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"Fundamental Studies of Fusion Plasmas"

by
R. E. Aamodt (P.I.), P. J. Catto, D. A. D'Ippolito, J. R. Myra, and D. A. Russell

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LODESTAR RESEARCH CORPORATION
2400 Central Avenue
Boulder, Colorado 80301

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MASTER
I. CONTRACTUAL INFORMATION

A. This report reviews the research results for the first six months of the present contracting period 11/1/91 to 10/30/92 for the main contract. Additionally, a contract mod of $85K was added at the end of the previous contract year and the results of that study are included herein (Sec. III, A-5).

B. The proposed budget for this project is presently in keeping with the funded proposal within allowed grant guidelines. There will be no remaining monies at the end of this project period.

II. IMPACT OF RESEARCH PROGRAM

The major portion of this program is devoted to critical ICH phenomena. The topics include edge physics, fast wave propagation, ICH induced high frequency instabilities, and a preliminary antenna design for Ignitor. This research was strongly coordinated with the world’s experimental and design teams at JET, Culham, ORNL, and Ignitor. The results have been widely publicized at both general scientific meetings and topical workshops including the speciality workshop on ICRF design and physics sponsored by Lodestar in April 1992. The combination of theory, empirical modeling, and engineering design in this program makes this research particularly important for the design of future devices and for the understanding and performance projections of present tokamak devices.

Additionally, the development of a diagnostic of runaway electrons on TEXT has proven particularly useful for the fundamental understanding of energetic electron confinement. This work has led to a better quantitative basis for quasilinear theory and the role of magnetic vs. electrostatic field fluctuations on electron transport. An APS invited talk was given on this subject and a collaboration with PPPL personnel was also initiated.

Ongoing research on these topics will continue for the remainder of the contract period and the strong collaborations are expected to continue, enhancing both the relevance of the work and its immediate impact on areas needing critical understanding.
III. SUMMARY OF TECHNICAL PROGRESS

A. ICRF-Edge Physics Studies

1. Introduction

One of the most significant new results of the ICRF program in the past five years is the great increase in our understanding of the interaction of high-power ICRF waves with the tokamak edge plasma.\(^1\) There is now a large body of experimental results and theoretical modeling showing that the ICRF-edge interaction greatly influences the success of ICRF heating and current-drive schemes in tokamaks. There is experimental evidence that the ICRF-edge plasma interaction produces low and high-Z impurity generation, anomalous loading (i.e. power dissipation) and decreased heating efficiency, and the generation of DC electric fields and \(E \times B\) drifts with resulting enhanced edge transport and profile modification. Both the impurities generated at the Faraday screen (FS) and the ICRF-driven transport have been shown to influence the ability to achieve \(H\)-modes on JET with ICRF heating alone, with the remarkable result that the \(H\)-mode threshold power, global confinement, and duration depend on the phasing of the antenna. Our recent work has shown that many of these effects can be attributed to the existence of rf sheaths, which have therefore been the major focus of the Lodestar edge physics research program. Our theoretical work has been done in close collaboration with M. Bures and J. Jacquino of the Joint European Torus (JET) experimental ICRF group, who have accumulated a large database on ICRF-edge physics and have played an active role in the development of rf sheath modeling. We have also maintained contact with the experimental ICRF programs on TFTR, DIII-D, MIT, Phaedrus and TEXTOR, and participated in the ICRF R & D committee for the U.S. ITER home team.

The interaction of ICRF waves with the material boundaries in the tokamak leads to the formation of rf sheaths when closed circuits linking substantial rf magnetic flux are formed by the magnetic field lines and various conducting surfaces (Faraday screen, limiters, etc.).\(^2\)-\(^{12}\) Equivalently, one can regard rf sheaths as being driven by the component of the rf electric field parallel to the equilibrium B field, i.e. the slow wave component. Thus, rf sheaths are driven quite strongly by slow wave couplers such as ion-Bernstein-wave antennas and also by fast wave antennas which are imperfectly aligned to the local magnetic field. A model\(^9\) based on the theory of rf sheaths has been successful in explaining many experimental dependences\(^8\)-\(^{11}\) of the metal impurity generation in JET (see Sec. 2). These models also predict substantial power dissipation in the sheaths when the induced voltage across the sheath and the local plasma density are sufficiently large.\(^4\) This effect has been observed in JET (as a significant reduction in heating efficiency) by
reversing the toroidal field and using monopole phasing to create unusually large rf sheath voltages. Sheath-driven parallel currents in the scrape-off-layer (SOL), which flow along field lines from the rf-sheath-biased antennas into limiters, have been measured on TEXTOR. Similar SOL modifications (sheath-induced impurities, currents, and DC potentials) were measured on Phaedrus for three antenna phasings (0, 90 and 180 degrees) and it was shown that the presence of an insulating limiter on each side of the antenna greatly reduced the edge modifications in agreement with the predictions of rf sheath theory. Thus, there is a growing body of experimental evidence supporting the importance of rf sheath effects in the SOL, and a growing effort to understand and control them.

In the past year, our work has concentrated on another important effect predicted by rf sheath theory, viz. the generation of ICRF-driven convective cells in front of the antennas. The spatial variation of the DC sheath potential \( \Phi_0(r,\theta) \) drives a convective flow pattern, which can penetrate radially a substantial distance into the SOL. This convection may explain the observed density profile flattening with monopole phasing on JET (see Sec. 3) and the observed phasing sensitivity of the ICRF-induced H-mode (see Sec. 4). It is particularly important to understand the physics of the monopole ICRF H-mode, and to discover whether the extended duration and reduced impurity radiation characteristic of this regime is applicable to reactor operation.

In large tokamaks, such as JET and TFTR, the typical heating regime has good single pass absorption, and the amount of reflected or transmitted wave energy impacting the SOL is small. Thus, rf sheaths driven by the near fields of the antenna are the most important concern in these experiments, and most of our work to date has concentrated on this problem. However, there are a number of interesting conditions where the single pass absorption can be low (ICRF current drive or direct electron heating, heating with a low minority concentration, or heating in smaller tokamaks) and rf sheaths driven by ICRF wave energy far from the antenna become important (see Sec. 5).

The following sub-sections describe our work in these areas in more detail.

2. ICRF Impurity Studies

One of the critical problems for ITER and other reactor-relevant machines is the control of impurities caused by auxiliary heating and by power flow to the divertor. Recently, an international effort has shown that the dominant ICRF-specific impurity generation mechanism is sputtering and self-sputtering by ions accelerated to hundreds of eV energies by rf sheaths. Our work on this subject with M. Bures and J. Jacquinot of JET has resulted in a detailed model of the metal impurity generation at the Faraday screen (FS) for the JET antenna geometry. An extensive comparison has shown good
agreement with JET impurity influx data for a variety of limiter and FS materials and for different operational scenarios. In the past year, we completed this effort with an analysis of the data from operation with the new Be Faraday screens on JET. The previous Be-gettered nickel screens were replaced by screens manufactured from solid Be with a more open design. The elimination of the residual Ni influx (from the gaps of the old gettered screens) allowed coupling to the H-mode for the first time in monopole phasing, emphasizing the importance of eliminating even trace amounts of high-Z impurities. The Be influx was less than $10^{19}$ atoms MW$^{-1}$ s$^{-1}$ (comparable to that with the previous screens) in agreement with the theoretical predictions. Another interesting aspect of the experiment was the investigation of the operating regime where the direction of the equilibrium magnetic field is reversed, so that the mismatch angle $\theta$ between the magnetic field and the screen bars increases from about 5 to 22 degrees. The model predicts that the voltage drop across the sheaths (and hence the energy dissipated by the sheath-accelerated ions) is proportional to $\sin \theta$. No change was observed for dipole phasing, where the rf sheath voltage vanishes by symmetry for either direction of B. With monopole phasing, it was found that the efficiency of ICRF heating is substantially degraded (about a 40% drop in central stored energy) when the direction of the toroidal field is reversed. The drop in coupled rf power was comparable to the theoretically predicted dissipation in the rf sheaths.

The beryllium screen experiments, especially the ICRF H-mode results, confirm the usefulness of the JET-Lodestar prescription for reducing ICRF-specific impurities to negligible levels in heating experiments. First, regimes with poor single pass absorption should be avoided, if possible, to prevent rf sheath formation away from the antenna which can pollute the SOL with limiter and wall impurities. Second, rf sheath formation at the antenna should be inhibited by proper antenna and Faraday screen design. A properly designed antenna should include the following features:

1) FS elements should be aligned as closely as possible to the local equilibrium magnetic field;
2) the density at the FS should be minimized by the use of close-fitting side protection tiles;
3) the antennas should be designed with two symmetric current straps and operated with the straps out of phase (toroidal dipole phasing);
4) the antenna voltage (and hence power per antenna) should be minimized;
5) a low-Z material, such as beryllium, should be used for the Faraday screen material because it has a relatively low self-sputtering coefficient in the energy range 0.5 - 1 keV.
The adaptation of these rules for ICRF current drive is a subject of future research.

3. ICRF Convective Cell Simulation

The radial and poloidal structure\(^9\) of the rectified or time-averaged rf sheath potential\(^4\) \(\Phi_{\text{rf}}\) produces \(E \times B\) convection in front of ICRF antennas.\(^{14}\) We have derived a generalized "convective cell" equation to describe the spatial variation of the electrostatic potential in the global SOL, and written a 2D fluid simulation code (CC2) to solve a model problem representing ICRF edge convection.\(^{14,15}\) The model consists of an equation for the slow-time evolution of the field-line-averaged electrostatic potential \(\Phi(x,y)\) in the SOL, assuming fixed density and temperature profiles. This convective cell equation includes the physics of advection, viscosity and parallel current flow into sheaths at material boundaries. It is solved subject to the B.C. that \(\Phi=\Phi_0\) at the antenna tile tangency surface (ATTS) and at the last closed flux surface (LCFS), where \(\Phi_0=\Phi_{\text{rf}}(x,y)+3T_e(x)\) is the DC sheath potential and \(x,y\) denote the radial and poloidal coordinates. This condition ensures no net parallel current \(J_\parallel\) leaving the plasma on field lines at the two radial boundaries of the SOL. The B.C.'s supply the drive for the \(E \times B\) flow in the simulation. In the SOL itself, a net \(J_\parallel\) flowing to the axial boundaries (limiters, diverter plates, etc.) is allowed and is balanced by \(J_\perp\) due to the ion polarization drift. The ratio \(J_\parallel/J_\perp=\lambda=(\rho_\parallel/L_\parallel)(L_y/\pi\rho_\phi)^4\) depends on the parallel connection length \(L_\parallel\), the poloidal scale length \(L_y\), and the Larmor radii based on \(T_e(\rho_\parallel)\) and \(\Phi(\rho_\phi)\). The limit of 1D sheath theory is recovered as \(\lambda \to \infty\), and as \(\lambda \to 0\) the convective cell equation reduces to the usual Navier-Stokes equation for 2D incompressible flow neglecting pressure gradients. The equipotential contours are the stream lines for the \(E \times B\) flow.

The rectified \(\text{rf}\) sheath potential \(\Phi_{\text{rf}}\) has a maximum value at the ATTS and decays radially in the SOL on the scale of the electron skin depth; it varies poloidally, and drives \(E \times B\) convective cells, with the periodicity of the Faraday screen structure and also on the scale length of the antenna. Convection on the scale of the poloidal antenna length will have the greatest interaction with the plasma edge and would be present even if the screen were removed. The smaller convective cells generated by the FS periodicity have less penetration, but increase the poloidal asymmetry of the equilibrium and drive local turbulence. The Bohm sheath potential \((\sim 3T_e)\) becomes the dominant term near the LCFS. \(T_e\) is assumed to have no poloidal variation so that the thermal sheath drives a radially sheared \(E \times B\) flow in the poloidal direction which opposes the \(\text{rf}\)-induced convective flow.

We solve the time-dependent coupled equations for the slow time evolution of the potential \(\Phi\) and vorticity \(\omega=\nabla^2\Phi\) using a combination of finite-differencing and
spectral techniques and study the interaction of the two counter-streaming flows. For sufficiently small antenna-plasma separation $\Delta x$, the interaction is strong and the convective cell disrupts the poloidally-uniform flow near the LCFS. The flow contours reconnect giving birth to new vortices which connect the LCFS to the antenna tile surface. The convection time around these vortices is much smaller than the time for sonic flow along field lines, so that the density and temperature profiles would be essentially flattened across the vortices if these profiles had been evolved in the simulation.

The numerical results correspond to the experimental observations on limiter operation in JET as follows: (1) the rf sheath potential driving the convective cells is large only in monopole phasing, in agreement with the profile flattening data on JET and other tokamaks; and (2) the rf-induced convection provides an explanation for the significant plasma density at the FS required in the modeling to account for the experimental impurity data. The comparison with JET H-mode data is given in the next section.

4. ICRF H-mode Modeling

An important development in the JET ICRF program, made possible by the great reduction in impurities, was the achievement of H-modes induced by ICRF heating alone. ICRF H-modes were obtained with both in-phase (monopole) and out-of-phase (dipole) operation of the 2 strap ICRF antennas. Dipole phasing yielded standard 3 MA H-modes of high quality: the threshold rf power was 5 MW, the energy confinement time was $\tau_E = 2.5 - 2.8 \tau_G$ (Goldston L-mode scaling), and the plasma density and impurity concentration typically increased with time until radiative collapse in the X-point region terminated the H-mode. By contrast, the H-mode obtained with monopole phasing had a larger threshold power (8 MW), smaller increase in confinement ($\tau_E = 1.7 \tau_G$), and a greatly reduced density rise. The early experimental results suggest that the monopole ICRF H-mode may provide a means of achieving density control and reducing the impurity accumulation in H-mode operation. The ICRF H-mode behavior suggests the existence of a phasing-dependent edge transport mechanism which degrades the effectiveness of the "transport barrier" but enhances the effectiveness of the divertor to withstand the power flow (perhaps by spreading out the heat load). The ICRF convective cells described in Sec. 3 provide a strong (and so far, the only) candidate mechanism to explain this phasing-dependent transport.

In the simulation rf-induced convective cells interact with the sheared flow layer at the plasma edge to produce secondary vortices which connect the last closed flux surface to the limiters, enhancing the edge transport and increasing the poloidal variation of the equilibrium density and temperature. All of these effects are likely to
inhibit the formation of the highly-sheared, poloidally uniform H-mode transport barrier, although this effect is not yet in our simulation. In addition to the phasing dependence, the simulation yields two other promising quantitative predictions: (1) for typical JET parameters, the rf convective flow velocity near the ATTS ($v_E \approx 10 - 20$ km/sec) is comparable to that observed experimentally in the H-mode on machines such as DIII-D, TEXTOR$^{18}$ and JFT-2M$^{19}$ and is rapid enough to broaden the density and temperature profiles; (2) the computed antenna-plasma distance for strong interaction between the convective cell and the sheared flow layer at the LCFS is comparable with that measured during ICRF monopole H-modes. Moreover, the measured scaling of $\Delta x$ with rf power and antenna phasing has been analyzed$^5$ by means of a simple model to compute the ratio of density scale lengths, $\lambda_n$(H-mode)/$\lambda_n$(L-mode), for dipole and monopole H-mode shots. The results show that significant profile flattening occurs for the monopole H-mode. Thus, the simulation and data analysis suggest that ICRF-driven convection is a strong candidate to explain the JET ICRF H-mode results.

Future work will incorporate self-consistent evolution of the SOL density and temperature profiles in the simulation, and will attempt to assess the effects of the edge convection on H-mode formation using presently evolving theoretical models.

5. Far Field Sheath Modeling

In situations where the single pass absorption is low, the potential of ICRF fast waves for both heating and current drive can be partially mitigated by difficulties with edge interaction including enhanced impurities and transport. While boronization provides control of metallic influxes, unfortunately the performance of the machine degrades as the boron wears off. So far, beryllium has proven to be the material of choice in ICRF experiments but US tokamaks in the foreseeable future will not be beryllium coated. Even when advanced coatings are employed to reduce the sputtering of dangerous impurities, perturbing effects of the ICRF on the edge, such as convective cells, still exist. Thus, there is a considerable need to exercise caution with regard to the deleterious effects of far field ICRF wave interactions with wall, limiters and the edge plasma.$^{20}$

Recent work at Lodestar on this topic$^{21}$ has been motivated by current drive experiments on DIII-D$^{22}$ as well as upcoming heating experiments on C-MOD.$^{23}$ In the case of current drive, the expected single pass absorption is low, resulting in large rf electric fields at the plasma edge. While ion heating allows for good single pass absorption in the best parameter regimes, it still may not be optimal in some of the planned modes of C-MOD operation.
To develop a modeling capability for these experiments, we have begun to delineate and examine the underlying physics issues, viz. the coupling of fast and slow waves at tokamak walls and limiters, and the coupling of slow waves to sheaths. When fast waves are incident on a conducting surface with sufficiently nontrivial geometry, image currents in the surface generally contain a component \( J_\parallel \), aligned parallel to the tokamak B-field. This component drives a slow wave which contains an \( E_\parallel \) that can drive sheaths near the conducting surface in the same way as for the near-field sheath problem that has been so extensively studied (see Ref. 9 and references therein).

A small special purpose code, describing just the plasma edge region near the limiters, is under development to determine the amplitude of the slow wave generated by the fast wave interaction with a limiter. Preliminary results show the expected generation of the slow wave, which is evanescent and radially confined to the region near the limiter tips. By running this code in tandem with the FREMIR single pass absorption code obtained from JET, it will be possible to relate the amount of generated slow wave to the fast wave power from the antenna, and to the critical plasma parameters that determine single pass absorption (such as line density, \( T_e \) and \( k_\parallel \)).

At a more advanced stage is the understanding of how the slow wave couples to the sheath. A set of unified equations has been developed which includes both the linearized physics of the usual cold fluid hydrodynamic waves and the nonlinear and adiabatic electron physics of the rf sheath. The unified equations have been employed to perform an asymptotic matching of the rf sheath solution to the slow wave fields. The sheaths couple directly to the slow waves because the sheaths are sensitive to the electron dynamics, which couple to \( E_\parallel \). This formalism determines the voltage drop across the sheath as a function of the slow wave amplitude.

Complementary to the view that the excited slow waves drive sheaths is the view that the sheaths provide a boundary condition for the wave fields. In many case, it is possible to show that a new set of effective boundary conditions at conducting surfaces results. The usual condition of vanishing tangential components

\[
E_{\text{tang}} = 0
\]

is replaced by new sheath boundary conditions on the normal component of the rf B field and the parallel component of the electric field

\[
B_{\text{norm}} = 0, \ E_\parallel = 0.
\]

In trivial geometries, it is worth noting that (1) and (2) are equivalent. The new boundary conditions arise from the rapid change in the parallel dielectric constant \( \epsilon_\parallel \) as one
passes from the plasma into the electron-poor sheath, a fact which has also played an important role in previous wave/sheath studies\textsuperscript{26} as well as in the Lodestar work on sheath plasma waves in the context of IBW launchers.\textsuperscript{25}

**B. Propagation of the Fast Wave Across a Minority Ion Cyclotron Resonance in a Toroidal Magnetic Field**

1. Introduction

Fast wave heating of tokamak plasmas in the ion cyclotron range of frequencies (ICRF) is well established in both the minority and second harmonic regimes. For both these schemes an understanding of the spatial dependence of the resonant interaction region is highly desirable since this would lead to greater control over the heating. Three important factors which influence the resonance width are the Doppler effect, the spatial variation of the equilibrium magnetic field (including the effects of trapped particles), and the anisotropy of the distribution function. All three effects are retained in work being performed in collaboration with Chris Lashmore-Davies of Culham. The work which is in the process of being written up represents a small but significant step towards obtaining a self-consistent description of wave propagation and quasilinear heating. For simplicity the work specifically considers only the fast wave minority heating case, even though the key techniques are valid more generally.

A fast (or compressional Alfvén) wave incident from the low field side and propagating in a minority heated cold plasma passes through the minority resonance before encountering the turning point at the ion-ion (or hybrid) cut-off and then tunneling through to the ion-ion resonance where it converts to a slow (or shear Alfvén) wave. Dissipative wave absorption by the ions does not occur in the vicinity of the minority resonance when kinetic effects such as Landau damping are not retained. Wave absorption in the cold limit is all due to mode conversion to the slow wave which is assumed to deposit its energy in the electrons which are also heated collisionally. When the thermal effect of Landau damping is retained the slow wave becomes an ion Bernstein wave which for $k_{\parallel} \neq 0$ damps on the electrons, where $k_{\parallel}$ is the parallel wavenumber.

Nearly all kinetic treatments of fast wave minority heating of inhomogeneous plasmas in the ion cyclotron range of frequencies assume the magnetic field only varies in the direction perpendicular to the magnetic field. This assumption is made for analytic tractability, but is normally inappropriate in a minority heated tokamak. The
rotational transform in a tokamak causes a minority ion moving along a magnetic field line to encounter only isolated points at which its local gyrofrequency $\Omega$ is equal to the wave frequency $\omega$. The wave–particle interaction time and the energy absorbed from the wave are controlled by the poloidal dependence of $\Omega$ and the pitch angle of the ion (the rotational transform causes a small number of trapped particles to reflect in the resonant layer). The non-local wave–particle interaction due to the parallel variation of the magnetic field introduces a dissipation which combines with Landau damping to give a more complete description of absorption.

2. Parallel Magnetic Field Effects on the Plasma Response

The model which has been used in most calculations of ion cyclotron absorption in a tokamak is that of an equilibrium magnetic field with straight field lines and a perpendicular gradient in strength. This model has produced much useful information and includes the two essential ingredients of Doppler broadening and perpendicular magnetic field inhomogeneity. However, a straight field line model cannot properly account for effects due to a parallel variation of the magnetic field.

In a tokamak magnetic field an ion (even in the zero Larmor radius limit considered) experiences a parallel variation in the equilibrium magnetic field due to the rotational transform. This feature has a profound effect on resonant ions compared with the straight magnetic field model. For the latter case an ion if resonant along a particular field line will remain in resonance since the field lines are straight. All other ions (i.e. with different values of $v_{||}$) will always be out of resonance on this particular field line. This situation also applies to the finite Larmor radius broadening where a given ion is either in resonance over the whole of its orbit or not at all. For a toroidal model the situation is quite different. Since the field lines are curved, ions with different values of $v_{||}$ can be in resonance on the same field line but at different spatial locations. Furthermore these ions do not remain in resonance permanently, as in the straight field model, but pass in and out of resonance. Thus, for the straight field line case there are fewer resonant ions interacting strongly and locally with the wave, whereas in the curved field line model there are many more resonant ions interacting weakly and non-locally with the wave.

In order to include the effect of rotational transform, a simplified concentric, circular flux surface model of a tokamak is employed. Using this model the rotational transform and the parallel variation of the magnetic field are retained in the derivation of the plasma response function. This parallel variation also means that trapped ions are present and must be accounted for in the derivation. The procedure employed retains the effects of the trapped particles that reflect in the resonance layer as well as the
remaining trapped and passing particles. To date, there are very few treatments of the effect of parallel field variation and trapped ions on fast wave heating at the ion cyclotron resonance. Particularly insightful is the work of Faulconer\textsuperscript{29} and Smithe et al.\textsuperscript{30} in which the parallel field variation model employed is partially motivated by the earlier work of Itoh et al.\textsuperscript{31} Their plasma response function is found by retaining the leading effect of the field inhomogeneity as well as Doppler broadening, but without introducing the full complications of tokamak geometry.

The model employed in the work presented here implicitly assumes a collisional disruption of the particle motion along the field relative to the wave occurs between interactions with the rf, as described in more detail in Refs. 30 and 32. As a result of the collisions, only the most recent details of the particle's trajectory need be retained in the vicinity of the resonance layer. This simplified treatment of the trajectory retains the non-local nature of the interaction by permitting the trapped and passing particles to pass through resonance, but dramatically simplifies the mathematical description.

The preceding simplifications allow a plasma response function to be derived for Maxwellian minority ions which generalizes and extends previous modified Z function forms\textsuperscript{29-31} obtained from a slab approximation to a tokamak (which also retain parallel inhomogeneity as well as Doppler broadening effects). The plasma response function obtained for a Maxwellian minority includes both passing and trapped ions and may be written as

\begin{equation}
Z(z;a,\sigma,\xi) = \int_{0}^{\infty} dx \exp\left\{ \frac{-izx}{4[1-i(a-\frac{1}{3}\xi)^2 x^2]} \right\}
\end{equation}

where

\[ z = \frac{(\omega-\Omega)R}{|n|v_t}, \quad a = \frac{R\alpha R/\alpha}{q|n|}, \quad \xi = \frac{\Omega R^2}{q^2 n^2 v_t^3} \hat{a} \hat{\beta}^2, \quad \text{and} \quad \sigma = \frac{|n|v_t}{2\Omega R} \ll 1, \]

with \( \omega = \) wave frequency, \( \Omega = \) minority gyrofrequency, \( R = \) major radius, \( q = \) safety factor, \( \beta = \) poloidal angle, \( n = \) toroidal mode number and \( v_t = \) minority thermal speed. For \( \sigma \to 0 \) Eq. (2) reduces to the form given in Refs. 29 and 30. The new inhomogeneity effects enter for off-magnetic axis heating (\( \xi \neq 0 \)) and at low toroidal mode numbers. The spatial dependence of \( Z \) enters via \( z \), with \( a, \xi \) and \( \sigma \) parameters to be specified.

In the limit of small minority concentrations Eq. (3) may be used to obtain the standard result for the transmission coefficient.\textsuperscript{33,34} The small minority concentration expression for the transmission coefficient is shown to be valid for more general unper-
turbed distribution functions that are arbitrary functions of pitch angle and speed on each flux surface provided $k_{\parallel} \rho \ll 1$, where $k_{\parallel}$ is the parallel wavenumber and $\rho$ the minority gyroradius.

For a bi-Maxwellian distribution function an explicit and compact expression for the plasma response function is derived which for strong anisotropy substantially modifies the Maxwellian result. Introducing the anisotropy parameter $A = T_\perp/T_\parallel > 1$, the plasma response function for the bi-Maxwellian case is

$$Z_b(z;\alpha,\sigma,\xi,k) = \int_0^\infty \frac{dx \exp \left\{ \frac{izx - (1 - \frac{1}{2}ax)^2x^2}{4[1-i(\alpha - \frac{1}{2}\alpha^2x - \frac{1}{2}\xi x)\sigma x^2]} \right\}}{[1 + i\frac{1}{2}(1 - \frac{1}{2}ax)kx^2][1-i(\alpha - \frac{1}{2}\alpha^2x - \frac{1}{2}\xi x)\sigma x^2]^{1/2}}$$

(4)

where the new parameter $k = (A\delta \Omega / \partial \Phi) / (4q n \Omega)$ can have a magnitude greater than or on the order of unity in some high power minority heating experiments. The derivation of Eq. (4) indicates that strong anisotropy ($A > 1$) modifies the influence of the inhomogeneity in the regions about $x \leq 1$ and $x \approx 2/a$ by altering the fraction of particles that resonate with the wave, where $x = |n|v/t$ with $t$ an earlier time along the trajectory.

Even though Eq. (4) is a significant improvement over the Maxwellian case, a bi-Maxwellian does not contain all the anisotropic structure necessary to model the distribution function of a minority heated plasma. Employing the QUICK.PRO off-axis minority heating code developed at Lodestar under a D.O.E. Phase I SBIR grant, solutions to the combined quasilinear and Fokker-Planck equation were examined for several cases. In nearly all cases a bi-Maxwellian was found to be an inadequate model because it did not properly account for the trapped particles reflecting in the vicinity of the minority resonance or the rf interactions of particles on flux surfaces nearly tangent to the minority resonance. At present it is not clear how to evaluate the plasma response function for a highly anisotropic distribution function with a complex pitch angle dependence in a numerically tractable and useful manner. Important physical effects associated with the severe pitch angle dependence may not be adequately described by a bi-Maxwellian approximation to the anisotropy. Ultimately what is required is a plasma response function for the actual minority distribution function obtained from a reliable quasilinear code which in turn is coupled to a wave propagation code which self-consistently evaluates the fields for the same plasma response function.
C. Anisotropic Minority Driven Instability

1. Introduction

Motivated by the possibility that minority ICRF heating of tokamak plasmas may give rise to unstable field fluctuations, we have solved the linear dispersion relation for a background deuterium plasma, including a tenuous but highly anisotropic minority-$^3$He component. Although both of the low-frequency normal modes, the Alfvén-Ion Cyclotron (AIC) wave and the Whistler or fast magnetoionic wave, may be destabilized by the helium, only the Whistler is unstable for parameters typical of a fusion reactor. The Whistler is de-stabilized near harmonics of the helium gyrofrequency and for angles of wave propagation to the magnetic field ranging from zero to ninety degrees, depending on harmonic number. For parameters appropriate to the JET reactor, typical growth rates are several percent of the $^3$He gyrofrequency and instability is found near each of the first 70 harmonics. We have studied the effect that a slowing-down distribution of alpha particles has on the instability. A sufficiently high concentration of alphas will quench the instability only if the phase velocity of the unstable wave is less than the alpha birth velocity.

Instabilities such as those studied here are related to the classic Weibel instability$^{35}$ of an unmagnetized plasma, in which the electrons are anisotropically distributed in velocity, and have since been discovered in magnetized plasmas for anisotropic majority ion distributions as well.$^{36}$ Relaxation has been studied numerically and analytically, using quasilinear theory, and proceeds so as to alleviate the anisotropy. In one case, the saturated state includes some remnant anisotropy sustained by self-consistent coherent waves.$^{37}$ Our case differs from those studied previously in that a minority ion species is responsible for the instability. We treat the minorities as a de-stabilizing perturbation of one of the normal modes of the background plasma.

The anisotropic distribution function can provide a source of free energy to fuel a potential instability. The minority ($^3$He) distribution function is constant on ellipses given by $v_{\perp}^2/r + v_{//}^2 = \text{constant}$, where $r = T_{\perp}/2T_{//}$ measures the degree of anisotropy, with $r = 1$ being the isotropic case. But the kinetic energy is constant on concentric circles in this space. Thus for $r > 1$, which is the appropriate choice for modeling ICRF heating, it is possible for a $^3$He ion to lose energy (predominantly in the perpendicular degrees of freedom) to a wave and move into a less populated region of velocity space. The inverse process, one in which initial and final particle states are interchanged, is equally probable. But the net rate of absorption is negative because there are more ions in the higher energy state. So the wave amplitude grows due to
This particular interaction. However the wave is also Landau-damped by the background plasma and by other $^3$He interactions. Only if the net transfer of particle kinetic energy to the wave is positive will there be instability and wave growth.

To include such effects as Landau damping accurately, we solved the dispersion relation (numerically, using Newton's method in the complex $\omega$-plane) for wave propagation in the homogeneous, warm magnetized deuterium plasma with anisotropic $^3$He minority, and with slowing-down alphas where indicated. We now summarize our results for the two low-frequency normal modes of the unperturbed plasma.

2. Dispersion Relation and Born Approximation

We assumed $r = 400$, corresponding to $T_\perp = 3.2$ MeV and $T_{/\parallel} = 4$ keV, and modeled the helium using a bi-Maxwellian distribution function. This high degree of anisotropy is consistent with experiments performed at JET. The background plasma was modeled using isotropic Maxwellian distribution functions for the electrons and for the deuterium ions, each at a temperature of 12 keV, unless stated otherwise (e.g., in Fig.(1)). All species were taken to be spatially homogeneous. Notice that we define the temperature so that isotropy corresponds to $T_\perp = 2T_{/\parallel}$; for example, if $T_{/\parallel} = 4$ keV then the corresponding isotropic temperature is 12 keV.

The well-known dispersion relation for a wave in this plasma is found by solving the linearized Vlasov-Maxwell equations for the fluctuating electric field vector. A linear homogeneous system of equations relating the Fourier coefficients of the electric field vector results which has a non-trivial solution provided that the determinant of the coefficient matrix vanishes. The vanishing of this determinant may be expressed as follows.

$$\det \left[ \varepsilon(k,\omega) - \left( \frac{k_c}{\omega} \right)^2 \left( 1 - \frac{k}{k_c^2} \right) \right] = 0,$$

where the dielectric tensor includes contributions from all of the charged particle species in the plasma,

$$\varepsilon = 1 + \sum_\sigma \varepsilon^\sigma$$

and each bi-Maxwellian species contributes
\[ \varepsilon^\sigma(k, \omega) = -\frac{2}{\omega_{p\sigma}^2} \sum_{m=-\infty}^{+\infty} \left[ \frac{\omega - m\Omega_\sigma}{\omega - m\Omega_\sigma} (1 - W(Z_m^\sigma)) + (1 - r_\sigma) W(Z_m^\sigma) \right] \pi(\beta_\sigma, Z_m^\sigma; m) + \hat{e}_r \hat{e}_z \frac{\omega_{p\sigma}}{k_1 T_{-}^{\sigma}} , \]

where \( Z_m^\sigma = \frac{\omega - m\Omega_\sigma}{|k||T_{-}^{\sigma}/m\sigma|^{1/2}} \), and \( \beta_\sigma = \frac{k_{-}^{2} T_{-}^{\sigma}}{2\Omega_\sigma} \).

In these expressions, \( \sigma \) labels particle species so that only for the helium is \( r_\sigma \) different from one. The matrix \( \pi \) depends on \( \beta_\sigma \) through modified Bessel functions. \( W \) is closely related to the plasma response function, and its argument \( Z_m^\sigma \) is in general complex. See Ichimaru39 for complete definitions. Notice that we have written \( \varepsilon \) in a Cartesian coordinate system in which the wave vector lies in the \( x-z \) plane and the ambient magnetic field points in the positive-\( z \) direction. When alpha particles are included, their contribution to the dielectric tensor is derived from the isotropic slowing-down distribution.40

In a warm plasma, a complex frequency is required to solve this dispersion relation. If it is assumed at the outset that the imaginary part of the frequency is much smaller than the real part, then an approximate solution may be found by assuming a real frequency and disregarding the imaginary part of the dispersion relation. Then the growth (or damping) rate is calculated by substituting this real frequency into the expression for the rate of Ohmic heating by the wave at (real) frequency \( \omega \). Using this method, the simple expression for the Landau damping of electrostatic plasma waves in the limit of long wavelengths may be derived, for example. We refer to this method of solution as the Born approximation.

The rate of Ohmic heating is related to the growth rate of the instability,

\[ \text{Re}(\mathbf{E}^* \cdot \mathbf{J}) / (|\mathbf{E}|^2 + |\mathbf{B}|^2) / 8\pi = -2 \cdot \text{Im}(\omega) \]

and is calculated in the Born approximation from the anti-hermitian part of the dielectric tensor \( \varepsilon \) evaluated, in our cases, using the real frequency of either the cold AIC wave or the cold Whistler for a given wave vector:
\[ \text{Re}(E^*J) = -i(\omega/4\pi)E^*\vec{\varepsilon} \cdot \vec{E} \]

The Born approximation is reasonably accurate near threshold where, by definition, the imaginary part of the frequency is small. However, the real dispersion must be treated accurately for the approximation to be quantitatively useful, and this proves challenging for the AIC wave and for the Whistler at phase velocities near the deuterium thermal speed where the simple cold plasma dispersion relation is invalid. As strong thermal effects must be included in the zero-order dispersion relation, simple expressions for the real frequency lose their accuracy and the Born approximation loses its computational advantage over solving the full dispersion relation. Nevertheless, we find the approximation, based on the assumption of a cold background plasma, useful for providing good guesses at the locations of roots of the full dispersion relation.

We solved the full dispersion relation numerically using Newton's method in the complex \( \omega \)-plane and confined our attention to a particular normal mode by seeding the search with that mode's frequency for a given (real) wave vector. Usually, the cold approximation sufficed for locating the mode even in the warm plasma so that initial seeds tended to be real. Occasionally, the Born approximation was used, along with Ohm's law, to provide seeds not on the real axis. The root finder bootstraps along a branch of the dispersion relation once it has found two nearby roots. It stops when successive approximations to a root fail a Cauchy convergence test (as determined using a preset tolerance) or when we stop it, for example, when the branch enters the lower-\( \omega \) half-plane where the wave is stable.

3. The AIC Wave

Consider first the AIC wave for propagation parallel to the background magnetic field. The instability is similar to one studied earlier but with the important difference that here it is a minority anisotropic distribution of ions that causes an instability due to the resonant interaction between the helium ions and the background wave. It is easy to calculate the rate of Ohmic heating in the Born approximation, ignoring the effect of the helium on the wave dispersion.

In the Born approximation, we assume that the usual cold magnetized plasma dispersion relation holds for the AIC wave and calculate the rate of Ohmic heating as the wave approaches resonance at the \( ^3\text{He} \) cyclotron frequency, \( \Omega_H^c \). (The unperturbed wave is resonant, \( k \to \infty \), at the deuterium cyclotron frequency, \( 3/4\Omega_H^c \).) If the electrons and deuterium are presumed cold then only the helium contributes to \( \vec{\varepsilon} \) and, enforcing the left-hand circular polarization of the cold AIC wave, we find that
\[ \text{Re}(E^*J) \sim (\omega - (1-1/r)\Omega_H) \exp\left[-(\omega-\Omega_H)^2 / 2k^2(\tau_{\text{mH}}) \right] + \\
+ (\omega + (1-1/r)\Omega_H) \exp\left[-(\omega+\Omega_H)^2 / 2k^2(\tau_{\text{mH}}) \right] \]

where we have suppressed a strictly positive coefficient. Notice that a negative rate of heating, and therefore wave instability, is possible whenever \( \omega < (1-1/r)\Omega_H \). (This observation generalizes: the instability is located just below \( N\Omega_H \), \( N = 1,2,... \), for non-parallel propagation, as described below.) But in order for there to be an instability, the destabilizing helium contribution must exceed the dissipative contributions of the deuterium and electrons. In fact, evaluating this expression using the zero-order AIC dispersion relation we find that if \( r > 10 \) then the Born approximation predicts instability for all frequencies less than \( 0.67\Omega_H \); if \( r > 100 \) then there is instability at all frequencies less than \( 0.74\Omega_H \), etc. (The AIC mode does not exist for frequencies larger than \( 0.75\Omega_H \) to zeroth order, i.e., neglecting the helium contribution to the cold dispersion relation.)

For the parameters of a fusion reactor, this approximate result is badly flawed by the assumption of cold deuterium. Solving the full dispersion relation numerically, including warm electrons at 12 keV and \( n_H / n_e = .01 \), we find that for deuterium temperatures greater than about 1.2 keV the AIC wave is stable. For smaller values of \( r \) this threshold temperature is lower. These results are summarized in Fig.(1). The growth rate of the AIC wave is maximized for parallel propagation. Thus, for \( r = 400 \) and \( T_D > 1.2 \) keV there is no AIC instability. (For low deuterium temperatures, we have also observed unstable sound waves propagating at non-zero angles to the magnetic field. But this wave too is easily stabilized by warm deuterium and would not be observed unstable in a fusion plasma.) Next we consider the Whistler wave for realistic fusion reactor parameters.
Fig. (1) The growth rate of the $^3$He-perturbed AIC wave, for propagation parallel to the background magnetic field, as a function of frequency for various deuterium temperatures. For temperatures greater than about 1.2 keV there is no instability. Parameters were chosen to simulate a JET ICRF-heated plasma: $B_0 = 34$ kG ($\Omega_H = 34.5$ MHz), $n_e = 9 \times 10^{13}$ / cm$^3$, $n_D / n_e = .98$, $n_H = .01$, $T_e = 12$ keV = $T_{//H}$, $T_{\perp H} = 3.2$ MeV ($r = 400$).

4. The Whistler Wave

The fast magnetosonic or Whistler wave is unstable only for angles of propagation strictly between $0^\circ$ and $90^\circ$ to the magnetic field. For parallel propagation, the polarization is right-hand circular so that the wave cannot resonate with the left-hand gyrating helium ions. For perpendicular propagation, the Whistler corresponds to the extraordinary wave that is resonant at the lower hybrid frequency. It is easily shown from the dispersion relation that the anisotropy affects only the dispersion of the ordinary wave at $90^\circ$; the Whistler is unaffected. An instability of the perpendicular propagating ordinary wave can be provoked by anisotropy opposite to that considered here (i.e., $T_\perp < 2 \cdot T_{//}$ or $r < 1$) in the ion distribution function of a two-component plasma.\textsuperscript{41} So we restrict our attention to propagation strictly between $0^\circ$ and $90^\circ$ to the magnetic field.

We plot the Born approximation to the growth rate of the $^3$He-perturbed Whistler in Fig. (2) for propagation at $85^\circ$ to the magnetic field and up to the 20th harmonic of the helium cyclotron frequency. In this case we include the contributions of the 12 keV
electrons and deuterium to the anti-hermitian part of the dielectric tensor in implementing Ohm's law. The instability persists at least through the 80th harmonic, in this approximation. At this angle, the lower hybrid resonance is near the 163rd harmonic, beyond which there is no Whistler. But the Born approximation badly over-estimates the growth rate and the extent of the instability in frequency.

Fig.(2) The growth rate of the $^3$He-perturbed Whistler wave in the Born approximation vs. frequency for propagation at 85° to the background magnetic field. Parameters are given in Fig.(1) but $T_D = 12$ keV.

Near the fundamental, the Whistler is destabilized by the helium over the widest range of angles of propagation - virtually all angles between 0° and 90°. The growth rate, obtained by solving the full dispersion relation numerically, is plotted for several angles in Fig.(3). As we seek instability at progressively higher harmonics, the Whistler dispersion relation forces us to use larger values of $k$ so that Landau damping on the deuterium ions squelches the instability at all but the larger angles of propagation. At the tenth and higher harmonics, there is no instability for angles less than 50°. But at 85° the instability persists at least through the 70th harmonic. With increasing harmonic number, the instability is constrained to lie in a very narrow wedge (±2°) of
$k$-space, approximately centered on $87^\circ$. For the parameters used in Fig.(3), the (cold) lower-hybrid resonance is at $\omega_{LH} = 101 \cdot \Omega_H$ for propagation at $87^\circ$.

Fig.(3) The growth rate of the $^3$He-perturbed Whistler wave just below the helium cyclotron frequency obtained by solving the full dispersion relation numerically for various angles of propagation. Parameters are as in Fig.(2).
We have observed that in a neighborhood of maximum growth near a given cyclotron harmonic the helium perturbation can be so strong as to create a new branch of zeroes of the dispersion relation distinct from the Whistler branch. As the angle of propagation increases from zero, the Whistler branch is first tangent to the real \( \omega \)-axis (at threshold) then crosses the axis at two distinct frequencies between which there is instability. With increasing angle the unstable interval increases in length as the maximum growth rate increases. Then at some angle this unstable Whistler branch bifurcates into two distinct branches such that it is no longer possible to move continuously from the weakly damped mode at frequencies farther below \( N\Omega_H \) onto the unstable branch at frequencies closer to \( N\Omega_H \) (e.g., by numerically bootstrapping along using incremental changes in \( k \)). With increasing angle the unstable branch reunites with the Whistler, the growth rate decreases, and the interval of instability shrinks until the mode is once again stable near 90°. This bifurcation is illustrated in Fig. (4a) at the fundamental and in Fig. (4b) at the 10th harmonic where the Born approximation to the growth rate is also plotted for comparison. Between bifurcations there is a new normal mode of the plasma, closely associated with but distinct from the Whistler.
Fig.(4) The growth rate of the $^3$He-perturbed Whistler propagating at a) $60^\circ$ near the fundamental and at b) $80^\circ$ near the 10th harmonic. In both cases, an unstable helium branch has broken off from the Whistler branch. In b) we compare the solution of the full dispersion relation to the prediction of the Born approximation.
5. Stabilization by Alpha Particles

To simulate the progressive stabilization of the Whistler by alpha particles in a burning plasma, we have solved the dispersion relation for successively greater concentrations of alphas using the (isotropic) slowing-down distribution function, with fixed birth and critical velocities, to describe them. For the slowing-down distribution, the anti-hermitian part of the dielectric tensor vanishes for phase velocities greater than the birth velocity (for real frequencies); there is no Landau damping and therefore no enhanced stabilization of the Whistler. Since the cold Whistler phase velocity \( \sim \cos(\phi) \) at high frequencies, the alphas are most effective at stabilizing the higher harmonics where the instability is confined, without alphas, to \( \phi \sim 90^\circ \). For example, the 70th harmonic is stabilized by \( n_\alpha \geq 0.005 \, n_e \) for all \( \phi \) if \( n_H = 0.01 \, n_e \), whereas \( n_\alpha \geq 0.07 \, n_e \) is required to stabilize the 10th harmonic at \( \phi = 80^\circ \). No concentration of alphas that is reasonably weak, so as not to profoundly alter the zero-order dispersion relation, can stabilize the Whistler at the 10th harmonic if \( \phi = 60^\circ \), because the Whistler phase velocity exceeds the alpha birth velocity at this angle. See Fig. (5).
Fig.(5) The growth rate at the tenth harmonic as a function of alpha particle concentration for propagation at a) 80° and at b) 60°. Parameters as in Fig.(2). Notice the slight change in $\gamma$ at 60° resulting from the change in the real part of the dispersion relation with increased $n_a$. 
The Whistler is strongly absorbed between adjacent harmonics. See Figures (2 and 4b). Thus, if low-frequency fluctuations are responsible for saturating the instability by scattering unstable waves into nearby sinks in \((\omega, k)\)-space, we anticipate that a sharp stationary power spectrum of Whistler waves concentrated near the \(^3\)He cyclotron harmonics will result from the instability. As the plasma burns, spikes in the power spectrum diminish first at the higher harmonics and larger angles of propagation. This could prove useful as an alpha particle diagnostic, for example, in the context of gyrotron scattering off electric field fluctuations.40

5. Conclusion

As a first step in studying the stability of ICRF-heated fusion plasmas, we have demonstrated the destabilizing effect of a tenuous, highly anisotropic bi-Maxwellian distribution of \(^3\)He ions on the low-frequency normal modes of a warm magnetized deuterium plasma. For parameters typical of a fusion reactor, we find that the AIC wave is always stable due to strong Landau damping on the deuterium. However, for the degree of anisotropy assumed, the Whistler or fast magnetosonic wave is unstable near each of (at least) the first 70 helium cyclotron harmonics at angles of propagation that approach 90° with increasing harmonic number. We have demonstrated the stabilizing effect that slowing-down alphas can have on waves with phase velocities less than the birth velocity.

A by-product of our research is the code STAB which solves the dispersion relation for wave propagation in a warm, uniformly magnetized multi-component plasma using Newton’s method in the complex \(\omega\)-plane. Once two poles are isolated starting from cold-plasma or Born-approximate first guesses, STAB bootstraps along a branch of solutions of the full dispersion relation. We anticipate generalizing STAB to handle anisotropic minority distribution functions other than bi-Maxwellians, consistent with our evolving understanding of the nature of ICRF heated plasmas, as described elsewhere in this report.

D. Conceptual Design of an ICRF Antenna for Ignitor by J. Jacquino: under the auspices of Lodestar Research Corporation

Summary

i) The design uses the recessed walls and port geometry drawings provided by B. Coppi. Following a discussion with him, the sides of the recess have been fanned to decrease the importance of the RF return currents which are deleterious for coupling.
The toroidal opening is now 85 cm wide. The depth of the recess and the port size are unchanged.

ii) I propose a broad frequency range: $100 < \nu < 200 \text{ MHz}$. This allows operation at $2 \omega_{CT}$ (140 MHz) as foreseen previously but also at $\omega = \omega_{CD}$ and $\omega = \omega_{CH}$ allowing interesting minority heating scenarios as well as operation below the maximum design field at $\omega = 2\omega_{CT}$.

iii) The FREMIR wave absorption code gives 52% absorption per pass at $2 \omega_{CT}$ for the final plasma parameters ($n_{eo} = 10^{21} \text{ m}^{-3}$, $T_{eo} = T_{io} = 10 \text{ keV}$, $k_i = 10 \text{ m}^{-1}$). Direct electron heating can be important depending on the launched spectrum. It reaches 50% of the coupled power for $k_i = 10 \text{ m}^{-1}$. The other half goes to the tritons with only 4% going to the $\alpha$'s. The single pass damping is low during the start-up phase.

At 200 MHz ($\omega = \omega_{CH}$), the damping is in excess of 70% and remains large during the start-up phase.

Both heating scenarios can be considered as appropriate with damping per pass similar to the normal JET operating regimes.

iv) Each antenna module is composed of 4 straps (12.5 cm wide, 40 cm long). The length of the straps is chosen to be less or comparable to the $\lambda/4$ resonance in the middle of the frequency range. The straps are grouped in poloidal pairs with a toroidal separation of 32.5 cm between strap centers. The strap currents are in-phase poloidally and out-of-phase toroidally. The structure creates the toroidally antisymmetric pattern which eliminates ICRF impurity effects. Angled beryllium bars form a Faraday screen ($\phi = 20 \text{ mm spaced every 30 mm}$). The antenna structure does not protrude from the outer wall.

v) The JET coupling code shows that 4 MW can be coupled by each antenna module provided:

a) the antenna withstands 40 kV

b) the separatrix is less than 4 cm away from the screen outward face.

The antenna is wide band (thanks to the 4 strap concept) with an optimum coupling around 140 MHz.

vi) I propose that the recessed wall and the vertical sides of the port carries the return RF currents. The concept is possible for the first time in a tokamak because the port and the vessel shape have been designed with ICRF in mind. This considerably simplifies the mechanical aspects of the antenna: The antenna short circuits are simply bolted on attachment blocks welded to the vessel and the outer conductors of the vacuum transmission lines have a square cross-section made with the vertical sides of the port and with additional horizontal plates welded between the port sides.
vii) Each antenna strap is fed by a separate coaxial line. The four vacuum transmission lines nearly fill up the port. Vacuum pumping is performed by the torus itself through the screen. Transition from rectangular to circular cross-section of the outer conductors of the transmission lines is performed at the port flange. A double conical feedthrough connects a 9 inch pressurized line to each vacuum transmission line near the port flange. The antenna module is driven by 4 amplifiers (1.5 MW each) phased to achieve the required dipole radiation pattern. This method is intrinsically wide-band. It allows automatic matching by frequency and tuning stub variation. It also allows arbitrary phasing to modify the toroidal spectrum. For instance 90° phasing will give wave directivity ranging from 30 to 70% depending on the density profile near the antennas. It is not clear at present whether this wave directivity can be used to drive electron or ion current in the high density regime foreseen in Ignitor.

E. Energetic Electron Transport Studies

During the grant reporting period, several aspects of energetic electron transport in tokamaks have been investigated. The motivation for these studies arises from a number of both practical and scientific considerations. One important practical consideration is that the confinement of suprathermal electrons is a key issue in the design of a steady steady tokamak. To efficiently drive current with suprathermal electrons they must remain spatially confined. Thus, the size of future devices is inextricably linked to suprathermal electron confinement. A separate, but related, practical reason necessitating an understanding of energetic electron confinement is in the context of runaway production during major disruptions in a tokamak reactor plasma. Runaways are a concern because of the damage they cause to plasma facing components. It is believed that runaway electrons are produced by parallel electric fields during an MHD disruption. The ultimate energy attained by these runaways is a function of their confinement time in the vicinity of the accelerating fields.

As a scientific tool for exploring important confinement physics for tokamaks, suprathermal and runaway electrons have proved to be very valuable. Studies at Lodestar, and in collaboration with TEXT, have investigated the energy dependence of electron transport. It was shown\textsuperscript{42} that a quasilinear theory containing both magnetic and electrostatic $E\times B$ driven radial transport is consistent with the available experimental data, and supports the notion that magnetic turbulence dominates runaway diffusion but electrostatic turbulence dominates thermal diffusion in TEXT. The conclusion that magnetic transport is irrelevant for explaining thermal electron transport in TEXT has been further substantiated by a recent compilation\textsuperscript{43} of a variety of experimental
data including: transient responses to plasma shifts and externally induced stochasticity,\textsuperscript{44-46} pellet ablation rates,\textsuperscript{47} and sawtooth pulse propagation experiments,\textsuperscript{48} (to be discussed separately below) all of which can be used to deduce magnetic fluctuation levels. During the past year, work on several remaining aspects of these topics was completed and reports have written with our collaborators.\textsuperscript{42,43,45,46} The observations made in these studies further expand the possibilities for employing runaway or suprathermal electrons as a diagnostic of the underlying transport mechanisms in a tokamak plasma.

In addition to our collaborations with TEXT, a collaboration with H.E. Mynick of PPPL has resolved an apparent difference between two calculations of drift-orbit modifications to the spatial diffusion rate of particles in stochastic magnetic fields.\textsuperscript{49,50} The apparent discrepancy occurred in the limit of radially broad modes so that radial drifts became unimportant and only poloidal drifts of the energetic particles needed to be considered. The calculations were reconciled by noting the relevance of a previously unexplored inequality which involved the spectral width of the turbulence relative to a finite Larmor radius parameter. Specifically, it was found that the results of Ref. 49 are recovered in the limit of broad spectra \( b \ll w \) where \( b \sim \rho_p/R, w = qR\Delta k_{\parallel} \), \( m \) is the poloidal mode number of the turbulence, \( \rho_p \) is the poloidal Larmor radius, \( R \) is the major radius, \( q \) is the safety factor and \( \Delta k_{\parallel} \) is the spectral width of the turbulence. In the opposite limit \( w \ll b \), the results of Ref. 50 are recovered. A technical report is in progress. Our previous estimates\textsuperscript{49} of drift-orbit modifications to the spatial diffusion of runaway electrons in TEXT remains unaffected by these more recent studies.

One of the difficulties in interpreting the runaway electron confinement experiments discussed at the beginning of this section arises because the source and energy distribution of runaways is not well understood on the longer time scales characteristic of the entire discharge. There are observations and arguments\textsuperscript{44} which suggest a continuous creation of runaways throughout the discharge; but there are also counter-\textsuperscript{44}ing arguments which lead to a model in which the runaways are produced at the initiation of the discharge. To interpret longer time scale experiments one must reconcile these different source scenarios with the observed energy distribution and confinement data. This challenge was one of the initial motivations for beginning a study of the runaway distribution function which included the important physics of acceleration by the induction field, Coulomb collisions and spatial confinement.

A second motivation arises from the consideration of pellet injection experiments. Pellet injection experiments on TEXT show that the ablation rate is sensitive to the presence of suprathermal and/or runaway electrons.\textsuperscript{47} Theoretically, it is expected
that the ablation rate should be controlled by the parallel electron heat flux impacting the pellet.

At Lodestar, we are in the process of developing an analytic model for the moments of the energetic electron distribution function which includes the crucial physics mentioned in the preceding two paragraphs. An equation for the one dimensional distribution function $f(v_{\parallel})$ (velocity parallel to B) has been derived and solved in terms of parabolic cylinder functions which are matched smoothly to a Maxwellian $f$ at low velocities by means of a WKB technique. The theory includes the usual critical speed at which drag and acceleration balance\textsuperscript{51-53} as well as a characteristic velocity of the highest energy runaways for which spatial confinement and acceleration balance.\textsuperscript{54}

The velocity moments of $f$ are dominated by the tail for sufficiently large powers of velocity $v_{\parallel}^p$. In TEXT, the density and current moment ($p = 0, 1$) are dominated by the Maxwellian, but the runaway tail may contribute to the electron beta ($p = 2$). The tail dominates the parallel heat flux ($p = 3$) in the pellet ablation experiments of interest, and it also dominates the contribution to the hard X-ray flux ($p \approx 5$) employed previously\textsuperscript{42,44-46} in the transient transport experiments.

Work on the runaway distribution function studies is still in progress. It is expected that the results will enable at least a partial resolution of the ambiguities associated with the runaway electron source in the confinement experiments. Additionally, as a result of this work, several novel techniques for employing pellets to conduct energetic electron transport studies\textsuperscript{47} may become amenable to more sophisticated theoretical modeling.

The Lodestar work on energetic electron transport in tokamaks described in the present subsection has integrated well with separate grant work done under the Small Business Innovation Research (SBIR) Program. In the SBIR study,\textsuperscript{48} it was demonstrated from TEXT data that the sawtooth crash could be used as a perturbation of both the thermal and runaway electrons, thereby enabling a simultaneous measurement of the thermal diffusivity and runaway diffusion coefficient. The thermal and runaway data points for the energy dependence of electron confinement are consistent with the notion that in TEXT thermal confinement is controlled by an $E\times B$ transport mechanism while energetic confinement is controlled by magnetic fluctuations, a feature that arises naturally in the theoretically modeling developed in Ref. 42 under the present grant. Further implications of these results are being explored as time permits.
References


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