) high β and hence high power density, 2) compactness, 3) high density (10^{15} - 10^{16} cm^{-3}) and readily hot (1-3keV) plasma formed by implosion, compatible with further (neutral- beam) heating, 4) simple linear geometry and hence easy modular production and maintenance, and 5) easily coupled direct energy conversion.^(1,2) An application of neutral-beam and pellet injection would compensate particle losses and hopefully suppress the rotational instability. As a result, neutral-beam heating and pellet refueling are assumed in our reactors, serving simultaneously for heating, refueling and maintaining the field reversal and the plasma moment of inertia.

Two operational schemes are considered. The first is called the "in-situ-burn" scheme, and the second called the "moving plasmoid" scheme. The first scheme (See Fig. 3) is a stationary one in which the plasma remains in the formation chamber throughout the plasma lifetime. The neutral beam and the pellet injection are applied during the heating (after implosion), and burning stages. The typical temperature history assumed for the plasmoid is shown in Fig. 4. The temperature obtained from implosion is usually 10-25% of the final burn temperature depending on the capacitor bank energy and implosion efficiency. After implosion, the neutral beam then takes over the task of further heating.

Figure 1. Spinning-up sequences of ion plasma. At $t = 0$, plasma does not rotate. At $t_1 > 0$, an initial rotation is set up due to angular momentum carried away by the lost particles. At t_2 , "loss" of moment of inertia further accelerates such a rotation.

Figure 2. Time history of the predicted and the measured ion rotation frequency (normalized by diamagnetic drift frequency, i.e., $\alpha = \omega_i / \omega_{D0}$) for FRX-B with initial conditions: $T_i = 200$ eV, $T_e = 120$ eV, $T_e = 120$ eV, $R = 4$ cm, $L = 36$ cm, $K = 0.6$ and $B_{ext} = 0.68$ T.

Figure 3. A "in-situ-burn" scheme of the Prototype F RTP reactor.

Figure 4. An assumed ion temperature history of the F RTP reactor plasma.

The neutral beam can be applied off-axis at the axial ends in the direction opposed to the ion rotation such that it may help stabilize the rotational instability. An alternative approach is to apply the neutral beam through the openings in the pinch coil. However, due to the limited cross-sectional area available in each opening, we might have several openings along the pinch coil.

The second scheme is based on the "Moving Plasmoid Heater (MPH)" concept, proposed by J. Galambos et al.,⁽³⁾ In this scheme, the plasma, once formed, is transported to a heating and compression chamber, where the neutral-beam injection (possibly with an additional adiabatic compression by external fields) is applied. After the plasma is sufficiently hot, it is then moved to the burn chamber specifically designed for fusion product energy conversion. At this stage the fuel pellets are injected to maintain the particle density in order to prevent the rotational instability, and hence to obtain a higher fractional burn-up. Finally, the plasma is moved to an expansion chamber, in which the plasma thermal energy and the associated field energy are recovered through electrostatic collection and magnetic coupling. The present scheme differs from the original MPH concept in that we add pellet injection in the burn stage.

The first scheme might be preferred as a prototype reactor due to its natural simplicity in engineering design and operation. However, the stationary scheme may be limited to lower burn-up design due to neutron or radiation wall loading problems. The moving plasmoid scheme is apparently much more complicated than the first one for control of the moving plasmoid, and may face thermal fatigue and other problems due to the intrinsic transient behaviors.

II.1.3 Plasma Burn Model and Reactor Design Considerations

A prototype D-T fueled FRTP power reactor is designed based on the "in-situ-burn" (stationary) operational scheme. The operation parameters are presented in Table II. In this design, we have assumed a neutron wall loading limit of 2MW/m^2 and a particle confinement scaling of "loss-cone-like" scattering transport (LCST)^(4,5). The particle confinement time τ_p^{LCST} thus scales as

$$\tau_p^{\text{LCST}} \sim T_i^{2.2} K^{2.7} r_c^{1.8} M_R^{1.3} B^{2.5} n_m^{-1} \quad (1)$$

where K is the rigid-rotor field-reversal factor, r_c is the coil radius, M_R is the mirror ratio for the coil configuration, B is the field at the coil and n_m is the particle density. The LCST scaling is in good agreement with present experiments (i.e., FRX-B).⁽⁶⁾

A plasma major radius of 30 cm is assumed with a blanket thickness of 45 cm. According to the 1-D exponential rigid-rotor equilibrium model we have a plasma average beta $\langle\beta\rangle$ of 0.89 and $K (\equiv \tanh^{-1} B(r=0)/B_\infty)$ of 0.6. Since the ion-electron equilibration time is less than 10% of the particle confinement time predicted, we have chosen identical electron and ion temperatures. On the other hand, the energy confinement time is assumed to be 80% of the particle confinement time which is consistent with experiment also.

The implosion energy W_{imp} is calculated with an assumed implosion efficiency of 4% (we have assumed an axial and radial compression) and the temperature obtained by implosion is 10% the final temperature. The injection power P_{in} is calculated to balance the total energy loss including both Bremsstrahlung and Cyclotron radiation energy loss and the energy required to heat the plasma to the final temperature after implosion. The injection period (τ_{in} , See Fig. 4) is chosen to equal one half the

energy confinement time. In addition, we assumed an 80% trapping efficiency for the neutral beam. The operating temperature is chosen to minimize the implosion energy and the neutral beam injection power while satisfying the neutron wall loading limit. Ignition during the burn phase is obtained at the final operating temperature when charged fusion energy balances the losses due to radiation, conduction and convection.

We assume a burn time $\tau_{\text{burn}} = 5\tau_p$ based on maintaining a stable equilibrium. The refueling rate N_r is calculated according to the requirement for compensating the particle loss, i.e., $N_r = N/\eta\tau_p$, where N is the total number of particles inside the separatrix and η is the trapping efficiency for the pellets assumed to be 90% in this design. Neither the buildup of fusion product ash or impurities are considered in the burn analysis.

The plasma Q value is defined as

$$Q = \frac{P_F \tau_{\text{burn}}}{P_{\text{in}} \tau_{\text{in}} + W_{\text{imp}}} \quad (2)$$

where P_F is the total fusion power. The net electric power output P_{net} is obtained by assuming a 36% blanket thermal energy conversion efficiency, with a neutron multiplication factor of 1.22, a 70% direct energy conversion efficiency of the fusion charged particles, and also a 25% recovery efficiency for the wasted implosion energy. In considering the time average, we have the net power output

$$P_{\text{net}} = \left[(36\% (P_N + P_R) + 22\% P_N + 70\% P_C) \tau_{\text{burn}} + 25\% W_{\text{imp}} - P_{\text{in}} \tau_{\text{in}} - W_{\text{imp}} \right] f_D / (\tau_{\text{burn}} + \tau_{\text{in}} + \tau_{\text{imp}}) \quad (3)$$

where P_C is the fusion power (charged particles) power, P_N is the neutron power, P_R is the total radiation power (including Bremsstrahlung and Cyclotron) and f_D is the duty factor, assumed to be 70% in the present design.

II.1.4 Results and Discussion

In the course of satisfying the neutron wall loading limit and minimizing the total input energy (implosion and beam-heating) we arrived at an operating temperature of 11 keV and an electron peak density of $0.51 \times 10^{15} \text{cm}^{-3}$. The corresponding external magnetic field is thus $\approx 2.2 \text{ T}$ and an average plasma $\langle \beta \rangle$ of 0.9. The required stored implosion energy is found to be 11.5 MJ, while the neutral beam injection power is 39 MW, which are quite manageable as near-term technology. In the meanwhile, we found a plasma Q value of 4.8 and a time averaged net electric power output of 12 MW. The constraint on this reactor is apparently due to the neutron wall loading. However, even with this limitation, the fusion power density of 12 MW/m^3 is quite large compared with the conventional Tokamak reactor designs.

The parameters in Table II represent only a preliminary reference design around which many tradeoff studies can be performed. For example, if larger neutron first-wall loadings could be tolerated (as high as 20 MW/m^3 in Ref. 2) then the fusion power density can be increased by increasing the external B field. The confinement time also increases which in turn greatly improves the plasma Q value and net power output.

In summary, the LCST confinement scaling model is favorable for near-term device success. Indeed, the burn analysis and reactor design for the F RTP indicates that it would be a suitable candidate for a compact, FED-type device.

Table II. Stationary D-T Fueled F RTP Reactor with
Neutral-Beam and Pellet Injection

Parameters	D-T
External field B_{ext} (T)	2.2
Peak density, $n_m (10^{21} \text{m}^{-3})^+$	.51
Ion temp., T_i (keV)	11
Coil/wall radius, r_c/r_w (m)	.9/.45
Major radius x length, $R \times L$ (m)	.3 x 4
Reversal factor, K	.6
Line density, $N_L (10^{21} \text{m}^{-1})$	2.2
Plasma $\langle \beta \rangle$ (average)	.90
Particle confinement time, τ_p (s)	.4
Plasma thermal/Implosion energy (MJ)	4.6/11.5 ⁺⁺
Injection power, P_{in} (MW)	39 [†]
Refueling rate, $N_r (10^{20} \text{sec}^{-1})$	41
Fusion power, P_F (MW)*	42
Power density, $P_D (\text{MW/m}^3)^{**}$	12
Neut./rad. wall loading $P_{\text{NW}}/P_{\text{RW}} (\text{MW/m}^2)^\S$	2/.08
First-wall neut. flux, $\phi (10^{18}/\text{m}^2 \text{sec})$	.9
Plasma Q value	4.8
Net elec. output, P_{net} (MW)	12 ^{††}

+ electron peak density;

++ Assuming 4% implosion efficiency and temperature after implosion to be 10% final temperature.

† Assuming 90% trapping efficiency.

* plasma volume = 2.1 m³;

** first-wall volume = 3.4 m³;

§ first wall surface area = 17 m².

†† Assuming 36% thermal energy conversion efficiency, 70% direct energy conversion efficiency, 25% recovery of implosion energy and a neutron multiplication factor 1.22.

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II.2 TRANSPORT, HEATING, AND REACTOR STUDIES OF HIGH β COMPACT TORI

The studies of transport, heating, and reactor feasibility of spheromak and compact tori concepts can be divided into three main categories: 1) The development and testing of a 1-1/2-D spheromak heating and transport code; 2) an analysis of inductive energy extraction from a compact torus plasma; and 3) the development of a gyroscopic stabilization theory for the spheromak tilt/shift problem. In the present section, we provide basic descriptions of this work and mention some of the key results and accomplishments that have been achieved in each of these three areas of compact torus research during the past year.

II.2.1 1 1/2-D Studies of Spheromak Heating and Transport

Recently, we have used simple global scoping techniques to study the viability of ohmic-compressive ignition of a spheromak plasma^(1,2). Studies of this type are important for two reasons: first, although auxiliary heating (i.e., beam injection) of a spheromak plasma may be possible, a spheromak reactor concept will be especially attractive if ohmic decay and compression alone are found to be sufficient to heat the plasma to an ignition temperature; secondly, it seems likely that the road to fusion demonstration within the spheromak/compact torus context might be greatly shortened and simplified if the ohmic-compressive ignition scheme is determined to be an acceptable route.

The interrelated effects of ohmic heating, adiabatic compression, and collisional transport on spheromak plasma evolution are, however, inherently "two-dimensional" in nature. Consequently, we have developed an arbitrary aspect ratio, poloidal flux-surface averaged 1 1/2-D computer code that will be capable of studying the ohmic-compressive heating question in a self-consistent manner. This algorithm involves an interaction between a 2-D

solution of the Grad-Shafranov equation and a 1-D, flux-surface dependent transport scheme that time steps adiabatic forms of the plasma pressure, temperature, etc. These two main sections of the code are connected with interfacing routines which calculate the necessary flux-surface averaged quantities and time-stepped equilibrium inputs and boundary conditions. This basic scheme is outlined in Fig. 1.

There are two basic differences between existing 1 1/2-D formulations and the code that we have developed at Illinois. The first difference is due to our use of poloidal (rather than toroidal) flux-surface labeling, and hence, its versatility beyond tokamak modeling. This is an important advantage since our code is being designed primarily for handling various aspects of high-beta plasma systems such as the spheromak and RFP. The second important difference is due to our use of semi-analytical, arbitrary aspect ratio solutions to the Grad-Shafranov equation (the "finite integral transform technique"). This technique is especially useful in the calculation of spheromak equilibria--for which large aspect ratio approximations cannot be justified.

Until this past year, all 2-D axisymmetric toroidal MHD equilibrium calculations done at Illinois have employed various sections of the PEST package.⁽³⁾ Although it can accurately reproduce plasma equilibria for arbitrary aspect ratio and arbitrary current density distribution, the PEST code requires iteration on a 2-D finite difference grid, and is too time consuming (computationally) to employ as the equilibrium-solver module of our 1 1/2-D code. Consequently, we have developed an integral transform technique for solving the Grad-Shafranov equation in arbitrary aspect ratio. In this scheme, surfaces of constant poloidal magnetic flux $\psi(r,z)$ are obtained by i) multiplying the Grad-Shafranov terms by a Bessel function kernel of the form $J_1(\xi r)$; ii) integrating with respect to radial component r over a finite

Fig. 1. Basic Flow Chart for 1 1/2-D Code.

interval encompassing the applied "vertical" field; iii) solving the 1-D second order (in z) differential equations for $\bar{\psi}(\xi, z)$; and iv) applying the appropriate Fourier-Bessel inversion theorem. This technique gives fully analytical Fourier-Bessel series solutions for $\psi(r, z)$ when a toroidal current distribution $J_\phi(r, z)$ is specified and semi-analytical, truncated Fourier-Bessel series approximations when $p(\psi)$ and $g(\psi)$ are specified. One of the important goals achieved during the past year has been the successful development and testing of this computational technique. Preliminary results demonstrating its equilibrium solving capabilities were presented at a computational session of the ANS San Francisco meeting.⁽⁴⁾ A typical equilibrium comparison calculation is shown in Fig. 2.

Utilizing a formulation similar to that of Grad⁽⁵⁾, we have developed a set of transport equations which have adiabatic, ψ -dependent variables and utilize poloidal flux and time as dependent variables. This formulation was chosen due to the peculiar nature of the RFP and spheromak field configurations in which poloidal flux is a monotonically increasing function of minor radius, while the toroidal flux may be double-valued or exceedingly flattened in the edge regions of the plasma.

Our transport scheme is based on an averaging of the moment equations for particle and energy balance over a poloidal flux surface. This procedure leads to moment equations in terms of adiabatic, ψ -dependent variables. These transport equations, in their final form, are a set of coupled parabolic partial differential equations. To solve them, we set up a finite-difference 4x4 block tridiagonal matrix system with "backwards" time centering. The physical processes that are modeled in this formulation include ohmic heating of the electrons and ions, collisional energy transfer between ions and electrons, electron thermal conductivity, adiabatic heating due to a time-dependent applied field, and particle and heat fluxes due to both classical and arbitrary aspect ratio Pfirsch-Schluter collisional transport coefficients. Impurity radiation cooling, charge

Figure 2. Comparison: ITT and Analytical Hill's Vortex.
 Matchup of poloidal flux surface is fairly good.

exchange loss of ion energy, enhanced transport, and auxiliary heating schemes can also be modeled within the framework of this code. Comparisons with a similar 1 1/2-D transport formulation being developed at Livermore for the modeling of Beta-II⁽⁶⁾ are presently underway.

Some preliminary results for the progression (in time) of flux surfaces for a typical Spheromak are illustrated in Fig. 3. Details of our 1 1/2-D formulation, the comparison calculations, and the Spheromak heating conclusions will be made available in the form of a PhD thesis⁽⁷⁾.

II.2.2 Inductive Energy Extraction From a Compact Torus Plasma

During the past year, we have developed a model for estimating the efficiency of the energy conversion associated with the inductive energy extraction from a compact torus plasma. Important relationships accruing from this study included expressions for the energy recovery efficiency as a function of the expansion chamber size and its magnetic field decrement. It was concluded from these relationships that conversion efficiencies ~80-90% can be obtained in a reasonably-sized expansion chamber. Questions relating to the tradeoffs between expansion chamber power surge and axial plasma velocity were also addressed.

An expression for the ideal inductive energy recovery efficiency (i.e., no losses in the pickup coils, walls, etc.) as a function of an expansion parameter, $Q_{ex} \equiv (\text{expanded plasma radius})/(\text{burn chamber plasma radius})$ is plotted in Fig. 4. Estimates for the ideal efficiency as a function of expansion chamber length are plotted in Figure 5 for burn chamber lengths of 20 and 50m and burn chamber radii of 1-5m. Other notable results of this study are summarized in Reference 8. These results and conclusions have been useful in our analysis of the spheromak reactor⁽²⁾ and, also, in the LLL-PGE analysis of the Moving Ring reactor concept⁽⁹⁾. Several basic relationships developed during the course of our study have also been useful in a

a) $t = 25 \mu s$

b) $t = 32 \mu s$

c) $t = 50 \mu s$

d) $t = 84 \mu s$

Fig. 3. Time Evolution of Poloidal Flux Surfaces

- Note: 1) Loss of separatrix-enclosed flux.
 2) Shrinkage of separatrix.
 3) Inward movement of magnetic axis.

Conclusion : Preliminary results appear to be qualitatively reasonable.

Figure 4. Ideal Recovery Efficiency as a Function of an Expansion Parameter, $Q_{\text{ex}} \equiv r_p/r_{p_0}$.

Figure 5. Ideal Efficiency vs. Axial Position (m) for $l_{DC} = 20$ and 50 m and $r_{bc} = 1, 2, 3$, and 5 m (Note: C_L indicates the burn chamber center).

complimentary theoretical study in which the question of mutual inductance coupling in the energy extraction process is analyzed in detail⁽¹⁰⁾.

II.2.3 Gyroscopic Stabilization of the Spheromak Plasma

In the course of recent experimental and theoretical investigations of the spheromak, it has become clear that tilting and shifting instabilities will be two of the formidable obstacles that must be overcome for the successful use of the spheromak concept. Although tilting and shifting can be individually stabilized by proper equilibrium field shaping, it can be demonstrated that the two unstable field index regimes overlap and that field shaping alone is not sufficient⁽¹¹⁾. Near-term remedies for this situation require the use of close-fitting conducting shells⁽¹²⁾, or close-fitting "figure-eight" coils⁽¹³⁾. Unfortunately, while these stabilization techniques may be useful in the S-1 experiment, they detract considerably from the engineering attractiveness of the spheromak reactor scenario and do not appear to be valid "long-run" solutions to the tilt/shift stability problem.

During the past year, we have developed a theory of gyroscopic stabilization of the Spheromak plasma⁽¹⁴⁾. Such a stabilization technique could be viewed as a "long-term" solution that does not detract from the engineering attractiveness of the spheromak. We have analyzed the stabilizing influence of a gentle toroidal plasma rotation ($1\% < V_{\text{rot}}/V_{\text{LH}} < 10\%$) via both force balance and energy principle techniques. Preliminary results indicate that the gyroscopic stabilization technique might be a reasonably successful method. In addition, recent experimental evidence indicates that a gentle bulk rotation may already exist in the present PS1a and c experiments⁽¹⁵⁾ and its possible connection with gyroscopic stabilization has been pointed out in Reference 14. A description of some of our stability analysis techniques, rotational control ideas, and a proposal for future work in this area are presented in our renewal proposal for the 1983 fiscal year.

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II.3 COLLISIONAL PARTICLE TRANSPORT IN AN F RTP

The transport of particles in F RTP devices is still not well understood. This is a key area, because scale-up studies for reactors must proceed with a reliable collisional transport model. In addition, the collisional particle transport represents the upper limit of the particle confinement, i.e., the most favorable situation one can achieve for particle confinement, in a field-reversed (FR) device. This provides direction for reactor design of FR devices where our steady-state calculations give important information for long time-scale operation.

As in the case of the loss cone in velocity space for a simple mirror, an unconfined region in (H, P_θ) space, or equivalently in velocity space, is expected to have significant effects on plasma behavior in field-reversed theta pinches (F RTPs). In present work, we adopt a Fokker-Planck kinetic approach to study the effect of this loss region on collisional particle transport in F RTPs. Although the original form of particle confinement criterion is given in (H, P_θ) space, (1) for the present application it is reformulated in (r, v) phase space.

II.3.1 Approach

First, we solved the problem in 4-dimensional phase space $(r, \theta, v_r, v_\theta)$:

Assuming

- (1) axisymmetric system, i.e. $\partial/\partial\theta=0$,
- (2) $L \gg R$, where L and R are the length and radius of the system, respectively, i.e., $\partial/\partial Z=0$,
- (3) $v_z = \text{constant}$, i.e., $\partial/\partial v_z=0$,
- (4) static magnetic field, i.e., $\underline{B}=B(r)$,

- (5) no macroscopic electric field, i.e., $E=0$,
- (6) external source to maintain a steady state,
- (7) fast loss of unconfined particles,

We then solve the full Fokker-Planck kinetic equation, i.e., Boltzmann equation with Fokker-Planck collisional term, for the ion distribution function $f_i(r, v_r, v_\theta)$. The approach used invokes transformations to obtain separation of variables. Procedures are as follows.

- (1) Shift origin of (v_r, v_θ) coordinates from $(0,0)$ to $(0, -qB_0 r/mc)$ and change variable v to $v/(q^2 B_0 (B_0 r^2/2\psi)/m^2 c^2)^{1/2}$. These transformations enable us to apply the separation-of-variables technique to the kinetic equation which is not separable in the usual (r, v) variables.
- (2) Use a model velocity distribution (modified Rigid Rotor) which satisfies the boundary condition in velocity space set by the particle confinement criterion to evaluate the velocity integrations of Rosenbluth potentials. (Later, the assumed model velocity distribution is tested for validity)
- (3) Use separation of variables to divide the kinetic equation into a set of simpler equations.
- (4) Solve the separated equations.

II.3.2 Results

Important results obtained from present model include the density profile, temperature profile, velocity distribution, scaling laws for local and averaged particle-confinement times, and source profile. All of these, except velocity distribution, are solved analytically. For example, the

explicit and analytic expressions are obtained for the closed-field region of a Hill's vortex magnetic configuration. The radial density profile obtained for this case, shown in Fig. 1, compares favorably with the experimental measurements,⁽²⁾ shown in Fig. 2. The density near axis ($r=0$) in Fig. 1 is smaller than the experimental value. This is understandable because of the assumption of fast loss in our model. The velocity distribution as a function of the dimensionless velocity variable v , which is normalized by $(q^2 B_0 (B_0 r^2 / 2\psi) / m^2 c^2)^{1/2}$, is shown in Fig. 3. In Fig. 3, the calculated velocity distribution $V(v)$ fluctuates due to numerical difficulties, yet compares favorably with the assumed velocity distribution $V_0(v)$. This implies that the simple analytical velocity distribution $V_0(v)$ is a good assumption, and can be used for further reactor studies, which are being carried on.

We then extend this approach to 6-dimensional phase space ($r, \theta, Z, v_r, v_\theta, v_z$). By employing suitable approximations, we have been able to complete this extension successfully. As in the case of 4-dimensional phase space, most of the results are analytic. Fig. 4 shows the density profile from our model, which compares favorably with the experimental radial density profile, Fig. 2, and experimental axial density profile,⁽²⁾ Fig. 5. The fundamental mode of velocity distribution, $V_1(v)$, and the assumed model velocity distribution, $V_0(v)$, are compared in Figs. 6 and 7. These comparisons again show that the model velocity distribution $V_0(v)$ is a very good assumption. Fig. 8 shows the first four higher modes of velocity distribution.

In conclusion, we have solved the steady-state Fokker-Planck kinetic equation in both 4-dimensional phase space and 6-dimensional phase space.

Figure 1. Radial density (in arbitrary unit) profile in the closed-field region of a Hill's vortex magnetic configuration from present model.

Figure 2. Radial density profile from experimental results.

Figure 3. Velocity distribution from present model.
 $V_0(v)$: the assumed model velocity distribution;
 $V(v)$: the calculated velocity distribution.
 v : the dimensionless velocity variable defined in the context.

Figure 4. Density profile in the closed-field region of a Hills vortex magnetic configuration from present model.

Figure 5. Axial density profile from experimental results.²

Figure 6. Velocity distributions. $V_0(v)$: the assumed model velocity distribution. $V_1(v)$: the fundamental model of the calculated velocity distribution. v : the dimensionless velocity variable.

Figure 7. Comparison between $4\pi v^2 V_0(v)$ and $4\pi v^2 V_1(v)$.
 The normalization, $\int_0^1 4\pi v^2 V_0 dv = \int_0^1 4\pi v^2 V_1 dv = 1$,
 has been made.

Figure 8. The first four higher modes of velocity distribution.

Most of the results are analytical. The assumed model velocity distribution is verified a good approximation in both cases. Density profiles are obtained analytically. The simple analytical expression for particle distribution function, which characterizes the system, is obtained, and should be very useful for further reactor studies of axisymmetric field-reversed devices, especially for scoping studies.

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II.4 ELECTRON CURRENTS IN FRM START-UP

A one-dimensional hybrid code FROST (Field Reversal One-dimensional Start-up code) has been developed to model the effects of electron currents on FRM start-up. Although "line-tying" eliminates electron currents on open field lines, electrons may heat up on closed field lines, increasing their diamagnetic current and decreasing ion drag. Also, charge separation between large orbit ions and fluid electrons produce a radial E-field on closed field lines, driving an $E_r \times B$ electron current. Thirdly, fluid equations indicate that electron currents can prevent the transition from open to closed field lines unless an ion impurity is present (Ohkawa currents^(1,2)) or azimuthal symmetry is broken ("magnetic pumping"⁽³⁾). Finally, once the field lines close, the Ohkawa current or "magnetic pumping" must be present for the degree of reversal to increase. If not, electrons at the null will be accelerated by collisions with ions, causing the net current there to vanish and the reversal to decay.

The goals of our study are to quantify the above theoretical predictions, to identify ways they may be minimized in an FRM experiment or reactor, and to find their scaling with increased plasma size. This work is in an intermediate stage; FROST, without an impurity species or an electrostatic field, has just become operational. A description of this hybrid code and its component fluid and ion sections follows, along with some preliminary results.

II.4.1 Code Description and Results

The plasma is modeled as azimuthally symmetric, with no axial dependence. Only the axial component of the magnetic field is assumed to be non-zero. Electrons and low energy ions are modeled using two-fluid equations, while higher energy ions are treated with a

numerical multi-group kinetic model. Each of these models have been developed separately; they were recently merged to form the hybrid code FROST.

The fluid code describes the electron and low energy ion dynamics using equations of continuity, energy and flux conservation. These are coupled to a time differentiated Grad-Shafranov equation which includes the non-fluid ion current. Flux coordinates are used, and the above equations are solved simultaneously, in time, with finite difference techniques. Because of high electron conductivity along field lines, the electron temperature and density are set equal on the two radii of a closed flux surface.

The fluid equations are at present collisionless, with no electrostatic fields. Collisions between the high energy ions and the fluid species will be reflected in the fluid model by Oct. 1982. Still later, the radial electrostatic field will be added. Operation and testing of the fluid code has been described in several past reports^(4,5,6) and so is not detailed here.

Because the ratio of plasma size to ion gyroradius is not large in FRM (2 to 20), the fluid equations are not valid for ions deposited by the neutral beam. Instead, the ion code models the high energy ion distribution with a numerical multi-group method in energy and canonical angular momentum. This distribution evolves due to collisions with electrons, the inductive electric field, and the neutral beam source. For simplicity the ions are assumed to have no axial velocity. The ion current and density profiles are constructed from ion orbits corresponding to each group. Since the ion gyroperiod is much shorter than other relevant timescales ($\tau_{\text{start-up}}$, τ_{ie}), orbit averaging is used. Amperes Law couples the ion current and the

magnetic field. Eulerian coordinates are used. The background is treated as a static target; electron currents are not modeled.

As a typical example of the operation of the ion code, we considered an 80-keV, 5 A/cm "pencil" neutral beam injected 10 cm off-axis into a 30-cm radius target plasma. The target density was uniform at $3.0 \times 10^{13} \text{ cm}^{-3}$, and the temperature was 100 eV. The vacuum magnetic field was 6 kG. A timestep of $8.7 \mu\text{sec}$ (≈ 250 gyroperiods) was used.

Figure 1 shows the reversal factor $RF \equiv B_z(r=0)/B_{\text{vac}}$, as a function of time. Incipient reversal occurs after 3.6 msec (about 420 timesteps). Figure 2 shows a plot of the normalized magnetic flux versus normalized radius for $t = 0$ ($\psi = \frac{1}{2} r^2$) and for $t = 3.83 \text{ msec}$, shortly after reversal. As expected, the flux $\psi(r)$ is flattened near the axis in the latter graph. Also shown in Figure 2 is a plot of the ion current versus normalized radius. This curve is peaked near the injection radius, as expected.

The hybrid code FROST, formed by the merger of the fluid and ion codes, is now operational. Magnetic coupling between the two models is provided by a time differentiated 1-D Grad-Shafranov equation which includes the non-fluid ion current:

$$\frac{\partial}{\partial \psi} \left[\frac{d}{dt} \left(\frac{B^2}{2\mu_0} \right) + \frac{dp_{\text{fluid}}}{dt} \right] = - \frac{d}{dt} \left[\frac{j_{\text{ring}}}{r} \right]$$

where j_{ring} is the high energy ion current. This equation resides in the fluid section of the code. Ampere's Law, formally used in the ion code, has been discarded. The fluid equations are collisionless at present. Therefore, flux is conserved by the electrons, and it is impossible to make the transition from open to

Figure 1. Reversal Factor = $B_z(r = 0)/B_{vac}$ Versus Time for a typical application of the kinetic ion code.

Figure 2. Initial Flux Function (parabolic) and Flux Function and Azimuthal Ion Current When $t = 3.83$ msec, from operation of the kinetic ion code.

closed field lines. This addition of collisions and an impurity species will allow this transition to occur.

Two new features have been added to FROST. These are:

- 1) Ionization and charge exchange of the neutral beam by high energy ions in the plasma.
- 2) A variable timestep adjusted to be less than a small fraction ($\sim .10$) of the average time for an ion to slow down through one energy bin.

The importance of these changes is made obvious by considering the particle composition of a typical start-up plasma. While the initial target is a low density plasma stream, the number of high energy ions from the beam soon becomes dominant. The associated rise in the electron density increases the slowing down rates as the buildup progresses.

As an example of operation of FROST, we have simulated tangential injection in 2XIIB. A 1.5 A/cm neutral beam was aimed 4 cm off-axis into a plasma stream of density $10^{13}/\text{cm}^3$, temperature 0.1 keV and radius 6 cm. The vacuum magnetic field was 4.65 kG and the wall radius was 12 cm. After 1.5 msec the average ion density was $6.6 \times 10^{13}/\text{cm}^3$, while the center line magnetic field was depressed by 30%. The circulating ion current at this time was 1.06×10^3 A/cm. Figure 3 is a plot of the reversal factor and the average ion density versus time. Longer operation of the code is planned to allow direct comparison with reversal experiments in 2XIIB, where peak values were reached after 4 msec of injection.

Figure 3. A 2XIIB simulation with FROST. Longer operation is needed to simulate the approach to reversal.

II.4.2 Conclusion

The fluid and ion codes have been merged to form the hybrid code FROST, and initial results simulating 2XIIB start-up have been produced. Work to be done includes programming of transport in the fluid equations and a radial electric field. Collisions between the high energy ions and the fluid species, including impurities added to supply the Ohkawa current, should be reflected in the fluid model by October, 1982. After this, the radial electrostatic fluid will be included. With these two features included, FROST should be valuable in revealing the role of electron currents in the creation of the reversed field geometry using neutral beams.

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II.5 RFP ANOMALOUS TRANSPORT SCOPING STUDIES

The purpose of the present work has been to develop a phenomenological anomalous transport model for RFP's (based on a braided magnetic field picture) which can be tested against current ZT-40 experimental results. Understanding of the transport behavior is crucial for (1) determining how the relatively short experimental confinement times will scale to reactor regimes for reactor feasibility and design studies, and (2) suggesting ways to eliminate or reduce the anomalous effect. Along with the transport model development (and necessary for evaluating the model), is the goal of constructing a computer code package which will reproduce ZT-40 results. This is being accomplished by modifying RFPBRN⁽¹⁾, a 1-D, two fluid, Lagrangian mesh, transport and stability code, developed by R. Nebel and R. Moses of LANL. This code uses a circular cylinder, quasi-static approximation and follows the time evolution of temperature, density and magnetic field profiles for a reactor-grade RFP plasma. This code is written in Fortran and is available through the MFE network on the Cray machine.

II.5.1 Status of Work (See Table 1)

1) ZT-40 Code Package: A thorough working knowledge of RFPBRN was achieved. The code was rewritten to operate in the convective frame ($\vec{E} \times \vec{B} = 0$). This entailed writing diffusive terms explicitly which permits a clearer description of transport effects along with the easy application of a diffusive edge condition necessary for describing the ZT-40 pumpout phenomenon. (See Table 2)

Table 1. Milestone Charts

CODE DEVELOPMENT:

THEORY DEVELOPMENT AND TESTING:

• Table 2. RFPBRN Code Development

RFPBRN = 1-D, time dependent, MHD, RFP transport and stability code.

<u>OLD FEATURES</u>	<u>NEW FEATURES</u>	<u>COMMENTS</u>
Lagrangian Frame ($V=0$)	Convective Frame ($\vec{E} \times \vec{B} = 0$)	allows a more explicit description of transport terms.
Ideal MHD Stability Package	Neglected	shortens transport run time.
Parameters Described: $n_i, n_\alpha, T_e, T_i, B_z, B_\theta$	$n_i, T_e, T_i, B_z, B_\theta$	α neglected for modeling ZT-40, shortens run time.
Boundary Conditions:		
no particle loss from plasma	particles & their energy can be lost	describes particle pump-out phenomenon.
$\vec{B} _{\text{Lagr.}} = 0$	$B_z(a)$ measured	match experiment.
fixed plasma radius	Compressive ($\vec{E}_\perp(a) \times \vec{B}$) b.c., $E_\perp(a)$ measured	match experiment, explain $I_z(t)$.
fixed no. of mesh pts.	variable no. of mesh points	allows plasma description near the wall.
_____	Neutral Reflux	explains $\langle \dot{n} \rangle$.
_____	Anomalous $\eta_{ }$	explains $I_z(t)$.
_____	Anomalous OH to ions	may help explain $T_i \geq T_e$.
_____	Simple Dynamo model	needed to explain $E_\perp(a,t), \Phi_z(t)$.

A preliminary parameter study (and comparison to ZT-40) was completed using Bohm transport coefficients. This work identified several important non-transport effects that must be included to predict ZT-40 behavior. One such effect is the importance of neutral reflux from the wall to explain the behavior of the average electron density over the last few msec of a ZT-40 shot. (See Fig. 1) A simple neutral reflux model was added to the code. We also added the effect of a perpendicular electric field at the wall. This field and the resultant compressive Poynting's vector flux is necessary to explain the observed increase in toroidal current. (See Fig. 2) Also, we observed the need for an anomalous parallel resistivity that increases with time (a behavior first estimated by Ron Moses of Los Alamos) to produce the rapid drop in toroidal current at the end of the shot. A fourth effect, the "dynamo," is needed to explain the maintained average toroidal magnetic flux and the value of the fields at the wall. This is related to the anomalous parallel resistivity and is presently being examined.

2) Anomalous Transport Theory: Of the three basic models (described in the original work proposal⁽²⁾ from which phenomenological transport coefficients can be derived, one, the fluctuating E-M field model, has been discarded because the characteristic fluctuating frequencies and step sizes of the resistive g-modes are too small to produce significant transport. Another model, the tangled discharge model, however, has been developed and seems to have the correct features for explaining ZT-40 transport, including pumpout and faster-than-Bohm electron thermal energy losses.

Contacts

This work is being conducted in close association with R. Nebel and R. Moses of the plasma theory group of LANL. K. Werley (UI) will be spending the upcoming summer at Los Alamos, to aid in the continuation of this work.

Figure 1. Average Density vs. Time. Curves are labeled to illustrate the difference between anomalous and classical predictions, and whether the compressive (C) boundary condition was used or not (NC). The dotted line represents exponential decay. ($D_{\text{anom}} = .05 D_{\text{Bohm}} \sim 75 D_{\text{clas}}$).

Figure 2. Toroidal Current vs. Time. Labels on curves specify the value of parallel resistivity and whether the compressive boundary condition was used. ($\eta_{\text{anom}} = 2 \times 10^{-4} \eta_{\text{Bohm}} \sim \frac{1}{2}$ to $10 \times \eta_{\text{clas}}$).

We will continue to better understand the role of the dynamo, eventually including a model for the dynamo and for the related parallel anomalous resistivity into the code. We will then introduce the anomalous thermal conductivities, diffusion coefficients and perpendicular resistivity from the tangled discharge model to the code and carefully debug the code, checking on energy and particle conservation, and identifying any numerical problems. Code predictions will then be compared to ZT-40 results and the model will then be evaluated.

II.5.2 Conclusion

This work is the first attempt at a complete phenomenological transport model for the RFP. At the very least, this work will provide a standard for comparing other models against, and it will suggest areas needing more detailed theoretical development.

The lack of detailed experimental profile information will make it difficult to verify the validity of a transport model. However, if our model accurately predicts the known data, it should be possible to suggest experimental parameter studies to help check on the model or to reduce the anomalous losses. It will also be worthwhile to compare our model to those used in present reactor designs.

II.5.3 Publications

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III. FUSION PRODUCTS EFFECTS

III.1 RADIAL ELECTRIC FIELDS FOR IMPROVED TOKAMAK PERFORMANCE

III.1.1 Electric Field Effects

The influence of externally-imposed radial electric fields on the fusion energy output, energy multiplication, and alpha-particle ash build-up in a TFTR-sized, fusing tokamak plasma was explored. In an idealized tokamak plasma, an externally-imposed radial electric field leads to plasma rotation, but no charge current flows across the magnetic fields. However, a realistically-low neutral density profile generates a non-zero cross-field conductivity and the species dependence of this conductivity allows the electric field to selectively alter radial particle transport.

The plasma burn time until reaching a given ash fraction (the "time-to-buildup") is used as a measure of plasma performance. Because the fuel species has a higher cross-field conductivity than the ash, imposition of an inward-directed field enhances fuel confinement more than ash confinement. An outward-directed field, on the other hand, drastically reduces the fuel density, but has little effect on the ash density. Therefore, the time-to-buildup is increased by an inward-directed field, but decreased by an outward-directed field.

To evaluate the spacial dependence of radial electric field effects on plasma performance, three charge-deposition profiles (shallow, moderately deep, and deep into the plasma) were taken to generate the field. Electric-field-induced transport is then self-consistently modeled⁽¹⁾. That is, the charge deposited inside the plasma is balanced by the current flow which results. The introduction of about one amp of charge current near the plasma edge allows nearly steady-state ($t > 100$ sec) operation with energy multiplication, Q , greater than unity (vs. $Q \rightarrow 0$ within tens of seconds without an imposed electric field). Deep penetration of the plasma by the charge source

leads to even greater improvements in Q ($Q > 4$ at $t > 100$ sec). Shown in Fig. 1 are typical enhancements in Q for a TFTR sized plasma.

Possible mechanisms for the injection of charge into the plasma were also explored. Since straightforward schemes (e.g., direct injection of electrons across the magnetic field) require excessive amounts of power, new and original techniques are needed. Two such schemes (Fat Banana Carrier Injection and Ripple-Trap Carrier Injection) were analyzed and found to be promising methods of low-energy-cost charge injection.

III.1.2 Fat Banana Carrier Injection

For Fat Banana Carrier Injection (FBCI) the width of a banana orbit (which is the largest penetration depth normally anticipable for a charged particle) is proportional to the square root of the mass. Therefore, high injection power in relation to electron injection may be avoided by increasing the mass of the injected particle. Injection of fuel negative ions along fat banana orbits also fails, however, to produce the desired electric field. For background D-T densities $\sim 10^{13} \text{ cm}^{-3}$, such as are expected near the edge of the TFTR plasma, the cross section for the reaction $D^- + D^+ \rightarrow D^0 + e + D^+$ is much shorter than the banana orbit time for energies of practical interest⁽²⁾. That is, the D^- penetrates only a small fraction of a banana width before losing its electron and transiting a (probably short) chord of the plasma cross section.

Therefore, a charge-carrier is desired which will delay giving up its electron to the plasma longer than a fuel negative ion, which will not lead to plasma contamination, and which requires minimal power throughout. Consider then the following sequence of events: ion X^{n+} (with species X to be chosen later) is injected into the plasma from outside, along a fat banana orbit. (See Fig. 2). At the deepest penetration of the plasma (i.e. after one half-orbit) the ion loses one electron and becomes $X^{(n+1)+}$. Assuming ionization

Figure 1. Time-integrated energy multiplication factor (Q = fusion energy output/gross energy input to maintain centerline temperature) vs. time for $\tau_p/\tau_E = 10$ and $\tau_p/\tau_E = 20$. D^- injection is assumed and electric-field-induced transport is included. Curve symbols correspond to shallow charge source (dotted curve), moderate penetration (dashed curve), and deep penetration (solid curve).

Figure 2. Fat Banana Carrier Injection. Charge carrier species X^{n+} is injected behind the limiter, deposits an electron inside the plasma and is recovered by the limiter. The net effect is the generation of an inward-directed radial electric field in the plasma. Note that the distance $n\Delta/(n+1)$ is the maximum possible penetration distance with FBCI. In practice, the collection area will need to be finite to account for spreads in the energy and pitch-angle of the carrier (X^{n+}) beam.

occurs due to electron impact, little change in the ion pitch angle should result. The banana width (inversely proportional to charge) decreases to $n/(n+1)$ times its previous value, but the ion will remain trapped. If the original injection point was situated properly outside the plasma, this new banana orbit barely skims the surface of the plasma at its largest radial excursion. At this point, the carrier ion is collected, leaving the plasma with a net single electronic charge.

Figure 2 shows the carrier ion being collected by a limiter. It should be noted, however, that this injection scheme should also be adaptable to a plasma with a divertor. The pitch angle at the most-outboard point of a typical fat banana orbit lies between 30 and 40 degrees⁽³⁾. Unless the divertor throat is quite strong, the carrier ion will not reflect off this magnetic mirror back into the plasma.

There are four requirements for successful Fat Banana Carrier Injection as described above. First, the bounce time of the fat banana orbit must be much less than the slowing down time of the injected ion. Otherwise, collisions with the ion may change its orbit such that it misses the collector at the end of its journey and becomes an impurity in the plasma. Second, the ionization time for the reaction $X^{n+} \rightarrow X^{(n+1)+}$ must be approximately half the bounce time in order to deposit the electron as far into the plasma as possible. Third, the ionization time for the reaction $X^{(n+1)+} \rightarrow X^{(n+2)+}$ must be long compared to half the bounce time to insure recovery of the carrier ion. Finally, the difference between the banana widths at charge state $n+$ and $(n+1)+$ must be large enough to give reasonable penetration of the plasma. Analysis of the preceding scenario⁽¹⁾ yields the banana width (Δ) as a function of mass (m), charge state (z) and ionization energy (I_n) for TFTR as

$$\Delta = 16.9 \sqrt{\frac{m}{zI_n}} \quad (1)$$

Prospective carrier species may be examined using Eq. 1 given an initial charge state and the carrier mass. The results for some sample carriers are shown in Table I. Highly-charged carriers like C^{3+} allow lower-energy injection, and the penetration depth is a higher fraction of the initial banana width (relative to singly-charged carriers). However, the banana widths are smaller, leading to net shallower injection. Higher-mass carriers like Pb have good penetration depths, but their bounce time is longer, increasing the likelihood that the carrier will overionize in the plasma, decrease its banana width, and miss the collector. Within these limits, however, there are several promising candidates, such as Mg^+ and Ca^+ . It should be further noted that refinements in the approximations above may significantly improve or diminish the viability of any given carrier species.

FBCI, then, appears to be an attractive technique for the introduction of charge into the plasma, across the magnetic field. If the carrier beam could be created at 100% efficiency and the energy in the recovered beam could be converted at 100% efficiency, then virtually no power would be required to impose a radial electric field on a plasma. For a more realistic efficiency analysis, label the efficiency with which the original carrier beam is created η_G , the efficiency with which the waste energy from this process is recovered η_{R1} , and the efficiency with which the energy of the returning carrier beam is recovered η_{R2} . Then, if the carrier-beam power required is P_{inj} , FBCI will cost the tokamak power balance

$$P_{inj} \left[\frac{1}{\eta_G} - \frac{\eta_{R1}}{\eta_G} (1 - \eta_G) - \eta_{R2} \right] \quad (2)$$

Assuming the quantity in the brackets in Eq. 2 can be limited to values around unity, only 3.7 kW of magnesium carriers (0.1 amp at 37 keV) would be needed to generate a 10 kV/m (maximum) radial electric field in TFTR.

TABLE I

Carrier	n	$\Delta(\text{cm})$ (Eq. 1)	penetration depth $\Delta(n/n+1)$
${}^4_9\text{Be}^+$	1	11.9	6.0
${}^{10}_5\text{B}^{2+}$	2	6.1	4.1
${}^{12}_6\text{C}^{3+}$	3	4.2	3.2
${}^{12}_{24}\text{Mg}^+$	1	21.4	10.2
${}^{20}_{40}\text{Ca}^+$	1	31.1	15.5
${}^{82}_{207}\text{Pb}^+$	1	62.8	31.4

With a centerline plasma density of $5 \times 10^{13} \text{ cm}^{-3}$, a density profile proportional to $1 - (r/a)^2$ and an energy confinement time of one second, this represents less than a 2% drain on the fusion power output at 8 keV. FBCI is, then, an energetically inexpensive method of improving tokamak plasma performance.

III.1.3 References

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III.2 FUSION PRODUCT DRIVEN MICROINSTABILITIES

III.2.1 Introduction

Direct coupling of alpha particle energy to plasma ions through quasi-linear wave-particle interaction has long been viewed⁽¹⁻⁴⁾ as an attractive means of enhancing the low classical alpha-ion energy split. One well known paper⁽⁵⁾ estimates an order of magnitude lower $n\tau_E$ assuming that all alpha energy is coupled to the plasma ions. Similarly, alpha particle interaction with shear Alfvén waves can potentially produce anomalous energy loss. It is the purpose of this work to quantify these effects at the plasma level and at the system level. Present status of these efforts is outlined below. Progress reported for this period includes:

- a. scoping of system level effects for both anomalous heating and transport of alpha particles,
- b. theoretical development of the quasi-linear diffusion equation,
- c. numerical solution of the thermonuclear driven fast magneto-sonic (FMS) wave heating of an ETF/FED like tokamak plasma,
- d. identification in n_e - T_i space of threshold contours for the FMS wave, i.e., stability regions.

III.2.2 System Level Scoping Studies

We have previously developed phenomenological models to scope the physics and plasma level effects of alpha driven instability heating⁽⁶⁾ and spatial transport.⁽⁷⁾ These models have been incorporated into the 1 1/2-D tokamak transport code WHIST and runs made for several sets of plasma conditions in a fusion engineering like device (FED). Results from representative runs are reported in Table I. In obtaining these results we start with identical initial conditions, turn on the effect and examine the

Table I

INITIAL PARAMETERS				Toroidal B-field kG	Applied Beam Power MW	FINAL PARAMETERS							
Run	n_e (axial) $10^{13}/\text{cm}^3$	T_i, T_e (axial) keV	Profile			n_e (axial) $10^{13}/\text{cm}^3$	$\langle n_e \rangle$ $10^{13}/\text{cm}^3$	T_i (axial) keV	$\langle T_i \rangle$ keV	$\langle \beta \rangle$ %	Total Power MW	On- Axis Enhance- ment	Net Power Change %
FED N 12	8.0	20.0	Parabolic	44.93	15	4.91	3.2	31.67	15.96	2.36	17.19	1.0	N/A
FED A 12	"	"	"	"	"	6.20	"	29.93	15.58	2.42	20.73	1.24	20.59
FED S 12	"	"	"	"	"	5.60	"	30.59	14.71	2.23	18.49	0.87	7.56
FED N 14	8.0	40.0	Parabolic	44.96	20	5.22	3.2	36.07	19.65	2.97	19.75	1.0	N/A
FED A 14	"	"	"	"	"	6.02	"	34.73	18.78	2.76	24.02	1.32	21.62
FED S 14	"	"	"	"	"	5.92	"	35.82	18.99	2.94	24.09	0.83	21.97
FED N 13	8.0	40.0	Parabolic	44.96	25	5.10	3.2	34.25	18.07	2.71	21.41	1.0	N/A
FED A 13	"	"	"	"	"	5.26	"	37.57	20.59	2.86	22.48	1.30	5.00
FED S 13	"	"	"	"	"	5.25	"	37.90	20.44	3.07	22.40	0.84	4.62
FED N 15	8.0	40.0	Parabolic	44.96	30	5.38	3.2	37.02	20.82	3.07	22.58	1.0	N/A
FED A 15	"	"	"	"	"	5.30	"	39.06	21.99	2.97	23.39	1.29	3.58
FED S 15	"	"	"	"	"	5.35	"	39.60	21.98	3.16	23.69	0.82	4.92

plasma after a real time simulation of 20 sec. Since the effect from anomalous heating will be weak, we have deliberately chosen to operate at high temperatures. For comparison, the effect of anomalous alpha particle transport is reported for the same initial conditions. Nomenclature used for the run category is: FED N = normal run, no anomalous effects; FED A = anomalous alpha heating effect⁽⁶⁾ turned on; FED S = anomalous alpha transport effect⁽⁷⁾ turned on.

We note from Table I that the anomalous alpha heating is indeed a mild effect resulting in a potential enhancement in reactor level power $\leq 20\%$. Examination of the detailed time histories indicates that the alpha pressure is reduced with an increase in the particle density; it is primarily this increase in plasma density (and not the variation of temperature profiles) which improves the device level, i.e., system level, power output. While a system-level enhancement $\sim 20\%$ is potentially interesting we shall show later that significant plasma level heating occurs only for $T_i \sim 5\text{-}10\text{ keV}$. At these temperatures, the increased collisionality compared to temperatures in Table I will result in only a marginal improvement in system level power for a tokamak device. Turning to the shear Alfvén mode we are surprised to note that even for a net decrease in on-axis heating $\sim 15\text{-}20\%$ that an improvement in system power takes place. This surprising result of increased system power in the presence of anomalous alpha spatial transport results from some flattening in alpha pressure with the increase in plasma fuel density.

Finally we comment on the plasma β for the device where the instability may make an effect. From the definition of β we have

$$\beta = \frac{8\pi}{B_0^2} \sum_s n_s T_s = \frac{\sum_s n_s T_s}{\frac{\langle n_i m_i \rangle}{2} \frac{B_0^2}{4\pi \langle n_i m_i \rangle}} = \frac{\sum_s n_s T_s}{\frac{1}{2} \rho v_A^2}$$

where $\rho \approx \langle n_i m_i \rangle$ = plasma mass density

v_A = Alfvén speed.

For $\sum_s n_s T_s = n_0 (T_e + T_i)$

$$\beta = \frac{T_e + T_i}{\frac{1}{2} \langle m_i \rangle v_A^2} = \frac{\sum_s T_s}{\frac{\langle m_i \rangle}{m_\alpha} \frac{v_A^2}{v_0^2} E_{\alpha 0}} = \frac{m_\alpha}{\langle m_i \rangle} \left(\frac{v_0}{v_A} \right)^2 \frac{\sum_s T_s}{E_{\alpha 0}}$$

where m_α = alpha mass

$E_{\alpha 0} = \frac{1}{2} m_\alpha v_0^2$ = initial alpha energy.

So far the last expression follows purely from definitions. From detailed physics studies (cf Figure 2) we require $v_0/v_A \sim 1.5$ for the instability

$$\rightarrow \beta \approx 2.25 \frac{m_\alpha}{\langle m_i \rangle} \frac{\sum_s T_s}{E_{\alpha 0}}$$

Inserting numbers we find $\beta \sim 2.3\%$ for alpha particles in a 10 keV plasma.

To maximize the effect, this β must occur on axis so that we will have an average plasma beta $\sim 1/3 - 1/2$ the axial value: $\langle \beta \rangle \lesssim 1\%$.

III.2.3. Theoretical Physics Development and Numerical Results

Theoretical analysis of the pertinent quasi-linear diffusion and growth rate equations reveals significant structure which can be profitably exploited. In contrast to Mikhailovskii,⁽⁸⁾ who worked in spherical $v-\theta$

coordinates, we analyze the problem in circular cylindrical $v_\perp$ - $v_\parallel$ coordinates. Evaluation of the perpendicular, parallel and cross diffusion coefficients reveals that the perpendicular quasi-linear diffusion coefficient, $D_\perp$, greatly dominates all other terms. Further, since both $D_\perp$ and the instability growth rate are, to leading order, independent of $v_\parallel$ we can actually reduce the problem to 1-D in tokamak geometry.

To find the required equation, we begin with the standard quasi-linear equation⁽⁹⁾ for infinite geometry. Taking $\omega_r \approx k_\perp v_A$, we note a small parameter in $v_\parallel$ to expand the Bessel functions in a Taylor series perform the integration over $k_\perp$ to find:

$$\begin{aligned} \frac{1}{2} v_A v_\perp \left[\frac{m_\alpha}{\pi q_\alpha} \right]^2 \frac{\partial f_\alpha}{\partial \tau} = & \frac{\partial}{\partial v_\perp} D_\perp \frac{\partial f_\alpha}{\partial v_\perp} \\ & + \frac{\partial}{\partial v_\perp} D_x \frac{\partial f_\alpha}{\partial v_\parallel} \\ & + \frac{\partial}{\partial v_\parallel} D_x \frac{\partial f_\alpha}{\partial v_\perp} \\ & + \frac{\partial}{\partial v_\parallel} D_\parallel \frac{\partial f_\alpha}{\partial v_\parallel} \\ & + \text{"sources"} \end{aligned}$$

$$D_\perp \approx k_{LD}^2 v_\perp^2 \sum_L \frac{1}{L \Omega_\alpha} \left\{ \frac{v_A^2}{v_\perp^2} E_r^2 J_L^2 + E_\theta^2 (J_L')^2 \right\}$$

$$D_x \approx \frac{1}{2} k_{LD}^2 v_\perp^2 \sum_L \left\{ \frac{v_A^2}{v_\perp^2} E_r^2 J_L^2 + E_\theta^2 (J_L')^2 \right\}$$

$$D_\parallel \approx \frac{1}{3} k_{LD}^3 v_\perp^3 \sum_L \frac{1}{L \Omega_\alpha} \left\{ \frac{v_A^2}{v_\perp^2} E_r^2 J_L^2 + E_\theta^2 (J_L')^2 \right\}$$

$$J_L = J_L \left(L \frac{v_\perp}{v_A} \right)$$

For the central portion of the plasma, the sum extends over 1 harmonic only to find in the resonance region;

$$D_{\perp} : D_x : D_{\parallel} \approx 1 : \frac{1}{2} \frac{v_A}{v_e} : \frac{1}{3} \frac{v_A^2}{v_e^2} .$$

Since the growth rate (equation 3.8, Ref. 4) is independent of $v_{\parallel}$ we can integrate out the $v_{\parallel}$ dependence in the above quasi-linear diffusion equation. This reduction to 1-D can be rigorously accomplished under the following conditions:

- a. wave propagation is almost entirely perpendicular to the applied magnetic field ($k_{\parallel} \ll k_{\perp}$) due to heavy ion cyclotron damping at $\omega_r \approx n\Omega_{\alpha}$ for finite $k_{\parallel}$, and
- b. neglect of inherently 2-D phenomena such as collisional effects (diffusion and slowing) and loss cones. These approximations should be acceptable for time scales short compared to the Coulomb slowing down time and close to plasma core.

Further, to bound the heating available we follow the time evolution of the birth alpha distribution function: i.e., an isotropic Gaussian shell with an energy spread given by⁽¹⁰⁾ $\Delta E \approx 177\sqrt{T_i}$ keV (T_i in keV).

Using the above formalism we have completed work on the following code packages:

- a. QLD1: A 1-D, time dependent, quasi-linear diffusion code for the birth alpha distribution function or a perpendicularly injected neutral beam.
- b. QLD1P: graphics package for QLD1.
- c. QLD1SVA: a cosmos job control language file to make production runs in n_e - T_i space for a given wave frequency, radial position and applied magnetic field.

The above codes run on either the CRAY or CDC 7600 machines with typical run times ~ 2 -3 minutes for QLD1. Typical run results appear in Figures 1-4.

III.2.4 FMS Threshold Contours

Using the linear growth rate we have also developed a code package ETAT which computes the alpha driven growth rate, γ_α ; the collisionless wave damping from the bulk plasma, γ_α ; and the associated alpha particle threshold fraction $\eta_T = n_\alpha/n_e$ required for the instability. A code package ETATP plots the contours; a typical run may be found in Figure 5 (ETF device at $B_0 = 42$ kG, radial position = 10 cm, $\omega_r \approx 6\Omega_\alpha$). Plotted contours represent $\log_{10}\eta_T$ with the outer boundary of 0.00 defined to be the stability limit. Contours are mapped at unit increments: -1.0 ($\eta_T \sim 10^{-1}$), -2.0 ($\eta_T = 10^{-2}$) through -6.0 ($\eta_T = 10^{-6}$). Also lines of local plasma beta are superimposed upon the contours.

Several remarks are in order. Firstly, it is quite apparent from Figure 5 that only certain regions in the n_e - T_i plane favor this instability. The lowest threshold $\eta_T \sim 10^{-6}$ occurs around $n_e \sim 4.5 \times 10^{13} \text{ cm}^{-3}$ and $T_i \sim 5$ keV. From that point the resonance detunes as T_i increases due to Doppler broadening of the birth alpha distribution function. Along the n_e axis the instability experiences successive tuning/detuning due to variation of the Alfvén speed in the argument to the Bessel function (cf Figure 2, toroidal resonance function).

From production runs with QLD1 we find that the potential enhancement in the classical α -ion energy split is maximized for the lowest threshold fraction, η_T . Values of this enhancement of up to 1.9 have been seen. For the run ETF.DBB in Figure 5, the plasma level effect is strongest for $n_e \sim 4.5 \times 10^{13} \text{ cm}^{-3}$ and $T_i \sim 5$ -15 keV. Thus, if the plasma beta on axis is to be $\sim 2\%$ we conclude that the effect will be a delicate one to exploit due to the narrow operating region for n_e and further that it will only contribute at low $\langle \beta \rangle \lesssim 1\%$.

THERMONUCLEAR DRIVEN
FAST MAGNETOSONIC WAVE HEATING
QLD1 SUMMARY

RUN: ETF.D32A

TOKAMAK PARAMETERS

RMAJOR (CM)	500.00
RMINOR (CM)	120.00
RADIAL POSITION (CM)	5.00
B-FIELD (K-GAUSS)	42.00

(LOCAL PARAMETERS FOLLOW)

PARTICLE DENSITY ($10^{13}/\text{CM}^3$)	4.50
ELECTRON TEMP (KEV)	10.00
ION TEMP (KEV)	10.00
ALFVEN SPEED (10^8 CM/SEC)	8.64
VALPHA / VALFVEN	1.51

INSTABILITY PARAMETERS

LOG 10(THRESHOLD FRACTION)	-5.51
LOG 10(-DAMPING RATE)	2.53
WAVE MODE NUMBER	6.00
PEAK ENHANCEMENT	1.44

Figure 1

1-D PHYSICS PARAMETERS

LEGEND

- = TOROIDAL RESONANCE FUNCTION, R
- △ = RESONANCE FUNCTION DERIVATIVE, RPRIM
- ◇ = V-PARALLEL AVERAGED ALPHA PROFILE
- ▽ = PERPENDICULAR DIFFUSION COEFFICIENT, DPERP

$\text{LOG } 10(\text{MAX. } R) = 1.0983$
 $\text{LOG } 10(\text{MAX. } R_{\text{PRIM}} (\text{CM/SEC})^{*-1}) = -7.2511$
 $\text{LOG } 10(\text{MAX. ALPHA PROFILE}) = -10.4374$
 $\text{LOG } 10(\text{MAX. } D_{\text{PERP}} (\text{CM/SEC})^{*3/\text{SEC}}) = 16.3392$

Figure 2

INSTABILITY HISTORIES

LEGEND

- = WAVE ENERGY DENSITY (ERGS/CM**3)
- △ = ENHANCEMENT OF CLASSICAL ENERGY SPLIT
- ◇ = TOTAL GROWTH RATE (SEC-1)
- ▽ = REL. ERROR: ALPHA NUMBER DENSITY
- × = REL. ERROR: SYSTEM ENERGY

LOG 10(MAX. WAVE ENERGY DENSITY) = 2.1648
 LOG 10(MAX. GAMMA/TOTAL/) = 2.9781

Figure 3

ALPHA PROFILE HISTORIES

LEGEND

- = FALPHA AT 10.0 MILLI-SEC
- △ = FALPHA AT 17.5 MILLI-SEC
- ◇ = FALPHA AT 48.5 MILLI-SEC
- ▽ = FALPHA AT 50.0 MILLI-SEC

LOG 10(VERTICLE SCALE) = -9.1174

Figure 4

LOG 10(THRESHOLD FRACTION)

LEGEND

□ = LOCAL PLASMA BETA = 1 %

○ = LOCAL PLASMA BETA = 3 %

△ = LOCAL PLASMA BETA = 5 %

◇ = LOCAL PLASMA BETA = 10 %

▽ = LOCAL PLASMA BETA = 15 %

Figure 5

III.2.5 Summary

We have completed work on scoping studies at the system level for both alpha driven instability heating and shear Alfvén spatial transport. Cases examined for a FED like device indicate only a mild and generally uninteresting improvement in total power output which occurs for $\langle \beta \rangle \lesssim 1\%$. Theoretical work indicated significant structure in the quasi-linear diffusion equation, linear growth rate and birth alpha distribution function. This structure has been exploited to rigorously reduce the problem to the 1-D time dependent case which treats all alpha particles - not just trapped ones. Subsequent numerical solution of extensive production runs have shown agreement between the plasma level engineering model (6) and the full quasi-linear problem. Threshold contours in the n_e - T_i plane are available which, for the first time, show candidate regions of stability/instability. Heating data have been generated from QLD1 production runs and a subroutine package inserted into the WHIST code to evaluate the heating potentially available. System level cases based upon the heating from the detailed, quasi-linear model are to commence shortly. Since the detailed physics has only confirmed the engineering model of anomalous heating, (6) we do not anticipate that the system level power will be improved by more than a few per cent.

While the thermonuclear driven FMS instability appears to be only mild for tokamaks the same instability can have a significant effect in mirror machines. Since the high energy alphas will be mirror rather than electrostatically confined, a loss cone will be established along the $v_{||}$ axis. Should the fast magnetosonic wave be destabilized by thermonuclear alpha particles, the alphas will be scattered into the loss cone due to quasi-linear $D_{\perp}$ diffusion and escape from the system before thermalizing.

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III.3 STUDIES OF FUSION PRODUCT TRANSPORT IN THE EBT AND FRM INCLUDING NES EFFECTS

III.3.1 Guiding Center Transport in the EBT

Initial estimates of alpha particle confinement and heating in EBT reactors based only on vacuum magnetic field considerations indicate poor confinement of these high-energy particles due to the poor centering of their μ -J surfaces.⁽¹⁾ Estimates show that due to the poor confinement of passing particles and additional losses during thermalization, as little as 60% of the total alpha energy may actually be deposited in the plasma. In order to evaluate means of enhancing alpha confinement, the Monte Carlo transport code, MCFRM, has been modified to represent EBTR (Elmo Bumpy Torus Reactor) geometry. The use of this transport code will serve as a stochastic check to the confinement calculations presented in the EBTR design study as well as introducing slowing-down dynamics into the problem of retaining fusion product energy.

The code MCEBTR tracks individual alpha particle trajectories in representative EBTR magnetic geometries. The toroidal magnetic field was modelled through the use of "current-loop" model developed at Oak Ridge.⁽²⁾ Because the magnetic field is steady-state, and since the gyroradius for the typical alpha particle is generally smaller than the scale length of interest ($B/|\nabla B|$), the guiding-center approximation is valid. A guiding center model based on the approximate invariants of the motion, μ and J , is used to determine fusion product motion.

To accurately model fusion product dynamics in EBTR and to keep computational time to a minimum a "unit cell" is utilized consisting of two adjacent toroidal-field coils as well as the aspect-ratio-enhancement (ARE) coils. The effect of the electron rings, which are needed for global stability, is not presently included. For good statistics, a sample size of

about 1000 particle histories is needed to accurately model the fusion product energetics. These particles, which are equally weighted, are initialized to reproduce the correct spatial birth distribution of the alpha particles. The initialization is isotropic in velocity space, and all test particles start on the midplane of the unit cell. In order to determine the "passing fraction," i.e., the fraction of newly born alpha particles whose pitch angles are too small to be trapped by the magnetic-mirror effect, a fourth-order Runge-Kutta integrator is used to trace magnetic field lines from the birth points on the midplane to the corresponding radial position at the coil plane. This integration allows the off-axis mirror ratio to be evaluated and each alpha particle to be classified as either "trapped" or "passing".

In the most general form, the guiding center equation of motion can be written (egs units) as:

$$\frac{d\vec{r}}{dt} = v_{||} \frac{\vec{B}}{B} + \frac{c\vec{E} \times \vec{B}}{B^2} + \frac{mcV_{||}^2}{eB^4} [\vec{B} \times (\vec{B} \cdot \vec{\nabla}) \vec{B}] + \frac{mcV_{\perp}^2}{2eB^3} [\vec{B} \times \vec{\nabla} B], \quad (1)$$

where the first term corresponds to the parallel motion along the field line, the second term to the $\vec{E} \times \vec{B}$ drift, the third term to the curvature drift, and the fourth term to the gradient B drift.

The core plasma of EBTR, away from the electron rings should be curl-free ($\vec{\nabla} \times \vec{B} = 0$), hence the above equation simplifies to:

$$\frac{d\vec{r}}{dt} = v_{||} \frac{\vec{B}}{B} + \frac{c\vec{E} \times \vec{B}}{B^2} + \frac{mc}{2eB^3} (2V_{||}^2 + V_{\perp}^2) [\vec{B} \times \vec{\nabla} B], \quad (2)$$

In addition, the equation of motion for the guiding center is solved:

$$\frac{dv_{||}}{dt} = \frac{-\mu}{MB} \vec{B} \cdot \vec{\nabla} B + \frac{e}{mc} \frac{\vec{E} \cdot \vec{B}}{B}, \quad (3)$$

where $\mu = mv_{\perp}^2/2B$ (magnetic moment), one of the adiabatic invariants of guiding center motion. The three components of Eq. (2) are expressed in a Cartesian coordinate frame with origin at the center of the torus.

The drift equations are solved numerically with a fourth-order Runge-Kutta solver. Although it is difficult to keep track of the local error accrued from the numerical integration with Runge-Kutta, a check on the global error is provided by keeping track of the conservation of the approximate invariant, μ , and the rigorous invariant, E , over collisionless time scales.

In a mirror-type magnetic field, the knowledge of the bounce, or turning point, is desirable. Consequently, a small algorithm using Eq. (3) was written to calculate a time-step that will predict the turning point as closely as possible to ensure μ -conservation. As the guiding center approaches the turning point, a variable time-step is calculated from the following:

$$v_{\parallel}(t) - \frac{dV_{\parallel}}{dt}(t) \Delta t = 0 \quad (4)$$

from which we get the predicted time-step, Δt , obtained by using:

$$\Delta t = C \cdot |V_{\parallel}/dV_{\parallel}/dt|$$

where $C = 1.05$ is an empirical constant.

For passing particles, the coil plane crossing point must be predicted to ensure μ -conservation. Since the EBTR is represented by a single unit cell, particles that are not trapped, i.e., that circulate through the coil plane, must be rotated from one end of the unit cell to the other while preserving μ and velocity direction. Hence, for this class of particles, a μ -conserving time-step is predicted near the point of passing as follows:

$$\Delta t = C \cdot (Z_{\max} - |Z|) / (dv_{\parallel}/dt) \quad (6)$$

where $C = 1.05$, and Z is the local coordinate which measures distance from the midplane of the unit cell along the torus. The variable $Z_{\max}$ is the maximum excursion from the midplane to the coil plane for each particular flux line as measured in a local coordinate frame centered on the minor axis at the midplane of a unit cell. The standard time-step used in MCEBTR is a multiple of the gyroperiod of a 3.5 MeV alpha particle in the on-axis magnetic field strength at the coil plane. For computational purposes, the default time-step is about two gyroperiods; however, this time-step is variable when a particle approaches its turning point or point of passing.

The alphas are initialized in configuration and velocity space according to an algorithm that assigns equal weight to each test particle. In configuration space, the starting grid is weighted to represent the alpha-particle birth distribution (i.e., proportional to the square of the background ion density). The background density and temperature profiles are assumed to be parabolic. The alpha-particles are initialized isotropically in velocity space. The Coulomb drag and pitch-angle scattering model used in MCEBTR is based on the Spitzer coefficients for drag and velocity dispersion ($\langle \frac{dv_{\parallel}}{dt} \rangle$, $\langle \frac{dv_{\parallel}^2}{dt} \rangle$, and $\langle \frac{dv_{\perp}^2}{dt} \rangle$).

The MCEBTR code was initially tested in an analytic representation of an infinite bumpy cylinder magnetic geometry. For 3.52 MeV alpha particles followed for 1500 gyroperiods, energy was conserved to about $\Delta\mu/\mu = 3 \times 10^{-3}$. For an infinite cylinder, the μ -J surfaces on which the alphas drift resemble cylinders. Because of cylindrical symmetry and lack of toroidal curvature, the μ -J drift surfaces are centered around the minor axis. Consequently, none of the μ -J surfaces intersect the wall, and 100% of

the alpha particles are confined on the first bounce, as would be expected. This is because the drift motion of the alphas is governed solely by mirror curvature in an infinite bumpy cylinder. There are no toroidal effects. The passing fraction was calculated to be 64% for this case.

Using a magnetics code which determines the vacuum magnetic field from a toroidal configuration of ring currents enables the effects of toroidal curvature on the shape of the drift surfaces to be investigated. The addition of toroidal curvature effectively distorts the symmetry of the drift surface around the minor axis, thereby pushing them off center. As a result the drift surfaces of certain classes of particles (mainly passing particles) are seen to intersect the first wall, thus adding to the prompt-loss fraction.

Using a single current-loop approximation in modelling each toroidal-field coil, and for the moment neglecting contributions from the ARE coils and electron rings, trapped and passing trajectories (Fig. 1 and Fig. 2) were generated in the magnetic geometry of a unit cell (Fig.3). Preliminary calculations using 500 particle histories estimated the passing fraction to be 61%. Running each particle history for at least one drift time around the minor axis and noting the number of alphas that intersect the first wall ($R_{\text{maj}} = 35 \text{ m}$, $r_{\text{minor}} = 1.28 \text{ m}$) leads to an estimate of the prompt-loss fraction equal to about 72%. The reason for such a large prompt-loss fraction is due to the fact that almost all the passing particles and a few of the trapped ones strike the first wall on their first azimuthal drift. This is because their μ -J drift surfaces are shifted inwardly toward the center of the torus. If we were to use aspect-ratio-enhancement (ARE) coils, the prompt-loss fraction would be considerably

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Fig. 1 Trajectory of a trapped alpha particle drift orbit born on the midplane of the unit cell at a dimensionless minor radius of .5 and pitch angle of 58° .

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Fig. 2 Trajectory of a passing alpha particle drift orbit born on the midplane of the unit cell at a dimensionless minor radius of .5 and pitch angle of 25° .

Fig. 3 The magnetic geometry of a unit cell used in transport calculations.

decreased because these extra coils would help nullify the detrimental effects of toroidal curvature allowing the μ -J surfaces to be more favorably centered. The loss fraction of 72% calculated with MCEBTR compares favorably with the 65% loss fraction predicted in the LANL/MDD EBTR study⁽¹⁾ in the absence of ARE coils.

To get a more rigorous estimate of the prompt-loss fraction it would be necessary to include ARE coils, electron rings, and an ambipolar potential in our analysis. However, because there is no analytic magnetic field representation available for EBTR the magnetics have to be generated from toroidally spaced current loops for each particle during each time-step. This unfortunately, leads to astronomical run times (~3 Cray hours for 500 particle histories). Attempts to fit EBTR magnetics by means of 3-D spline routines have not proved accurate enough to successfully track fusion product orbits. Consequently, this line of research has been terminated. However, we feel that this research along with the results from Ref. 1 indicate the problems of confining alpha particles in large aspect-ratio EBT reactors. It appears that the prompt-loss fraction to be expected in EBTR (~35%) is much larger than that expected in tokamaks (~10%). As a result ignition in an EBT may prove more difficult to achieve than in a tokamak.

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III.3.3 Evaluated Nuclear Elastic Scattering Data from an R-Matrix Formulation

The proper evaluation of the transport effects of nuclear elastic scattering (NES) requires a data base for the light charged nuclei: P, D, ^3He , ^4He , and T over a wide range of energy (0-20 MeV). The best available evaluated data base up to this time for NES has been the compilation by Perkins and Cullen.⁽¹⁾ Their work involved the integration of the differential scattering data over solid angles to determine the reactivity, σ_v , the average fractional energy loss, $\overline{\frac{dE}{E}}$, and the average lab scattering cosine, $\bar{\mu}$. Unfortunately, they did not tabulate the differential scattering data needed to perform rigorous transport calculations. There are also alarming gaps in the cross section data for important reactions such as α -T, α -D and ^3He - ^3He . In an effort to improve the NES data base and to provide some differential scattering information, research was performed in collaboration with G. Hale and D. Dodder in Group T-2 at LANL to produce, CPINT, a processing code used to calculate differential and integrated NES data from the LANL R-Matrix⁽²⁾ codes.

Elastic scattering of charged particles includes components from Coulomb scattering, nuclear scattering, and the Coulomb-nuclear interference between them. The Coulomb scattering is represented by the Rutherford formula and electronic screening is ignored. For distinguishable particles define $\frac{d\sigma}{d\Omega}_{\text{Rd}}$ as the differential Rutherford scattering cross section (barns/steradian) and define $\frac{d\sigma}{d\Omega_i}$ as the differential scattering cross section for identical particles. The lab energy (eV) is denoted as E; μ is the COM scattering cosine, and m_1 is the incident particle mass (amu).

Define the following: Z_1 and Z_2 are the charge numbers of the incident particle and target, s is the spin (needed for identical particles only), A is the target/projectile mass ratio, k is the particle wave number, and η is the dimensionless Coulomb parameter. The differential cross sections(3) can then be written as

$$\frac{d\sigma(\mu, E)}{d\Omega_{Rd}} = \frac{\eta^2}{k^2(1-\mu)^2} \quad (1)$$

and

$$\frac{d\sigma(\mu, E)}{d\Omega_{Ri}} = \frac{2\eta^2}{k^2(1-\mu^2)} \frac{1+\mu^2}{1-\mu^2} + \frac{(-1)^{2s}}{2s+1} \cos\eta \ln \left(\frac{1+\mu}{1-\mu} \right) \quad (2)$$

where

$$K = \frac{A}{1+A} \sqrt{4.7845E-6m_1E}, \quad (3)$$

and

$$\eta = Z_1 Z_2 \frac{\sqrt{(2.48E+4) m_1}}{E} \quad (4)$$

The net elastic scattering cross section for distinguishable particles may be written as

$$\begin{aligned} \frac{d\sigma(\mu, E)}{d\Omega_{Ed}} &= \frac{d\sigma(\mu, E)}{d\Omega_{Rd}} - \frac{2\eta}{1-\mu} \operatorname{Re} \left\{ e^{i\eta \ln \frac{(1-\mu)}{2}} \sum_{\ell=0}^L \frac{2\ell+1}{2} a_{\ell}(E) P_{\ell}(\mu) \right\} \\ &\quad + \sum_{\ell=0}^{2L} \frac{2\ell+1}{2} b_{\ell}(E) P_{\ell}(\mu) \end{aligned} \quad (5)$$

and the cross section for identical particles is

$$\begin{aligned}
\frac{d\sigma(\mu, E)}{d\Omega E_i} = & \frac{d\sigma(\mu, E)}{d\Omega R_i} - \frac{2\eta}{1-\mu^2} R_e \left\{ \sum_{\ell=0}^L \left[(1+\mu) e^{i\eta} n \frac{(1-\mu)}{2} \right. \right. \\
& + (-1)^\ell (1-\mu) e^{i\eta} n \frac{(1-\mu)}{2} \left. \left. \right] \frac{2\ell+1}{2} a_\ell(E) P_\ell(\mu) \right. \\
& + \sum_{\ell=0}^L \frac{4\ell+2}{2} b_\ell(E) P_{2\ell}(\mu),
\end{aligned} \tag{6}$$

where the a_ℓ are complex coefficients for expanding the trace of the nuclear scattering amplitude matrix and the b_ℓ are real coefficients for expanding the nuclear scattering cross section. The value of L represents the highest partial wave contributing to nuclear scattering. In Eqs. (5) and (6), the first term is the pure Rutherford differential cross section, the second term is the Coulomb-nuclear interference differential cross section, and the third term is the pure nuclear differential cross section.

To evaluate the pure nuclear and Coulomb-nuclear differential cross sections simply subtract out the pure Rutherford differential cross sections from Eqs. (5) and (6) as follows:

$$\frac{d\sigma^d}{d\Omega_{CN}}(\mu, E) = \frac{d\sigma(\mu, E)}{d\Omega_{Ed}} - \frac{d\sigma(\mu, E)}{d\Omega_{Rd}} \quad (\text{for unlike particles}) \tag{7}$$

and;

$$\frac{d\sigma^1}{d\Omega_{CN}}(\mu, E) = \frac{d\sigma(\mu, E)}{d\Omega_{E_i}} - \frac{d\sigma(\mu, E)}{d\Omega_{Ri}} \quad (\text{for like particles}) \tag{8}$$

The differential cross sections (Eqs. (7) and (8)) can then be used to calculate the integrated nuclear data $(\sigma_v, \frac{dE}{E}, \mu_L)$ useful for transport calculations.

Since Eqs. (7) and (8) are nonintegrable at the cosine extreme ($\mu=-1$), it is necessary to define a new cross section,

$$\frac{d\sigma}{d\Omega_N} \equiv (1-\mu) \frac{d\sigma}{d\Omega_{CN}} \quad \begin{array}{l} (\mu > 0 \text{ for like particles}) \\ (\mu \geq -1 \text{ for unlike particles}), \end{array} \quad (9)$$

which can be integrated over the appropriate cosine range. Using the newly defined cross section of Eq.(9) allows the calculation of the integrated nuclear data in a manner similar to that of Perkins and Cullen(1) as follows:

$$\sigma_{CN}(E) = 2\pi \int_{\mu_1}^{\mu_2} \frac{1}{(1-\mu)} \frac{d\sigma}{d\Omega_N} d\mu \quad (10)$$

where $\sigma_{CN}(E)$ is the average nuclear and Coulomb nuclear elastic cross section. The average fractional energy loss per collision is

$$\frac{dE}{E} = \frac{4\pi m_1 m_2}{(m_1 + m_2)^2} \frac{1}{\sigma_{CN}(E)} \int_{\mu_1}^{\mu_2} \frac{d\sigma}{d\Omega_N} d\mu, \quad (11)$$

and the average lab scattering cosine, $\mu_L(E)$, can be expressed as

$$\mu_L(E) = \frac{2\pi}{\sigma_{NI}(E)} \int_{\mu_1}^{\mu_2} \mu_L(\mu) \frac{d\sigma}{d\Omega_n} d\mu. \quad (12)$$

The lower limit of integration has the following value:

$$\mu_1 = \begin{cases} -1, & \text{for unlike particles} \\ 0, & \text{for like particles,} \end{cases}$$

and the upper limit of integration is evaluated as

$$\mu_2 = \begin{cases} \max(\mu_+, \mu_{c0}) & \text{for unlike particles} \\ 1, & \text{for like particles,} \end{cases}$$

where μ_+ is the COM cosine where $\frac{d\sigma}{d\Omega_N}$ becomes negative and μ_{c0} is the arbitrary cutoff cosine to represent the boundary between large-angle scattering (for $\mu > \mu_{c0}$) and small-angle scattering (for $\mu < \mu_{c0}$).

This integrated cross section, $\sigma_{CN}(E)$ becomes divergent at $\mu = 1$ because of the Coulomb-nuclear interference term which has a pole at the positive cosine extreme (see Eqs. (5) and (6)). Another argument for using the cutoff cosine is that the interference term oscillates as μ approaches 1 so that $\frac{d\sigma}{d\Omega_N}$ may not be well defined at this limit. For the $p\text{-}^3\text{He}$, reaction the interference term as seen to be increasingly destructive as μ went to 1 where it negated some of the contribution from the pure Rutherford and pure nuclear cross section. In fact, it appears that reactions with a large dimensionless Coulomb parameter, η , such as $p\text{-}^3\text{He}$, $D\text{-}^3\text{He}$, $p\text{-}^6\text{Li}$ etc. will see substantial destructive Coulomb-nuclear interference; however, the exact amount of negation is hard to calculate because of the oscillatory behavior of the interference cross section at small scattering angles. Yet, it appears that the interference effects may be significant for some reactions in contrast to the claims of Perkins and Cullen(1) who did not recognize the oscillatory nature of the interference term and assumed its effects to be entirely washed out by the extremely large Rutherford cross section at small angles.

Preliminary calculations for the $p\text{-}^3\text{He}$ integrated nuclear cross section σv , and the average fractional energy loss (lab), $\frac{dE}{E}$ are seen in Fig. 1 and Fig. 2 along with similar calculations by Perkins and Cullen. The integrated data obtained from the processing code, CPINT, was calculated with a cutoff cosine of $\mu_2 = .94$ which corresponds to a COM scattering angle of 20° . The LLNL data used either a cutoff cosine of $\mu_2 = .94$ or the largest value of μ up to which $\frac{d\sigma}{d\Omega_N}$ was still positive. Comparing the σv values (Fig. 1) generated by R-Matrix parameters with the values generated at LLNL shows good agreement over the entire energy range with the R-Matrix values a little larger than the LLNL values between 0-3 MeV. Likewise, the R-Matrix values of dE/E are greater than the LLNL calculations between 0-4 MeV, but are in close agreement for proton energies greater than 4 MeV. The disagreement between the R-Matrix and LLNL results at low proton energies may be due to discrepancies in the choice of cutoff cosine at low proton energies.

In summary, the processing code, CPINT, has been developed in conjunction with nuclear physicists at LANL. This code evaluates the integrated nuclear data from a R-Matrix formulation which may be useful in NES transport calculations. The R-Matrix method has the following advantages over the formulation of Perkins and Cullen at LLNL:

- (1) differential NES cross section can be easily calculated and tabulated,
- (2) analytic forms of the evaluated nuclear cross sections can be calculated in some cases, and (3) the Coulomb-nuclear interference term can be separated from the pure nuclear and pure Rutherford terms. The disadvantage of the R-Matrix method lies in the unfortunate fact that it has not been applied in the analysis of several key reactions such as $p\text{-D}$, which are crucial in performing fusion product transport calculations for $D\text{-}^3\text{He}$ and Cat-D plasmas.

Fig. 1. P-HE3 NES REACTIVITY .VS. ENERGY

Fig. 2. P-He3 FRACTIONAL ENERGY LOSS .VS. ENERGY

III.3.4 References

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III.3.5 Spatial Effects of NES on Fusion Product Transport in the FRM

Previous studies⁽¹⁾ have shown that close to 45% of 14.7 MeV D-³He protons, can be absolutely confined in a typical FRM (Field Reversed Mirror) plasmas despite the relatively small size of the plasma (plasma radius $\sim 2 f_p^+$ gyroradii). This surprising result lead to the prediction of near-ignited operating conditions for a D-³He fueled FRM. However, recent work ⁽²⁾ has indicated that NES (nuclear elastic scattering) events induce significant perturbations to the confined trajectories of 14.7 MeV, D-³He protons that alter the fp confinement and energy deposition correlations previously calculated for Coulomb slowing down. Ongoing work is focusing on the stochastic calculation of the percent change in fp confinement and heating induced by NES in the closed field region of the FRM.

In order to rigorously estimate the changes in fp confinement and heating, the stochastic transport code, MCFRM⁽¹⁾, has been upgraded to track a statistically significant sample (~ 1000 histories) of 14.7 MeV protons as well as their knock-on progeny. A Hill's Vortex model of field reversal is used to simulate the magnetic topography of a FRM, an elongated version of Hill's Vortex can be used in MCFRM to model fp dynamics in a long-pulsed, quasi-steady, Field Reversed Theta Pinch (FRTTP). Spitzer coefficients are used to simulate Coulomb drag and pitch-angle scattering, and a particle pusher is used to track fp trajectories.

The probability of undergoing a NES event is sampled each time step for every particle history from a cumulative probability distribution. The NES data used is taken from the ENDL library at LLNL⁽³⁾; however, recent work in collaboration with nuclear physicists at LANL has produced additional differential and integrated NES data from a R-Matrix formulation.⁽⁴⁾ The knock-on ions resulting from NES events are

⁺ fp $\equiv$ fusion product

angularly distributed according to a P_1 approximation which assumes isotropic scattering (in the COM system) together with an anisotropic correction factor involving forward and backward scattering.

As a first approximation to the fp and knock-on energy deposition in the closed-field region, the knock-ons are not followed discretely but have their energy smoothed out over their energy surfaces with the use of a confinement criterion developed to explain particle motion in the FRM.⁽⁵⁾ Accordingly, after a knock-on ion is born following a NES event, its dimensionless energy, E_D , and canonical angular momentum, P_θ , are used to determine if the knock-on trajectory is closed-field confined, marginally confined (samples both open and closed field lines) or unconfined. This information together with other energy deposition correlations⁽¹⁾ allow the knock-on energy to be smoothed out over configuration space as a first approximation.

The MCFRM code is being statistically tested to determine the accuracy of the Monte Carlo calculation, or more concisely, the number of particle histories, N , that are needed to provide computational estimates of the energy deposition to within an error tolerance, ϵ . After this statistical analysis the energy deposition will be calculated with NES effects over a range of plasma parameters together with the steady-state fp distribution function.

The P_θ distribution of thermalized fps distorted by NES will be calculated thus allowing estimates of ash deposition. Various variance reduction methods such as the exponential transform, Russian roulette and angular biasing techniques are being utilized in MCFRM to enhance computational efficiency.

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III.4 ION-ELECTRON EQUILIBRATION STUDIES

Much of the enhancement of the percentage of fusion product energy deposited to background ions through NES (and anomalous microinstabilities), is lost by strong equilibration between the background ions and electrons; thus reducing the effectiveness of these mechanisms on increasing T_{ion}/T_e , and lowering $n\tau_e|_{\text{ignition}}$. Not only are the effects of NES (and fusion product driven microinstabilities) enhanced by reducing the equilibration power, P_{ie} , but ignition requirements are reduced by the lowering of P_{ie} alone (in particular advanced fuel cases are studied here, where $T_i/T_e > 1$).

Effects which occur in plasmas, without external injection of energy, which have been investigated include 1) Rosenbluth depletion effect(1,2) mentioned in the previous progress report, 2) Inclusion of non-dominant terms in the collision integral(3,4,5) used for energy transfer calculations, and 3) the effect of a magnetic field on reducing the Coulomb log for ion/electron collisions.(6) The investigation of these effects, and the corresponding lowering of ignition requirements are discussed below.

III.4.1. Rosenbluth Depletion Effect

In a Maxwellian electron distribution, only the low energy electrons receive energy from ions, on the average. For the case of a Maxwellian ion distribution, the low energy electrons receive energy from the ions slightly faster than they can redistribute it amongst themselves, causing a depletion of the low energy "foot" of the electron distribution. This depletion causes a reduction in P_{ie} , since there are fewer electrons available to receive energy. Rosenbluth(1), derived an analytical model for this

process which predicted modest reductions ($\sim 2\%$) in P_{ie} for $T_{ion}/T_e \lesssim 2$. Previous numerical calculations⁽²⁾, however, predicted a much stronger effect. These numerical stimulations were redone and found to be in error. Furthermore, the present numerical calculations agree with Rosenbluth's model (see Fig.1.) for $T_i/T_e \lesssim 50$, above which Rosenbluth's model overestimates the effect.

III.4.2 Numerical Method

The method for numerical simulation of the depletion effect was the same as in reference 2, except a different Fokker-Planck code was used, (namely the FPPAC⁽⁷⁾ code was used - from MFE library). The method involves starting a Maxwellian ion and electron distribution at some $T_i/T_e (< 1)$. The ion distribution is held constant, and the electrons are allowed to heat up. The foot of the electron distribution was quickly distorted (on a faster time scale than the ion/electron equilibration time), due to the depletion effect. The reduction in P_{ie} is presented (Fig. 1) as the actual energy exchange rate, normalized to the energy exchange rate for a Maxwellian electron distribution with the same average energy as the calculated distribution (T_e is considered to be $2/3$ of the average energy of the actual distribution). The calculation of the reduction of P_{ie} checked at particular T_{ion}/T_e 's by starting the electrons at different T_e 's and noting the reduction was the same at the particular T_i/T_e .

III.4.3 Non Dominant Terms

In the usual evaluation of the collision integral for a test particle interacting with field particles "phenomenological" cutoffs at the maximum and minimum impact parameters are introduced to avoid divergences. It is possible to avoid the cutoffs by 1) accurately treating the dynamics of

Figure 1. Rosenbluth Depletion Effect on Equilibration Power P_{ie} .

close collisions at small impact parameters, and 2) including the shielding properties of the plasma through the plasma dielectric constant ϵ . Both these modifications lead to additional terms (to the usual Coulomb $\log_e \Lambda$) called non-dominant terms.

The usual(3,4,5) solution procedure, including these modifications, is to divide the integral over impact parameters (or scattering angle) into two parts. For small impact parameters (where plasma shielding effects are small) the large angle scattering events are accurately modeled introducing "Type 1" nondominant terms. For large impact parameters, the collisions may be correctly assumed to be small angle, but the shielding properties of the plasma must be included leading to "Type 2" nondominant terms. (The plasma shielding terms are accounted for by using the Balescu-Lenard form of the kinetic equation). The arising nondominant terms are, in general, functions of X ($X \equiv$ test particle speed/average electron thermal speed for Maxwellian electrons). To determine the effect of the non-dominant terms on P_{ie} , the energy exchange rate of a test ion with a Maxwellian electron distribution is averaged over a Maxwellian ion distribution for the test particles (including the non-dominant term at each particular test particle speed).

The "type 1" N.D. correction to $\log_e \Lambda_{ie}$ encountered in calculating P_{ie} was analytically predicted by Fankel(3) to be $= -.577$. Numerical evaluation of this quantity provided a slightly larger reduction of $\log_e \Lambda_{ie}$, $= -.635$, over the range of X which predominantly contributes to P_{ie} ($0 < X < 10 \times \sqrt{m_e/m_i}$).

Previous evaluation of "type 2" nondominant terms had been restricted to cases of energetic ions slowing on electrons, and neglected terms $\sim m_e/m_{ion}$, which significantly contribute to P_{ie} (i.e., they become important for $X \lesssim 10 \times \sqrt{m_e/m_i}$). Including these terms, the reduction in the usual $\log_e \Lambda_{ie}$, due to "type 2" N.D. terms was $-.50$, over the range of X , $0 < X < 10 \times \sqrt{m_e/m_i}$.

Thus the total reduction to $\ln\Lambda_{ie}$ (here P_{ie}) is - 1.14. This amounts to only a small ($\sim 7\%$) reduction in equilibration power for typical magnetic confinement conditions ($\ln\Lambda_{ei} \sim 18$), or $\sim 15\%$ decrease in P_{ie} for typical inertial confinement conditions ($\ln\Lambda_{ei} \sim 8$).

III.4.4 Effect of a Magnetic Field on Reducing P_{ie}

McNally⁽⁵⁾ has argued that the maximum impact parameter in $\ln\Lambda_{ie}$ should be (average thermal ion speed)/(electron cyclotron frequency), whenever this quantity is less than the usual Debye length. As an example, consider a plasma with $n_e = 5 \times 10^{14} \text{ cm}^{-3}$, $B = 50 \text{ kG}$, $\beta \approx 3\%$. Under these conditions, McNally's model predicts a 25% reduction in P_{ie} , (for $\ln\Lambda_{ei} \sim 18$).

Combined with the previously mentioned effects, this amounts to a total reduction of $\sim 35\%$ in the equilibration power. In the following analysis a reduction of only 25% is assumed ($\ln\Lambda_{ie}^* = 75\%$ classical $\ln\Lambda_{ie}$). This conservative approach is taken since the effects may not combine linearly.

III.4.5 Effect on Global Plasma Properties

The effects of reducing P_{ie} , on global plasma characteristics such as n_i/n_e , were investigated with the same Hot Ion Mode analysis used for NES. In the following examples, the inclusion of NES and reduced P_{ie} simultaneously is also shown.

Figure 2 shows the T_{ion}/T_e ratios for a Cat-D plasma. Note that the reduction in P_{ie} has a stronger effect than NES, on increasing T_i/T_e (at 60 keV NES alone causes $\sim 3\%$ increase in T_i/T_e , vs $\sim 10\%$ rise through reduced P_{ie}). Figure 3 shows the percentage of fusion product energy deposited to background ions. Note that when P_{ie} is reduced, the enhancement from NES is also reduced. This is because with

Figure 2. T_{ion}/T_e for Cat-D Plasma with Reduced Equilibration Power (P_{ie}) and NES.

Figure 3. Percentage of Slowing Down Power to Ions for Cat-D Plasmas, with Reduced Equilibration Power and NES.

lower P_{ie} the electrons tend to have a lower T_e , which causes the fusion products to slow down faster, decreasing their NES probability. Thus, the combination of the two effects is non linear. Figure 4 shows the effects of lowering P_{ie} (and including NES) on reducing ignition requirements, for Cat-D. Note that the combination of these effects independently overestimates the actual reduction (at 60 keV NES alone causes ~35% reduction, and lowering P_{ie} causes ~55% reduction, but only ~65% reduction occurs with both considered simultaneously).

III.4.6 Summary

Reduction of equilibration power has been shown to reduce $n\tau_e|_{\text{ignition}}$ for Cat-D plasmas by $\geq 50\%$ in certain circumstances. The combination with NES results in further reductions, but not as much as a linear combination of the independent effects would predict. Investigation of this effect for D-T use in high field tokamaks (where $T_i/T_e \geq 1$) could show further reductions in $n\tau_e|_{\text{ig}}$ for these devices, by raises in the T_{ion}/T_e ratio.

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Figure 4. Ignition Requirements (including cyclotron losses) for Cat-D Plasma, with Reduced Equilibration Power and NES.

IV. STAFF CONTRIBUTORS AND ACKNOWLEDGMENTS

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V. REPORTS, THESES, AND PUBLICATIONS

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