HIGH BETA CAPTURE AND MIRROR CONFINEMENT OF LASER PRODUCED PLASMAS

MASTER

ALAN F. HAUGHT
RICHARD G. TOMLINSON
WILLIAM B. ARD
LAURENCE R. BOEDEKER
THOMAS L. CHURCHILL
WALTER J. FADER
RAYNARD A. JONG
ARTHUR E. MENSING
DONALD H. POLK
JOHN H. STUFFLEBEAM

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Reported By:

A. F. Haught
Manager
W. J. Fader
R. G. Tomlinson
R. A. Jong
W. B. Ard
A. E. Mensing
L. R. Boedeker
D. H. Polk
T. L. Churchill
J. H. Stuffebeam

Approved by:

Arthur S. Kesten
Manager of Energy Research

Date: December 31, 1977

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Final Report Under Contract EY-76-C-02-2277.*000
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High- β Capture and Mirror Confinement
of Laser Produced Plasmas

The LITE fusion plasma research program at UTRC has been investigating the stabilization and confinement physics of a mirror plasma created by energetic neutral beam heating of a confined target plasma. The dense mirror confined target plasmas for these experiments are produced by laser irradiation of solid LiH particles electrically suspended in vacuum at the center of a minimum-B magnetic field or by injection from a Ti washer gun. An energetic neutral beam is employed to heat the plasma by charge exchange collisions and, through ionization of beam neutrals, to provide a source of hot ions to counteract the loss of particles from the mirror field by decay processes. During the period covered by this report work has been concentrated on the investigation of hot ion losses in a warm target plasma, development of a cryocondensation pump for the LITE beam line neutralizer, theoretical studies of ECRH modification of the ambipolar potential in mirror plasmas, and analysis of the effects of localized cold plasma on DCLC stabilization. The results of these investigations are summarized below and detailed in four papers which comprise the body of this report.

Measurements of the lifetime of hot ions in a mirror confined warm plasma have been carried out by observations of the hot ion buildup time obtained with energetic neutral beam injection. The background vacuum was sufficiently low as not to be a limitation on the hot ion plasma lifetime in these experiments. The hot ion lifetime scales inversely with the warm target plasma density and thus is determined by interaction with the target plasma or accompanying cold neutrals from the washer gun source. Assuming electron drag as the primary interaction gives an unreasonably low electron temperature of $\lesssim 1$ eV for the warm target plasma. Observations of scattered neutrals at intermediate energies and of the variation in the scattering magnitude with decay time suggest that the hot ion lifetime may be due to other mechanisms than simple electron drag.

A cryocondensation pump of novel design has been constructed and incorporated in the neutralizer chamber of the LITE neutral beam line. Both the helium and nitrogen shield panels are constructed using extruded members with integral cooling tubes which greatly reduce the number of

welds required. Special brazed joints are used for the transition between the aluminum panels and the respective helium and nitrogen dewar reservoirs. The cryocondensation pump will be used to provide increased pumping speed in the LITE beam line whose length has been shortened for greater beam intensity and for experiments with a reverse flow neutralizer being developed for minimum cold neutral contamination of the beam.

Calculations have been carried out to evaluate the sizes and shapes of ambipolar potential modification produced by electron cyclotron resonance heated electrons and to determine the spatial distribution and densities of cold ions trapped in the potential wells. The calculations employ a velocity distribution for the hot ECRH electrons derived from a resonant diffusion model of cyclotron heating and include the self-consistent effect of the trapped cold ions on the potential distribution. The results obtained show that, as has been suggested by Ioffe, ECRH electrons can produce an electrostatic potential well of sufficient depth to trap an appreciable fraction of cold ions in a mirror confined plasma.

The effects of the spatial distribution of the cold ions on their effectiveness for stabilizing the drift cyclotron loss cone instability has been studied numerically using the formulation of Pearlstein in which the dispersion relation for the DCLC mode is solved for finite-sized plasmas containing hot and cold components. These calculations indicate that the trapping of small amounts of cold ions near the midplane of a mirror confined plasma can reduce the amount of uniformly distributed cold plasma required to stabilize the plasma against the DCLC instability. Furthermore, the calculations show that cold ions trapped on the midplane are more effective than cold ions trapped off of the midplane.

SECTION I

HOT ION LIFETIME IN WARM PLASMA BACKGROUND

by

A. F. Haught, W. B. Ard, D. H. Polk, J. H. Stufflebeam
and R. G. Tomlinson

HOT ION LIFETIME IN WARM PLASMA BACKGROUND

I. INTRODUCTION

The LITE facility was developed to examine the generation and properties of mirror confined plasma resulting from energetic neutral beam heating of a dense target plasma which is produced by laser irradiation of solid LiH particles electrically suspended in vacuum at the center of a minimum-B mirror magnetic field. An alternative, which provides a pure hydrogen target plasma at the expense of increased neutral background gas density, is to produce the target plasma by capture of streaming plasma from a hydrogenated Ti washer gun located outside the mirror field on the field axis. This paper describes the results obtained in initial experiments to examine the hot ion buildup and loss in the confined plasma produced by injection of a 30 kV neutral hydrogen beam into a washer gun target plasma of $\sim 10^{13} \text{ cm}^{-3}$ density. The apparatus used in these experiments is described in Section II. In Sections III and IV, a simple analytical model of the injection beam/target plasma interaction is described and the experimental results are compared to the predicted behavior. The results of these comparisons are summarized in Section V.

II. EXPERIMENT SYSTEM

The experiment system, shown in Fig. 1, consists of a central vacuum chamber containing the minimum-B mirror magnet coil and a titanium washer gun mounted on axis above the magnet, the neutral beam source, and a beam dump

tank. Initial evacuation and cw pumping of the vacuum chambers at pressures $\lesssim 5 \times 10^{-8}$ torr are provided by three 10-inch diffusion pumps on the beam neutralizer chamber, two 6-inch pumps on the experiment chamber, and a single 10-inch pump on the beam dump tank. Prior to experiment runs, titanium getter elements in the top and bottom sections of the experiment chamber and in the dump tank can be flashed to coat the surfaces of liquid nitrogen dewars lining the cylindrical chamber walls. With these surfaces acting as getter pumps, a base pressure of $2-4 \times 10^{-9}$ torr is achieved. In addition, a titanium ball element can be inserted into the magnet to coat the interior surfaces of the magnet shell with getter material to provide pumping for transient gas loads during the experiment runs.

To remove large divergence components of the neutral beam and to reduce the streaming gas load from the source on the volume of confined plasma inside the magnet, beam defining apertures sized to pass only the focused components of the beam are placed at the end of the neutralizer chamber (12 cm dia) and at the outer face of the magnet shell (9.25 cm dia). Reflux from the beam dump is reduced by means of an aperture (17.8 cm dia) and a beam dump plate of water-cooled titanium set at angle of 20° relative to the beam axis. To limit the conductance of neutral gas from the washer gun into the magnet interior, a 0.6 cm wide by 15 cm long rectangular aperture, aligned to pass all field lines passing through the plasma confinement volume, is located at the top of the magnet.

A. Magnet

The confinement magnet is a spherical, baseball seam coil using liquid nitrogen prechilled copper windings enclosed in a vacuum tight 316 stainless steel shell. The minimum-B mirror magnetic field produced has mirror ratio and well depth of 2, distance between mirror points of 31 cm, and a maximum center field intensity of 15 kG. The effective plasma confinement volume of $\sim 400 \text{ cm}^3$ has a diameter of $\sim 6 \text{ cm}$ in the midplane and a cross-section area normal to the neutral beam of $\sim 100 \text{ cm}^2$. At 15 kG center field, the adiabatic energy limit for confinement of hydrogen ions is 25 keV; hence, only the half and one-third energy components of the 27-30 kV neutral beam are adiabatically confined. As a result, for these experiments the magnet was operated at 12 kG center field (16 keV adiabatic limit) to minimize Joule heating of the coil and the consequent cooling time between pulses.

B. Titanium Washer Gun Target Plasma

The initial confined target plasma for these experiments was produced from the streaming plasma from a hydrogenated titanium washer gun mounted on the magnetic field axis above the magnet (Fig. 1). The washer gun, shown in Fig. 2, is driven by a pulse forming network (Fig. 3) which, for these experiments, was configured to deliver a current pulse of $\sim 1 \text{ msec}$ duration. Using a separately triggered ignitron as a crowbar switch, the gun arc current could be terminated at any desired point with a current decay time of $\sim 100 \text{ } \mu\text{sec}$. The confined target plasma produced by the washer gun has ion energies of $\sim 200 \text{ eV}$, electron temperatures of $\sim 10 \text{ eV}$, and density $\sim 10^{13} \text{ cm}^{-3}$.

C. Neutral Beam

The source for the neutral beam is an ORNL 15-cm ion source and cylindrical gas neutralizer cell. The ion source extraction grids, containing an array of 861 holes of 0.375 cm diameter each, were designed to focus the beamlets onto the target plasma. Using a 50-cm long neutralizer, measurements showed that a 10 amps equivalent neutral current of one-half and one-third energy components of a 27 kV beam were focused onto a 100 cm^2 target area at a distance of 230 cm from the source. Annular zone calorimeters on each side of the target region (Fig. I-1) and a 2.5 cm diameter calorimeter probe were used to determine the focal distance and beam current distribution on the target based on extracted ion species ratio of 0.5/0.4/0.1 for H^+ , H_2^+ , and H_3^+ , respectively. For the 50 cm neutralizer, the calculated neutral species ratios are 0.21:0.56:0.23 for full, one-half and one-third energy components. Because more recent measurements at ORNL suggest that the proportion of the full energy component may be higher for this source, an electrostatic energy analyzer to measure the neutral species distribution has been constructed and installed, but no measurements have yet been made.

In the course of these experiments, a recheck of neutral beam current density at the target showed that the current density was approximately a factor of 10 lower than previously measured. In addition, it was found that the beam aperture plate outside the magnet was intercepting $\sim$ two-thirds of the beam power. Upon disassembly, it was found that the curvature of extraction grids had increased sufficiently to shift the beam focus away

from the target plasma toward the beam source by more than 50 cm. Reforming of the grids to the correct curvature is in progress, but, as a result of the change in focus, the target plasma/neutral beam interaction data presented herein were obtained with injection beam currents on the target of only ~ 10 percent of the 10 amp beam capability.

D. Diagnostics

Since the objective of these experiments is to establish a hot ion dominant plasma by the interaction of the high energy neutral injection beam with the confined target plasma, the primary measurements of interest are the time variation of the confined plasma density, the fraction of the neutral beam that interacts with the confined plasma, and the buildup rate and density of hot ions trapped in the mirror field.

A 4-mm microwave interferometer, in which the measuring signal is transmitted by a horn/lens combination across the midplane of the magnetic mirror field to a horn/lens receiver as shown schematically in Fig. 4, provides a measure of the electron density of the confined plasma for $N_e \lesssim 2 \times 10^{13} \text{ cm}^{-3}$. The upper limit of the measuring range is less than the cutoff density for 4 mm microwaves ($7 \times 10^{13} \text{ cm}^{-3}$) because of refraction of the microwaves by the plasma at high densities. The microwaves provide density decay measurements for the cold target plasma, and, for a sufficient reduction in decay rate in the presence of the neutral injection beam, would provide a direct measure of hot ion dominance in the confined plasma.

A secondary emission detector (SED) located at the end of the beam dump tank 300 cm from the plasma (530 cm from the ion source) is used to measure the attenuation of the neutral beam in the presence of the confined plasma and, thereby, provide a measure of the fraction of the beam trapped by the plasma.

The hot ion density is inferred from observations of the high energy neutrals escaping from the plasma as a result of charge exchange losses of trapped hot ions. A silicon barrier detector (SID) of 8 mm diameter aperture located at 1.5 m from the plasma center (Fig. 4) observes the escaping neutrals directly and provides a measure of the total number of fast ions in the plasma. In addition, a small fraction of the escaping neutrals are reionized in a nitrogen gas stripping cell and are energy analyzed in a 90 degree, three channel electrostatic analyzer. The three energy channels are positioned to detect ions having energies $3/2 V$, $3/4 V$, $1/2 V$ where V is the analyzer deflection potential. Hence, if V is set equal to two-thirds of the beam energy, the three detector channels are centered on the full, one-half, and one-third energy components of the neutral beam. An earlier version of the energy analyzer used channel electron multipliers as the detectors, but to provide linear response at the higher particle fluxes of interest in these experiments, the analyzer detectors were changed to silicon barrier detectors of the same type used for the neutral particle detector.

To interpret the measurements in terms of absolute measurements of hot ion population, both the SID and the fast atom energy analyzer can be calibrated by injecting the neutral beam onto a nitrogen gas target in the magnetic field

and relating the observed signal to the calculated rate of production and loss of hot ions on the known density of nitrogen in the target volume viewed by the detectors. The gold coated silicon detectors used have not yet been so calibrated since, in these experiments, only relative signal amplitudes are required in examining the rate of buildup and decay of the fast neutral signal.

III. MODEL

The hot ion plasma density produced by injection of an energetic neutral beam onto a mirror confined warm target plasma is given by

$$\frac{dn_H}{dt} = n_w J \sigma_{cx} - \frac{n_H}{\tau_H} \quad (1)$$

where n_w = warm target plasma density

J = neutral beam injection current density

σ_{cx} = beam charge exchange cross-section on target plasma

τ_H = hot ion loss time

and it is assumed that ionization of the beam on the confined hot ions can be neglected. Since $\sigma_i \ll \sigma_{cx}$, ionization on the warm target ions will always be small compared with $n_w J \sigma_{cx}$. For injection conditions such that $n_H \ll n_w$, with a square pulse injection beam

$$J = \begin{cases} 0 & t < t_0 \\ J & t_0 \leq t \leq T \\ 0 & T < t \end{cases} \quad (2)$$

incident on a steady state (externally maintained) target, the hot ion density has the form

$$n_H = n_w J \sigma_{cx} \tau \left[1 - e^{-\frac{(t-t_0)}{\tau}} \right] \quad t_0 \leq t \leq T$$

$$= n_w J \sigma_{cx} \tau e^{-\frac{(t-T)}{\tau}} \quad T < t \quad (3)$$

The hot ion density is observed using an energy discriminating neutral particle detector to measure the fast atom flux from the plasma. For injection current and vacuum background conditions such that charge exchange scattering by the energetic beam is much greater than that due to the low energy neutrals, i.e.,

$$n_w J \sigma_{cx} V \gg \frac{n_0 V}{4} S \quad (4)$$

where V = target plasma volume

n_0 = cold neutral background density

v = cold neutral velocity

S = target plasma surface area,

then the fast neutral flux to the detector is proportional to

$$F \propto n_H J \sigma_{cx}$$

and the fast neutral detector signal has the form

$$S_F \propto n_w (J \sigma_{cx})^2 \tau \left[1 - e^{-\frac{(t-t_0)}{\tau}} \right] \quad t_0 \leq t \leq T \quad (5)$$

Equation 5 for the fast neutral detector signal is plotted in Fig. 5 along with the target plasma density, the injection beam current pulse (Eq. 2), and the confined hot ion buildup density (Eq. 3). Measurement of the risetime of

the fast neutral signal provides a direct measure of the hot ion lifetime and from Eq. 1 can be used to determine the equilibrium hot ion buildup density

$$n_H = n_W J \sigma_{cx} \tau; \frac{dn_H}{dt} = 0 \quad (6)$$

IV. DISCUSSION OF RESULTS

Typical data for the experiments is shown in Fig. 6. The injection current, measured with a 1 in. dia calorimeter target located at the target plasmas position, was 40 mA/cm², and, for the shot shown in Fig. 6, the background gas density was 3 x 10⁻⁸ torr. From the microwave interferometer measurements, the equilibrium target plasma density established by the washer gun plasma source was 10¹³ cm⁻³, and the charge exchange scattering by the beam

$$n_W J \sigma_{cx} V = 1.6 \times 10^{17} \text{ sec}^{-1}$$

compared with the cold neutral scattering

$$\frac{n_0 V}{4} S = 1.8 \times 10^{15} \text{ sec}^{-1}$$

meets the conditions for beam dominated scattering of Eq. 4. Even considering the effect of the ~ 1 percent cold neutrals accompanying the energetic beam, the cold neutral charge exchange

$$J_0 A = 4.7 \times 10^{16} \text{ sec}^{-1}$$

is less than that due to the energetic beam neutrals.

From the microwave transmission measurements the target plasma density reaches its equilibrium value at ~ 400 μsec. At earlier times the calculated electron density is displayed at larger values but it is clear from the transmission data that this is an artifact and is the result of the inability of the interferometer to accurately track the phase shift at high attenuations as the

electron density is evaluated working backwards in time from the asymptotic value in the plasma decay. At 400 μsec the accel voltage of the injection source is pulsed on. Secondary emission detector measurements show that the injection current rises rapidly in $\sim 6 \mu\text{sec}$ to a value which remains nearly constant throughout the pulse and when turned off at 750 μsec decreases to zero in $\sim 10 \mu\text{sec}$. The washer gun target plasma source is turned off using the same timing pulse as the injection beam termination; however, from the microwave data the target plasma decays on the same time scale as the washer gun arc current, $\sim 100 \mu\text{sec}$, long compared with the rate of decrease of the injection beam.

The fast atom flux is measured with the silicon barrier detector which is sensitive only to particles with energies above $\sim 0.5 \text{ keV}$ and thus selectively discriminates against scattered warm target plasma ions. The silicon detector signal shows a 10-90% risetime of 31 μsec which corresponds to an e-folding time of 14 μsec . Using the experimental e-folding time in Eq. 6 yields $n_H = 3.4 \times 10^{10} \ll n_w$, as assumed in the derivation of Eq. 3. Measurements of the e-folding time have been carried out for different values of background pressure with the results shown in Fig. 7. At pressure below 10^{-6} torr, the e-folding time approaches a limiting value and the background pressure has little effect. Thus, we conclude that for the conditions of these experiments the hot ions within the target plasma are subject to a loss with a time constant of $\sim 15 \mu\text{sec}$, independent of background pressure. Additional measurements have been made at high vacuum with different target plasma densities produced by operating the washer gun at reduced voltage. Although only limited data is available, it is clear from Fig. 8 that the limiting hot ion loss

time is inversely dependent on the warm target plasma density. Assuming a 100 eV ion temperature for the warm target plasma, the charge exchange mean free path for cold background neutrals incident on the target plasma is ~ 1 cm. If the limiting loss time was the result of background neutrals, the screening provided by the target plasma would be reduced for lower target plasma densities, and the inverse dependence of Fig. 8 would not be obtained. It appears, therefore, that the limiting loss is the result of interaction of the hot ions with the warm target plasma or with streaming cold neutrals produced by the washer gun along with the warm plasma.

Because the hot ion component is a small fraction of the warm target plasma density, it may be assumed that the electron temperature is unaffected by the injection and is fixed at the value determined by the washer gun stream. Assuming that the primary interaction of the hot ions is with the electrons of the warm target plasma, i.e., electron drag, the temperature of the target plasma electrons can be determined from limiting loss time. From Spitzer, neglecting the effects of the warm plasma ions, the energy decay time is

$$\tau = \frac{3m_i m_e k^{3/2}}{8(2\pi)^{1/2} n z^2 z_0^2 e^4 \ln \Lambda} \left[\frac{T_e}{M_e} \right]^{3/2} = \frac{2.2 \times 10^7}{n} T_e^{3/2}$$

and the loss time for ions to be dragged into the ambipolar hole is

$$\tau = \tau \ln [E_i / E_n] = \frac{2.2 \times 10^7}{n} T_e^{3/2} \ln [T_i / 3 T_e]$$

using the assumption that the ambipolar space potential $\phi \approx 3 T_e$. For an average ion injection energy of 13.5 keV (half energy beam component),

$T_i = 9$ keV, and setting τ_d equal to the limiting loss time gives an electron temperature of .89 eV. This value of the electron temperature appears unreasonably low, however, since electron temperatures of ~ 10 eV were measured for washer gun target plasmas in earlier experiments using Thomson scattering diagnostics.

Observations were also made of the fast neutral flux in discrete energy channels using the fast atom detector. The signals at the full, half, and one-third energies displayed the same time history as the silicon barrier detector signal with virtually identical risetimes. Preliminary measurements were also carried out at other energies. As expected, for detector energies above the full injection energy no scattered signal was observed. However, in a limited number of measurements at intermediate energies the anticipated time delay for the ions to be dragged down in energy was not found, a result which would be consistent with low values of electron temperature but which may also be the result of overlap of the detection channels with the injection components.

V. SUMMARY

Measurements of the lifetime of hot ions in a mirror confined warm plasma have been carried out by observations of the hot ion buildup time obtained with energetic neutral beam injection. The background vacuum was sufficiently low as not to be a limitation on the hot ion plasma lifetime in these experiments. The hot ion lifetime scales inversely with the warm target plasma

density and thus is determined by interaction with the target plasma or accompanying cold neutrals from the washer gun source. Assuming that the hot ion lifetime is limited primarily by electron drag gives a warm plasma electron temperature less than 1 eV, a factor of 10 lower than obtained in earlier experiments from Thomson scattering measurements. Observations of scattered neutrals from the plasma at energies between the full, half, and one-third energy components of the beam have not shown the initial time delay which would be expected with an electron drag dominated loss with electron temperatures of ~ 10 eV, for which the energy decay time is > 50 μ sec.

A number of puzzling aspects of the data were noted in these experiments. For the conditions of the experiments at high vacuum, it is expected that charge exchange scattering with the energetic beam neutrals provides the primary fast atom signal, and, from the model, the signal should decrease abruptly when the injection beam is turned off. However, a decaying tail is observed on the fast atom signals, both with the silicon detector and with the energy selective fast atom detector, with a decay time constant approximately one-half the time constant measured for the signal risetime. The same factor of two relationship is obtained even with high background pressures and thus is not the result of background accumulation from the washer gun, neutral injection beam, or plasma, which would be independent of ambient vacuum and result in a fixed decay time at low background pressures. This suggests that in these experiments the hot ion lifetime may be due to other mechanisms than simple electron drag. On the other hand, the variation of signal amplitude, displayed

in Fig. 9, generally conforms to the model which predicts a signal proportional to $n_w \tau$.

Additional experiments are planned to investigate the lifetimes of hot ions in a warm plasma background. These experiments will involve higher energetic beam injection currents to provide a more nearly beam dominated charge exchange scattering; a factor of 10 greater beam currents is available with improved beam focusing. Silicon fast neutral detectors which view the entire plasma volume and measure the fast neutral flux off the midplane to take into account pitch angle scattering will be incorporated in the future experiments and additional studies will be made to resolve the time history of the scattering at intermediate energies to resolve the effects of electron drag on the hot ion lifetime.

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FACILITY FOR 10A INJECTION ONTO A TITANIUM WASHER GUN PLASMA TARGET

FIG. 1

WASHER GUN CROSS-SECTION

PFN ARC POWER SOURCE AND TRIGGER FOR WASHER GUN

CONFIGURATION OF MICROWAVE INTERFEROMETER AND FAST ATOM DETECTOR
(NOT TO SCALE)

VARIATION OF HOT ION DENSITY (n_H) AND FAST
NEUTRAL SIGNAL (S_F) FROM MODEL

TYPICAL DATA FOR NOMINAL 30kV BEAM INJECTION ON WASHER GUN TARGET PLASMA

$$P_0 = 3.0 \times 10^{-8} \text{ TORR} \quad t = 0 \text{ @ WASHER GUN TRIGGER}$$

A. BEAM PARAMETERS

B. MICROWAVE INTERFEROMETER

C. FAST NEUTRAL DETECTORS

D. SID RISE AND DECAY TIME

FIG. 6

RISE AND DECAY TIMES (e-FOLD) OF SILICON DETECTOR SIGNAL AS A
FUNCTION OF INITIAL AMBIENT PRESSURE

FIG. 7

RISE TIME OF SILICON DETECTOR SIGNAL AS A FUNCTION OF TARGET PLASMA DENSITY

FIG. 8

SID SIGNAL AMPLITUDE : VARIATION WITH RISETIME
AND TARGET PLASMA DENSITY

SECTION II

LITE CRYOCONDENSATION PUMP

A. E. Mensing, D. J. McFarlin, R. F. Dondero

LITE CRYOCONDENSATION PUMP

A. E. Mensing, D. J. McFarlin and R. F. Dondero

I. INTRODUCTION

A liquid helium cooled cryocondensation pump was installed in the neutralizer chamber of the LITE neutral beam line to provide increased pumping speed in this chamber and reduce the flux of cold gas from the neutralizer into the experiment chamber. The sources of cold gas are a) the ion source itself and b) a counter flow neutralizer into which hydrogen is injected to raise the pressure (thus increasing the neutralizer efficiency) as well as reducing, through molecular collisions, the directional streaming cold gas flux (Ref. 1). A sketch of the LITE facility is presented in Fig. 1 showing the principle components including the cryopump.

The design for the LITE cryocondensation pump was a modification of that employed in the cryocondensation pumps of PLT beam lines. Two optically shielded liquid-helium-cooled flat panels are suspended vertically in the chamber, one on each side of the beam (see Fig. 2). Each liquid helium cooled panel (Fig. 4) is 28 cm x 58 cm and is shielded from radiation by optically black LN_2 cooled chevrons (see Fig. 3) which permit the flow of gas particles but block all optical radiation. A cross-sectional sketch of the panel-chevron assembly is shown in Fig. 5.

II. DESIGN CHARACTERISTICS

The UTRC cryocondensation pump design incorporates integral liquid nitrogen - liquid helium dewars and two chevron shielded condensation panel assemblies. The dewars, liquid supply lines, valves and panel feed plumbing are constructed of Type 304 stainless steel, whereas the panels and manifold assemblies are constructed of aluminum. The helium system, except for the fill valve, is a completely welded or brazed structure. The nitrogen system is separable at the panel manifold/feed line interface through hard seal flange joints. The pump is instrumented for the determination of helium dewar liquid level and surface temperature at several locations on the chevrons and condensation panels.

The stainless steel dewar assemblies are mounted within a vacuum chamber and are arranged so that the liquid nitrogen dewar forms an annular radiation shield around the cylindrical liquid helium dewar. The dewars are supported by long, high length-to-diameter-ratio rods using low thermal conductivity ceramic insulators at interface points. The helium dewar and supply valve are arranged so that the liquid helium is shielded by a cold gaseous helium layer. In addition, the pintle of the supply valve utilizes machinable ceramic (low thermal conductivity) to minimize heat leaks. The helium dewar is designed to store a charge of liquid, whereas the nitrogen dewar is designed for continuous feed.

The panel assemblies, including their respective manifolds, are constructed of aluminum. The conceptual arrangement of the shielded condensation panel is depicted in Fig. 5. The central helium cooled condensation panel is shielded on both sides by optically dense liquid nitrogen cooled chevron panels. Both the helium and nitrogen shield panels are constructed using extruded members. The helium panel extrusion is a 2.4 mm thick by 15 cm wide plate with a 1.2 cm I.D. integral cooling tube. The liquid nitrogen cooled chevron extrusion includes an integral 3.8 mm inch diameter nitrogen cooling passage. The chevrons have a length of 4.13 cm, an included angle of 80 degrees and are stacked at 2.5 cm spacings. These extrusions can be obtained in lengths up to 5 meters. The use of extruded sections in contrast to the PLT type design results in an order of magnitude reduction in the number of aluminum welds.

Each helium panel is rigidly attached to the helium dewar by the two outer helium return lines (see Fig. 4). The other two return lines, as well as the liquid supply line, contain bellows joints to accommodate thermal contraction and alignment tolerances. All the helium supply and return lines are fabricated of Type 304 stainless steel. The transition from steel to aluminum is made at the panel supply manifolds through use of a Bi -Braze^R connector. To facilitate assembly and maintenance, the chevron panel manifolds are flange jointed to the nitrogen supply and return lines. Again, Bi -Braze^R connectors are used to effect the aluminum/steel transition, and bellows joints are utilized to allow for thermal contraction and alignment tolerances.

Cryopump instrumentation includes a helium dewar liquid level sensor and chevron and condensation panel temperature sensors. The liquid helium level is determined through use of an American Magnetics helium level sensor. Measurement of the chevron shield temperature is made at six locations using chromel-alumel thermocouples. The liquid helium cooled condensation panel temperature is measured at two locations using a calibrated carbon resistor and bridge network (see resistive sensor sketch in Fig. 6). The carbon resistor, which exhibits a large resistance increase between room temperature and 4°K is epoxy bonded in an aluminum block. The aluminum block can be screw attached to the condensation panel. This arrangement facilitates calibration since the sensor can be submerged easily in an open dewar of liquid nitrogen and/or liquid helium. The resistors used in the LITE cryopanel have $57\ \Omega$ at 300°K (room temperature), $70\ \Omega$ at 77°K and $193\ \Omega$ at 4.2°K .

The geometric factors controlling the molecular flow conductance through the chevron panel were evaluated to determine pumping speed. The experimental measurements of transmission probability (Refs. 2 and 3) formed the basis of this evaluation. The effect of chevron included angle on the experimentally determined transmission probability is shown in Part A of Fig. 7. The transmission probability increases with increasing angle to an angle of 90° above which it remains essentially unchanged with increasing angle. The transmission probability characteristic shown is for values of chevron length to spacing ratio greater than 10. Part B of the figure shows the effect of chevron included angle on the open area through the chevron bank. This shows the flow area to decrease in an approximately linear fashion with increasing angle.

The conductance through the chevron shield panel is proportional to the product of the transmission probability and the open area through the chevrons along with other factors not related to the geometry. An analysis of the trade-off between chevron angle and open area was made for the UTRC extruded chevron configuration to determine the optimum value of the product of the transmission probability and open area. The result of the optimization study is depicted by Part C of the figure which shows that the optimum occurs at approximately 80°. The value of the effective area for the UTRC design is calculated to be 308 sq cm.

The theoretical pumping speed per unit area of a condensation panel in $\ell/\text{sec-cm}^2$ is $S = (TP) \times \frac{\bar{v}}{4} \times \frac{1}{1000}$ where $\bar{v}$ is the mean molecular velocity and TP = transmission probability. For hydrogen at 300°K, $\bar{v} = 2.21 \times 10^5$ cm/sec, and for 80° chevrons, TP = 0.21, then $S = 11.6 \ell/\text{sec-cm}^2$. The open area of the chevrons (one side) is 1470 cm^2 ; thus the total theoretical pumping speed of the cryopump (including all four surfaces) is $6.8 \times 10^4 \ell/\text{sec}$. Experiments are currently in progress to measure the actual pumping speed of the LITE cryopump by flowing a known amount of hydrogen (or deuterium) into the neutralizer chamber and measuring the pressure (as shown in Ref. 4, a much lower ultimate pressure can be obtained with deuterium). These experiments will also permit the measurement of LHe use rate to determine how effective the LHe cooled panels are shielded from room temperature radiation and heat leaks.

REFERENCES

1. Haught, Alan F., et al.: High Beta Capture and Mirror Confinement of Laser Produced Plasmas. Semiannual Report, Contract EY-76-C-02-2277.*000, UTRC R77-953500-27, September 1977
2. Optimization of Molecular Flow Conductance, Leveson, Milleron and Davis, Lawrence Livermore Laboratory, Transactions of the 7th National Vacuum Symposium, 1960.
3. Journal of Applied Physics, Vol. 35, p. 529, 1964.
4. Pittenger, L.C.: Vacuum Engineering for Fusion Research and Fusion Reactors. Lawrence Livermore Laboratory, UCRL 78501. Presented at Annual Symposium of American Vacuum Society, Chicago, Illinois, September 1976.

SECTION II

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FACILITY FOR 10A INJECTION ONTO A TITANIUM WASHER GUN PLASMA TARGET

FIG. 1

LITE CRYOCONDENSATION PUMP

LN₂ COOLED CHEVRON ASSEMBLY

LHe COOLED PANELS ASSEMBLY

UTRC CRYOPANEL

TEMPERATURE SENSOR

CHEVRON PANEL CHARACTERISTICS

SECTION III

AMBIPOLAR POTENTIAL MODIFICATION BY CYCLOTRON RESONANCE

HEATED ELECTRONS IN MIRROR-CONFINED PLASMAS

Walter J. Fader

AMBIPOLAR POTENTIAL MODIFICATION BY CYCLOTRON RESONANCE
HEATED ELECTRONS IN MIRROR-CONFINED PLASMAS

W. J. Fader

It is well known that mirror confined plasmas are susceptible to loss cone instabilities that shorten plasma confinement time (Ref. 1). The growth of such instabilities can be suppressed by injection of cold ions into the loss cones as in the 2XII experiments (Ref. 2) and in the PR-6 experiment (Ref. 3). Alternatively, Ioffe and his collaborators (Ref. 3) have observed the quenching of instability in the PR-6 experiment as the result of cyclotron resonance heating of electrons. In their interpretation of this result, the resonance-heated electrons become mirror confined near the center of the plasma and create a local depression of the ambipolar potential in which cold ions become trapped and stabilize the loss cone modes.

In this paper are presented results of calculations of the sizes and shapes of ambipolar potential modifications produced by a group of cyclotron resonance heated electrons and the spatial distributions and densities of the ions trapped in the potential wells. A study of the effectiveness of the cold ions in suppressing the instabilities has been carried out by Jong (Ref. 4).

I. ELECTRON CYCLOTRON RESONANT HEATING

Because of the spatial variation of the magnetic mirror field confining the plasma, the applied microwave field is in resonance with the electron cyclotron frequency only at points of intersection of the field lines with a surface of constant $|B|$. If, in passing through a resonance point, the perpendicular component of the electron velocity is in phase with the perpendicular component of the local electric field vector of the microwaves, the electron will acquire an increment of perpendicular energy. In successive bounces between turning points, the electron passes through the resonance points with randomized phase, and the change in perpendicular energy can be characterized as a random walk in perpendicular velocity $v_{\perp}$. Because changes in parallel energy are of higher order, the perpendicular velocity $v_{\perp}$ becomes much larger than its parallel velocity $v_{\parallel}$, and the turning points of the cyclotron resonance heated electrons tend to bunch near the resonance points.

Eldridge (Ref. 5) has calculated this resonant diffusion in velocity space from the electron orbital motion, treating the interaction with the microwave field as a small perturbation. In doing so, he derived a perpendicular velocity diffusion coefficient $B_{\perp}$ in the form

$$B_{\perp} = 2\pi \left(\frac{e}{m}\right)^2 |E|_R^2 \left\{ L \left(\frac{d\Omega}{ds}\right)_R \right\}^{-1} \quad (1)$$

where $|E|_R^2$ is the mean-squared value of the microwave field at a resonance point, L is the separation of the two resonance points on a field line, and

$(d\omega/ds)_R$ is a derivative of the cyclotron frequency in the direction of the local magnetic field at the resonance point. The time evolution of the distribution function of resonant heated electrons can be calculated for collisionless heating as the solution of the diffusion equation

$$\frac{\partial f}{\partial t} = \frac{1}{v_{\perp}} \frac{\partial}{\partial v_{\perp}} v_{\perp} B_{\perp} \frac{\partial f}{\partial v_{\perp}} \quad (2)$$

Starting with Maxwellian electrons at temperature kT_e , their distribution evolves into the bimodal form

$$f(v_{\perp}, v_{\parallel}, t) = C(t) \exp(-\beta_{\parallel} v_{\parallel}^2) \exp(-\beta_{\perp}(t) v_{\perp}^2) \quad (3)$$

where

$$\beta_{\parallel} = m/2kT_e$$

and

$$\beta_{\perp} = \beta_{\parallel} / (1 + 4\beta_{\parallel} B_{\perp} t) \quad (4)$$

Hence, at the resonance points, the distribution of resonance-heated electrons remains Maxwellian in $v_{\parallel}$ with the temperature of the bulk of electrons, and Maxwellian in $v_{\perp}$ with a temperature linearly increasing with time. For $4\beta_{\parallel} B_{\perp} t \gg 1$, the distribution is characterized by a perpendicular temperature or energy given by

$$kT_{\perp} = E_{\perp} = 2mB_{\perp} t = m/2\beta_{\perp}(t) \quad (5)$$

The cyclotron-heated electrons become mirror-confined as they gain perpendicular energy, and the increased concentration of hot electrons near the

resonance points perturbs the ambipolar potential distribution, depressing it to form a potential well for cold ions. The cold ions trapped in this well have a stabilizing influence on the mirror confined plasma which can be studied by choosing model distributions of the densities of the hot and cold ions along a field line and then solving the dispersion relation for the growth constants for unstable modes, as was done by Jong (Ref. 4). The spatial distribution of the cold ions must be consistent with the width, depth, and location of the potential well for the combined population of hot electrons and cold ions. A self-consistent calculation of particle density and ambipolar potential distributions is described in the next section.

II. PLASMA POTENTIAL AND PARTICLE DENSITIES

To find the distributions of ambipolar potential and particle densities along a field line, the density of each species is calculated with the local mirror ratio $\zeta = B(s)/B(0)$ and the ambipolar potential $\phi(s)$ as parameters. The quasineutrality condition provides an equation which may then be solved for $\phi(s)$ as a function of $\zeta(s)$:

$$s_+ n_+(\zeta, \phi) + s_c n_c(\zeta, \phi) = n_-(\zeta, \phi) + n_H(\zeta, \phi) \quad (6)$$

On the left hand side of Eq. (6), n_+ and n_c are the densities of the energetic mirror confined ions and the cold ions trapped in the potential well. On the right of Eq. (6), n_- and n_H are the densities of the bulk

electrons and cyclotron resonance heated electrons. The density $n_i(s)$ of the i th plasma species is given by the integral

$$n_i(s) = 2\pi \int_0^\infty v_\perp dv_\perp \int_{-\infty}^\infty f_i(v_\perp, v_\parallel, s) dv_\parallel \quad (7)$$

over velocity space. The distribution $f_i(v_\perp, v_\parallel, s)$ at any value of s on the field line is obtained from the distribution $f_{i0}(v_{\perp 0}, v_{\parallel 0})$ at some reference point on the field line provided that the guiding centers of all the particles passing through s also pass through the reference point. The integral (Eq.(7)) is easily transformed to an integral over $v_{\perp 0}, v_{\parallel 0}$, namely

$$\int f_i(v_\perp, v_\parallel, s) v_\perp dv_\perp dv_\parallel = \frac{\zeta(s)}{\zeta_0} \int \frac{v_{\parallel 0}}{v_{\parallel 0}} f_i(v_{\perp 0}, v_{\parallel 0}) v_{\perp 0} dv_{\perp 0} dv_{\parallel 0}, \quad (8)$$

where ζ_0 is the local mirror ratio at the reference point. In obtaining Eq. (8), use has been made of relations between the variables $v_\perp, v_\parallel$ and $v_{\perp 0}, v_{\parallel 0}$ provided by energy conservation,

$$v_\perp^2 + v_\parallel^2 + \frac{2Z_i e \phi(z)}{m_i} = v_{\perp 0}^2 + v_{\parallel 0}^2 + \frac{2Z_i e \phi_0}{m_i} \quad (9)$$

and the adiabatic invariance of $v_\perp^2/B$,

$$v_\perp^2 = v_{\perp 0}^2 \zeta / \zeta_0. \quad (10)$$

These relations are used again in eliminating $v_\parallel$ on the right of Eq. (8) to obtain

$$n_i(\zeta, \phi) = \frac{4\pi\zeta(s)}{\zeta_0} \int_{v_{\perp 0}^-}^\infty dv_{\perp 0} v_{\perp 0} \int_{v_{\parallel 0}^-}^{v_{\parallel 0}^+} \frac{v_{\parallel 0} dv_{\parallel 0} f_{i0}(v_{\perp 0}, v_{\parallel 0})}{[v_{\parallel 0}^2 + v_{\perp 0}^2(1 - \zeta/\zeta_0) - 2Z_i e(\phi - \phi_0)/m_i]^{1/2}} \quad (11)$$

The limits of integration in Eq. (11) must be chosen consistent with the locations of the boundaries of the region of velocity space occupied by the i th species of particles.

1. Energetic Ions

The model distribution chosen for energetic ions is a loss cone distribution containing a degree of anisotropy characterized by the parameter $\gamma = E_{\perp}/E_{\parallel} = \alpha_{\parallel}/\alpha_{\perp}$. The center of the mirror-confined plasma is chosen as the reference point where the energetic ion distribution function is assumed to be

$$f_0(v_{\perp 0}, v_{\parallel 0}) = A_+ (v_{\perp 0}^2 (R-1) - v_{\parallel 0}^2) \exp(-\alpha_{\perp} v_{\perp 0}^2 - \alpha_{\parallel} v_{\parallel 0}^2) \quad (12)$$

Since the kinetic energies of these ions are much larger than their potential energies, $2 \alpha_{\perp} Z_i e \phi / m_i$ and $2 \alpha_{\parallel} Z_i e \phi / m_i$ are small quantities and the potential terms in the integrand and integration limits of Eq. (11) can be neglected. The integration limits in Eq. (11) are set by the minimum $v_{\parallel 0}$ of ions accessible to the point ζ , $v_{\parallel 0}^- = v_{\perp 0} (\zeta-1)^{1/2}$, and by the loss cone boundary, $v_{\parallel 0}^+ = v_{\perp 0} (R-1)^{1/2}$. The lower limit $v_{\perp 0}^-$ is set equal to zero. The axial distribution of the density of energetic ions is, in this approximation,

$$\begin{aligned} n_+(\zeta) &= \pi^{3/2} A_+ \left(\frac{\partial}{\partial \alpha_{\parallel}} - (R-1) \frac{\partial}{\partial \alpha_{\perp}} \right) \frac{\zeta (R-\zeta)^{1/2}}{(\alpha_{\perp} + \alpha_{\parallel} (\zeta-1)) (\alpha_{\perp} + \alpha_{\parallel} (R-1))^{1/2}} \quad (13) \\ &= \frac{n_+(1) \zeta}{(1+\gamma(\zeta-1))^2} \left(\frac{R-\zeta}{R-1} \right)^{3/2}. \end{aligned}$$

The ϕ -dependent part of the $n_+(\zeta)$, which has been neglected, gives a small correction to Eq. (13) in the vicinity of the mirror points ($\zeta = R$), where $n_+(\zeta)$ is itself small. Near the midplane, $n_+(\zeta)$ is virtually independent of ϕ and perturbations of ϕ caused by the hot electrons.

2. Bulk Electrons

The bulk electrons provide the neutralizing background charge for the energetic mirror-confined ions and are essentially electrostatically confined by the plasma ambipolar potential. The electron distribution is well-approximated by Maxwellian with a temperature T_e of about one-third to one-fifth of $|e\phi_m|$, where $|\phi_m|$ is the magnitude of the potential at the mirror points relative to the midplane. The electron escape rate, which must equal the ion loss rate for quasineutrality, is just the rate at which electrons in the Maxwellian tail are scattered above the energy $|e\phi_m|$. The model distribution function for the bulk electrons is therefore chosen as Maxwellian and isotropic:

$$f_0(v_{\perp 0}, v_{\parallel 0}) = B_- \exp(-\beta(v_{\perp 0}^2 + v_{\parallel 0}^2))$$

$$\beta = m_e / 2kT_e \quad (14)$$

In Eq. (11), we set $Z = -1$ and, for convenience, chose $\phi_0 = 0$ at the midplane.

Then ϕ becomes negative, dropping to a value $-|\phi_m|$ at the mirror points.

Since the electron loss boundary intersects the $v_{\parallel 0}$ axis at $v_{\parallel 0} =$

$(2e|\phi_m|/kT_e)^{1/2}$, the range of $v_{\perp 0}$ integration is 0 to ∞ , the upper limit $v_{\parallel 0}^+$ is

$$v_{||0}^+ = (v_{\perp 0}^2 (R-1) - \frac{2e\phi_M}{m_e})^{1/2}, \quad (15)$$

while $v_{||0}^-$ is that value for which $v_{||} = 0$ at ζ, ϕ , namely, $v_{||0}^- = (v_{\perp 0}^2 (\zeta-1) - 2e\phi/m_e)^{1/2}$. Upon substituting $f_0(v_{\perp 0}, v_{||0})$ into Eq. (11) and integrating parts, the expression for the density $n_-(\zeta, \phi)$ is derived in agreement with the published result of Yushmanov (Ref. 6). In terms of the variables $\Phi = \beta 2e\phi/m_e = e\phi/kT_e$ and $\phi_m = e\phi_m/kT_e$, this density may be written as

$$\begin{aligned} n_-(\zeta, \Phi) = n_-(0) \exp(\Phi) & \left\{ \text{Erf}((\Phi - \Phi_m)^{1/2}) + \left(1 - \frac{\zeta}{R}\right)^{1/2} \exp\left(\frac{\zeta(\Phi - \Phi_m)}{R - \zeta}\right) \right. \\ & \times \left. \text{Erfc}((R/(R - \zeta))^{1/2} (\Phi - \Phi_m)^{1/2}) \right\} \left\{ \text{Erf}(-\Phi_m)^{1/2} + \left(1 - \frac{1}{R}\right)^{1/2} \exp\left(-\frac{\Phi_m}{R-1}\right) \right. \\ & \times \left. \left. \text{Erfc}(-\Phi_m R/(R - \zeta))^{1/2} \right\}^{-1} \end{aligned} \quad (16)$$

Now, $|\Phi_m|$ is on the order of 3 to 5 and, as will be seen, the range of $|\Phi|$ of interest near the midplane is 0.5 to 1, so that negligible error is involved in approximating the density of bulk electrons, $n_-(\zeta, \phi)$, by

$$n_-(\Phi) = n_-(0) \exp(\Phi). \quad (17)$$

Hence the bulk electron density is nearly independent of local mirror ratio, as expected for electrostatically confined electrons.

3. Cyclotron Resonance Heated Electrons

In constructing a model distribution function for the cyclotron resonance heated electrons, it was assumed that the heating is collisionless. That is, the heated electrons gain perpendicular energy faster than their parallel energies can change by collisional randomization of the velocities.

Since their parallel energies continue to be of the order of kT_e while their perpendicular energies become much larger than $|e\phi_m|$, they are magnetically rather than electrostatically confined. The perpendicular energy is acquired at two points on the field line where the electron cyclotron frequency is resonant with the applied microwave frequency. At the resonance points ($|s| = s_1$), the hot electron distribution is assumed to be of the form of Eq. (3). Then the reference function from which the density of resonant-heated electrons is calculated at other points is

$$f_1(v_{\perp 1}, v_{\parallel 1}) = C - \exp(-\beta v_{\parallel 1}^2) \exp(-\beta_{\perp} v_{\perp 1}^2). \quad (18)$$

In Eq. (11), the limits of integration are $v_{\perp}^{\pm} = 0$,

$$v_{\parallel 1}^{+2} = v_{\perp 1}^2 (R/\zeta_1 - 1) - \frac{2e}{m} (\phi_m - \phi), \quad (19)$$

and

$$v_{\parallel 1}^{-2} = \text{MAX} \left\{ v_{\perp 1}^2 (\zeta/\zeta_1 - 1) - \frac{2e}{m} (\phi - \phi_1); 0 \right\}. \quad (20)$$

Then

$$\begin{aligned} n_H(\zeta, \phi) = 4\pi C - \frac{\zeta}{\zeta_1} \int_0^{\infty} v_{\perp 1} dv_{\perp 1} \exp(-\beta_{\perp} v_{\perp 1}^2) \exp(\beta v_{\perp 1}^2 (1 - \frac{\zeta}{\zeta_1}) + \frac{2e\beta}{m} (\phi - \phi_1)) \\ \times \int_{u^-}^{u^+} \exp(-\beta u^2) du \end{aligned} \quad (21)$$

where

$$u^{\pm} = (v_{\perp 1}^2 (R - \zeta)/\zeta_1 + \frac{2e}{m} (\phi - \phi_m))^{1/2} \quad (22)$$

and

$$\bar{u} = 0 \quad \text{for} \quad \zeta \geq \zeta_1$$

$$= (v_{\perp 1}^2 (\zeta_1 - \zeta) / \zeta_1 + \frac{2e}{m} (\phi - \phi_m))^{1/2}, \quad \zeta < \zeta_1. \quad (23)$$

In approximating the upper limit u^+ by $u^+ = \infty$. The integrations give

$$n_H(\zeta, \Phi) = n_H(\zeta_1, \Phi_1) \frac{\zeta}{\zeta_1} \frac{\exp(\Phi - \Phi_1)}{1 + (\zeta/\zeta_1 - 1) E_{\perp} / kT_e} \equiv n_H(\zeta_1, \Phi_1) G(\zeta, \Phi) \quad (24)$$

in the case of $\zeta \geq \zeta_1$. For ζ smaller than ζ_1 , that is, at points between the center of the plasma and the resonance points, we make the approximations

$$u^+ = \infty \quad \text{and} \quad u^- \approx v_{\perp 1} (1 - \zeta/\zeta_1)^{1/2}, \quad (25)$$

and obtain the result

$$\begin{aligned} n_H(\zeta, \Phi) &= n_H(\zeta_1, \Phi_1) \frac{\zeta}{\zeta_1} \frac{\exp(\Phi - \Phi_1)}{1 + (1 - \zeta/\zeta_1)^{1/2} (E_{\perp} / kT_e)^{1/2}} \\ &\equiv n_H(\zeta_1, \Phi_1) G(\zeta, \Phi) \end{aligned} \quad (26)$$

for $\zeta < \zeta_1$.

4. Cold Ions

From inspection of Eqs. (24) and (26), it can be seen that the maximum in the density of resonant heated electrons occurs at a resonance point. This maximum produces a local minimum in the ambipolar potential at $\zeta = \zeta_1$, and there is, necessarily, a local maximum $\phi = \phi_2$ at $\zeta = \zeta_2$, between the resonance point and the mirror point when $\zeta = R$. The value ϕ_2 at the local maximum

may be larger or smaller than the value at the plasma center where $\zeta = 1$.

If larger, cold ions can be trapped electrostatically between the two points where $\phi = \phi_2$ on opposite sides of the plasma center. If ϕ_2 is smaller than the central value of ϕ , cold ions can still be trapped near the resonance points but excluded from the vicinity of $\zeta = 1$.

The cold ion velocity distribution function is specified at a resonance point as a Maxwellian with temperature T_c . Thus,

$$f_1(v_{\perp 1}, v_{\parallel 1}) = D_+ \exp(-\alpha_c (v_{\perp 1}^2 + v_{\parallel 1}^2)) , \quad \alpha_c = m_+ / 2kT_c . \quad (27)$$

In Eq. (11), the appropriate limits of integration are

$$\begin{aligned} v_{\perp 1}^- &= 0, \quad v_{\perp 1}^+ = \infty \\ v_{\parallel 1}^- &= (v_{\perp 1}^2 (\zeta / \zeta_1 - 1) + 2Z_c e(\phi - \phi_1) / m_+)^{1/2} , \end{aligned} \quad (28)$$

or zero, if the quantity in parenthesis is negative, while

$$v_{\parallel 1}^+ = (v_{\perp 1}^2 (\zeta_2 / \zeta_1 - 1) + 2Z_c e(\phi_2 - \phi_1) / m_+)^{1/2} . \quad (29)$$

Partial integration of Eq. (11) gives

$$\begin{aligned} n_c(\zeta, \phi) &= \pi^{3/2} D_+ \alpha_c^{-3/2} \exp(-2\alpha_c Z_c e(\phi - \phi_1) / m_+) \\ &\times \left\{ \text{Erf}(y) + \left(1 - \frac{\zeta}{\zeta_2}\right) \exp\left(\frac{\zeta y^2}{\zeta_2 - \zeta}\right) \text{Erfc}((\zeta_2 / (\zeta_2 - \zeta))^{1/2} y) \right\} \end{aligned} \quad (30)$$

in which

$$y = (2\alpha_c Z_c e(\phi_2 - \phi) / m_+)^{1/2} .$$

In terms of the ratio $A = T_e/T_c$ of bulk electron and cold ion temperatures and $A\Phi = 2Z_c e \phi / kT_c = a_c 2Z_c e \phi / m_+$, Eq. (30) may be rewritten

$$n_c(\zeta, \Phi) = n_c(\zeta_1, \Phi_1) \exp(-A(\Phi - \Phi_1)) J(\zeta, \Phi) / J(\zeta_1, \Phi_1) \quad (31)$$

in which

$$J(\zeta, \Phi) = \text{Erf}(A^{1/2}(\Phi_2 - \Phi)^{1/2}) + (1 - \frac{\zeta}{\zeta_2})^{1/2} \exp\left(\frac{\zeta}{\zeta_2 - \zeta} A(\Phi_2 - \Phi)\right) \times \text{Erfc}(A^{1/2}(\Phi_2 - \Phi)^{1/2} (\zeta_2 / (\zeta_2 - \zeta))^{1/2}) \quad (32)$$

for $\zeta \leq \zeta_2$. For values of Φ larger than Φ_2 , the cold ion density also vanishes. This equation is the analog of Eq. (16) for the bulk electrons, but the approximation that led to Eq. (17) is not valid here, because kT_c is comparable with the well-depth, $e(\phi_2 - \phi_1)$.

5. Ambipolar Potential

After substitution of $n_+(\zeta)$, $n_-(\phi)$, $n_H(\zeta, \phi)$ and $n_c(\zeta, \phi)$ from Eqs. (13), (17), (24) or (26), and (31) into the charge balance or quasineutrality equation (Eq. (6)), ϕ is defined as an implicit function of the local mirror ratio ζ . If both the energetic ions and cold ions are protons, $Z_+ = Z_c = 1$ and Eq. (6) may be rewritten as

$$n_+(1)F(\zeta) + n_c(\zeta_1, \Phi_1) \exp(-A(\Phi - \Phi_1)) J(\zeta, \Phi) / J(\zeta_1, \Phi_1) - n_-(0) \exp(\Phi) - n_H(\zeta_1, \Phi_1) G(\zeta, \Phi) = 0. \quad (33)$$

The function $F(\zeta)$ was defined in Eq. (13), while the two possible forms of the function $G(\zeta, \phi)$ are defined in Eqs. (24) and (26). The densities of cold ions

and hot electrons at $\zeta = \zeta_1$ can be specified as fractions Δ_+ and Δ_- of the ion density $n_+(\zeta_1)$ and electron density $n_-(\Phi)$, respectively. Thus,

$$n_c(\zeta_1, \Phi_1) = \Delta_+ n_+(\zeta_1) = \Delta_+ n_+(1) F(\zeta_1) \quad (34)$$

while

$$n_H(\zeta_1, \Phi_1) = \Delta_- n_-(\Phi_1) = \Delta_- n_-(0) \exp(\Phi_1). \quad (35)$$

After dividing through Eq. (33) by $n_+(\zeta_1)$ we have

$$\begin{aligned} \frac{F(\zeta)}{F(\zeta_1)} + \Delta_+ \exp(-A(\Phi - \Phi_1)) J(\zeta, \Phi) / J(\zeta_1, \Phi_1) \\ - \frac{n_-(0) \exp(\Phi)}{n_+(\zeta_1)} (\exp(\Phi - \Phi_1) + \Delta_- G(\zeta, \Phi)) = 0 \end{aligned} \quad (36)$$

and by setting $\zeta = \zeta_1$, $\Phi = \Phi_1$, we find that

$$\frac{n_-(0) \exp(\Phi_1)}{n_+(\zeta_1)} = \frac{1 + \Delta_+}{1 + \Delta_-}. \quad (37)$$

For any value of ζ in the range $1 \leq \zeta \leq .9R$, the scaled ambipolar potential ϕ can be calculated as the root of the function

$$\begin{aligned} H(\zeta, \Phi) \equiv \frac{F(\zeta)}{F(\zeta_1)} + \Delta_+ \exp(-A(\Phi - \Phi_1)) \frac{J(\zeta, \Phi)}{J(\zeta_1, \Phi_1)} \\ - \frac{1 + \Delta_+}{1 + \Delta_-} [\exp(\Phi - \Phi_1) + \Delta_- G(\zeta, \Phi)] = 0 \end{aligned} \quad (38)$$

The ambipolar potential wells created by mirror-confined, cyclotron resonance heated electrons in the absence of cold ions are easily calculated after

setting $\Delta_+ = 0$ in Eq. (38). First, Φ_1 is calculated using the function $G(\zeta, \Phi)$ from Eq. (26) and setting $\zeta = 1$, $\Phi = 0$ to obtain

$$\exp \Phi_1 = \left\{ 1 + \frac{\Delta_-}{\zeta_1 + (\zeta_1^2 - \zeta_1)^{1/2} (E_{\perp}/kT_e)^{1/2}} \right\} \frac{(1 + \gamma(\zeta_1 - 1))^2}{\zeta_1(1 + \Delta_-)} \left(\frac{R-1}{R-\zeta_1} \right)^{3/2} \quad (39)$$

Thus, in the range $1 \leq \zeta < \zeta_1$, we have

$$\begin{aligned} \Phi = \Phi_1 + \ln \zeta / \zeta_1 + 2 \ln \frac{1 + \gamma(\zeta_1 - 1)}{1 + \gamma(\zeta - 1)} + \frac{3}{2} \ln \frac{R - \zeta}{R - \zeta_1} \\ + \ln(1 + \Delta_-) - \ln \left(1 + \frac{\zeta}{\zeta_1} \frac{\Delta_-}{1 + (1 - \zeta/\zeta_1)^{1/2} (E_{\perp}/kT_e)^{1/2}} \right), \end{aligned} \quad (40)$$

and for $\zeta \geq \zeta_1$,

$$\begin{aligned} \Phi = \Phi_1 + \ln \zeta / \zeta_1 + 2 \ln \frac{1 + \gamma(\zeta_1 - 1)}{1 + \gamma(\zeta - 1)} + \frac{3}{2} \ln \frac{R - \zeta}{R - \zeta_1} \\ + \ln(1 + \Delta_-) - \ln \left(1 + \frac{\zeta}{\zeta_1} \frac{\Delta_-}{1 + (\zeta/\zeta_1 - 1) E_{\perp}/kT_e} \right). \end{aligned} \quad (41)$$

In the case of $\Delta_+ \neq 0$, the ambipolar potential may be calculated numerically by application of an iterative procedure such as Newton-Raphson iteration to Eq. (38) to find the roots Φ for specified values of ζ . The function $J(\zeta, \Phi)$ defined in Eq. (32) contains an additional parameter Φ_2 . The value of the local maximum between $\zeta = \zeta_1$ and $\zeta = R$. If Φ'_1 and Φ'_2 are the corresponding values calculated for the same value of Δ_- but for $\Delta_+ = 0$, then Φ_2 and Φ_1 are related by

$$\Phi_2 = \Phi_1 + (\Phi'_2 - \Phi'_1) \ln(1 + \Delta_+) \quad (42)$$

The parameter Φ , is then calculated from an analog of Eq. (39).

III. RESULTS

Before the variations of quantities with distance, s , along a field line can be displayed, the local mirror ratio $\zeta(s) = B(s)/B(0)$ must be specified. Figure 1 shows a plot of $\zeta(s)$ on the axis of the minimum $-B$ magnetic field of the LITE experiment (Ref. 7). It is typical and, in what follows, this dependence of $\zeta(s)$ has been assumed. The axial mirror ratio is approximately $R = 2$, and the axial distance between mirror points is approximately 32 cm.

The axial distributions of the densities of energetic, mirror-confined ions, corresponding to Eq. (13), scaled by the value at $s = 0$ are shown in Fig. 2 for two values of γ , the ratio of the perpendicular to parallel temperatures of the ions. The curve labeled $\gamma = 1$ represents the flattest distribution attainable for an axial mirror ratio, $R = 2$. As the anisotropy parameter increases to $\gamma = 3$, the turning points of the bounce motion of the ions move closer to $s = 0$, bunching the ion density distribution closer to the plasma center. It seems clear that, if the cold ions introduced into the plasma for stabilizing the loss-cone modes could also be bunched near the center by electrostatic trapping, their effectiveness could be enhanced. The distribution of mirror confined ions depends on γ and the local mirror ratio $\zeta(s)$ and is essentially independent of the ambipolar potential or its perturbations. In the absence of cyclotron resonance heated electrons and cold ions charge balance requires that the bulk electron density equal the ion density

everywhere, so that the axial distributions of bulk electron density are also shown in Fig. 2. The corresponding axial distribution of the ambipolar potential may, in turn, be obtained from Eq. (17) by equating n_- to n_+ , having chosen $\phi = 0$ at $s = 0$. In the approximations leading to Eqs. (13) and (17), the ambipolar potential varies as the logarithm of the plasma density. Plots of ϕ versus s are shown for $\gamma = 1$ and 3 in Fig. 3 for the first five centimeters from the plasma center. In this limited range of s , the plasma densities exceed half their central values and cold ions should be most effective in their stabilizing function.

The effects on these potential distributions of a group of mirror-confined cyclotron-resonance heated electrons are illustrated in Fig. 4 in the case in which the two resonance points coincide at $s = 0$. The ratio of the hot electron density to the bulk electron density at $s = 0$ is $\Delta_- = 0.5$. The shape of the axial distribution of the hot electron density is given by Eq. (24) and is represented by the bell-shaped dashed curve in the lower half of the figure. It drops to half its central value at $S = 1.1$ cm where $\zeta(s) = 1.01$, consistent with the ratio $E_{\perp}/kT_e = 100$ assumed for the hot electrons. As a result of the limited axial extent of hot electron distribution, they constitute a total fraction of all electrons of approximately 8 percent. As described in Section II, above, $\phi(0)$ was chosen to be zero for convenience in calculation. For the purpose of comparing the results with the unperturbed $\phi(s)$ in Fig. 3, curves for the same value of γ are matched at $s = 5$ cm. The dashed curves extending to $s = 0$ represent the results from Fig. 3. At $s = 0$

the perturbed potential is depressed by an amount $\Delta\phi = /n(1 + \Delta_-) = 0.4$ independent of the value of γ for the mirror-confined ions. On the other hand, the well depth, which is the difference between the value of ϕ at $s = 0$ and at the maximum is considerably diminished with increasing γ . The greater the anisotropy of the mirror-confined ions, the shallower the potential well created by a group of cyclotron resonant heated electrons, simply because the unperturbed ϕ falls off more steeply with increasing s .

The effects of changing the ratio of perpendicular energy of the hot electrons to their parallel temperature are illustrated in Fig. 5 for $\Delta_- = 0.5$ and $\gamma = 1$. The three bell-shaped curves in the lower half of the figure representing the axial distributions of density of the hot electrons correlate with $E_{\perp}/kT_e$, as expected with increased concentration of turning points close to the center with increasing $v_{\perp}/v_{\parallel}$. As $s = 0$, $\Delta\phi$ is the same for all three ratios of $E_{\perp}/kT_e$, but narrower hot electron distribution leads to a narrower but deeper potential well when the $E_{\perp}/kT_e$ ratio is increased. A ratio of 50 produces a broad, shallow well which is almost a flat plateau. From Eqs. (4) and (5), it should be possible to estimate the time required for the hot electron $E_{\perp}/kT_e$ ratio to reach each of the values 50, 100 and 1000 as $t = kT_e / emB_{\perp}(E_{\perp}/kT_e - 1)$, in which $B_{\perp}$ is the perpendicular velocity diffusion coefficient of Eq. (1). A problem arises in that the product $L(d\Omega/ds)_R$ in the denominator of Eq. (1) was meant to represent the ratio of the separation distance of the resonance points to the coherent length (times the gyro-frequency). In the present case, in which the resonance points are closer than a wavelength, the ratio of the separation distance to the coherence

length is approximated by unity and $L(d\bar{\Omega}/ds)_R$ is set equal to the gyrofrequency at the plasma center. For a bulk electron temperature of 10 eV and an rms microwave electric field strength of 50 V/cm, it is found that $t = 1.6 \mu\text{sec}$ for $E_{\perp}/kT_e = 50$, $t = 3.2 \mu\text{sec}$ for $E_{\perp}/kT_e = 100$ and $32 \mu\text{sec}$ for $E_{\perp}/kT_e = 1000$.

Cold ion trapping is illustrated in Fig. 6 for the case of $\Delta_- = 0.5$, $E_{\perp}/kT_e = 100$, $\gamma = 1$ and a cold ion temperature which is four times the bulk electron temperature. The curves at the bottom of the figure represent the axial distributions of cold ion density for cold ion fractions of $\Delta_+ = 0.1$ and 0.2 at $s = 0$. The potential well tends to "fill in" with the addition of cold ions, the change in Φ being on the order of the change in Δ_+ .

When the resonance points are moved away from $s = \pm .875$ cm, the hot electron density distribution still has peak values at the two resonance points as shown in Fig. 7. The interior solution of Eq. (26) for the hot electron density must be used for $\zeta < \zeta_1$. The two branches of the hot electron distribution are evident in the dashed curve at the bottom of Fig. 7. The ambipolar potential curves show corresponding cusp-like minima at the resonance point. The maximum in the $\gamma = 1$ curve near $s = 2$ cm exceeds the values of Φ at $S = 0$, so that cold ions trapped in this potential well may occupy the space between the maximum and the corresponding point on the opposite side of the plasma center. In the case of $\gamma = 3$, Φ everywhere to the right of $s = .875$ cm is less than its value at the plasma center. Hence, in this case, trapped cold ions would be excluded from the center of the plasma but would be localized in bands near the resonance points. These effects are

illustrated in Figs. 8 and 9. In Fig. 8, a fraction $\Delta_+ = 0.1$ of cold ions with a temperature $T_c = 4 T_e$ has been added at $s = .875$ cm to the $\gamma = 1$ potential well of Fig. 7. In the lower half of the figure, the resulting cold ion distribution is represented by the solid curve, while the dashed curve represents hot electrons. The cold ions range from $s = -2$ to $s = +2$ cm and their density distribution peaks at the resonance points. The dashed curve of $\Phi(s)$ in the upper half of the figure is the $\gamma = 1$ curve of Fig. 7 and comparison with the solid curve plotted for $\Delta_+ = 0.1$ shows that Φ is increased by approximately Δ_+ at the resonance point. In Fig. 9, the resonance points have been moved to $s = \pm 1.75$ cm and γ has been increased to 3. The effective width of the potential well has been reduced to a narrow band between 1.5 and 2.5 cm, and the cold ions are restricted to this range as shown by the solid curve at the bottom of the figure. The dashed potential curve shown at the top of the figure represents the potential before cold ions have been added. Again the change in Φ at the resonance point is approximately Δ_+ .

It appears that, as Ioffe and his collaborators (Ref. 3) have suggested, a group of cyclotron-resonance heated electrons can produce an electrostatic potential well of sufficient depth to trap an appreciable fraction of cold ions in a mirror-confined plasma. The results presented in the foregoing discussion are in agreement with the dependencies of the potential perturbation and well depth on hot electron fraction Δ_- found in Ref. 3. The present work

is based on a model velocity distribution of hot electrons derivable from a resonant diffusion model of cyclotron heating and, in addition, includes the effects of the trapped cold ions on the potential distribution by a self-consistent calculation. These refinements have not changed the picture of ambipolar potential modification presented in Ref. 3 in any fundamental way. However, the present results do show that the effectiveness of the hot electrons in trapping the cold ions and the spatial concentration of cold ions are both sensitive to the locations of the resonance points when they are moved away from the plasma center. The effect of spatial concentration of cold ions on their stabilizing effectiveness is the subject of a study made by Jong (Ref. 4).

REFERENCES

1. Post, R. F. and M. N. Rosenbluth: Phys. Fluids, Vol. 9, p. 730 (1966).
2. Coensgen, F. H., W. F. Cummins, C. Gormezano, B. G. Logan, A. W. Molvik, W. E. Nexsen, T. C. Simonen, B. W. Stallard, and W. C. Turner: Phys. Rev. Lett., Vol. 37, p. 143 (1976). Also: Logan, B. G., J. F. Caluser, F. H. Coensgen, D. L. Correll, W. F. Cummins, C. Gormezano, A. W. Molvik, W. E. Nexsen, T. C. Simonen, B. W. Stallard, and W. C. Turner: Phys. Rev. Lett., Vol. 37, p. 1468 (1976).
3. Ioffe, M. S., B. I. Kanaev, V. P. Pastukhov, and E. E. Yushmanov: Sov. Phys. JETP, Vol. 40, p. 1064 (1975).
4. Jong, R. A.: To be Published.
5. Eldridge, O.: Plasma Physics, Vol. 15, p. 676 (1972).
6. Yushmanov, E. E.: Sov. Phys. JETP, Vol. 22, p. 409 (1966).
7. Haught, A. F., et al.: High Beta Capture and Mirror Confinement of Laser Produced Plasmas. UTRC R77-953500-27, p. 28, September 1977.

SECTION III

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LOCAL MIRROR RATIO ON THE AXIS OF THE MINIMUM-B MAGNETIC FIELD OF LITE

AXIAL DISTRIBUTIONS OF THE DENSITY OF MIRROR-CONFINED IONS

AXIAL DISTRIBUTIONS OF AMBIPOLAR POTENTIAL

POTENTIAL WELLS PRODUCED BY A GROUP OF CYCLOTRON RESONANT HEATED ELECTRONS

$$E_{\perp}/kT_e = 100 \quad \Delta_- = 0.5$$

FIG. 4

DEPENDENCE OF THE POTENTIAL WELL SHAPE ON $E_{\perp}/kTe$

$\Delta_- = 0.5$

$\gamma = 1$

TRAPPING OF COLD IONS IN THE POTENTIAL WELL

$\Delta_- = 0.5, \gamma = 1$

$E_{\perp}/kT_e = 100$

$T_c/T_e = 4$

FIG. 6

POTENTIAL WELLS PRODUCED BY A GROUP OF CYCLOTRON RESONANT HEATED ELECTRONS

RESONANCE POINTS AT $S = \pm 0.875$ CM $\Delta_- = 0.5$, $E_{\perp}/kT_e = 100$

FIG. 7

TRAPPING OF COLD IONS IN THE POTENTIAL WELL

FIG. 8

TRAPPING OF COLD IONS IN THE POTENTIAL WELL

$\Delta_- = 0.5$

$E_{\perp}/kT_e = 100$

$s = \pm 1.75 \text{ CM}$

$\Delta_+ = 0.1$

$T_c/T_e = 4$

FIG. 9

SECTION IV

EFFECT OF LOCALIZED COLD PLASMA ON

DCLC STABILITY IN MIRROR PLASMAS

by

Raynard A. Jong

EFFECT OF LOCALIZED COLD PLASMA ON
DCLC STABILITY IN MIRROR PLASMAS

Raynard A. Jong

The suppression of the drift cyclotron loss cone (DCLC) instability in mirror plasmas has been demonstrated in a series of experiments¹⁻⁴ in which a cold plasma stream was injected through the mirrors into the central plasma region. While effective in decreasing the fluctuation levels produced by the DCLC mode, this method introduced a serious energy drain, and depressed the electron temperature.⁵ Using an alternative technique, electron cyclotron resonance heating (ECRH) to produce a small fraction of hot electrons, Ioffe has produced apparent stabilization for 50-100 μsec ² in the PR-6 mirror trap. Ioffe has proposed a model, subsequently analyzed by Aamodt⁶ and Kanaev⁷ in which the stabilization arises by cold ions which are trapped in local electrostatic wells formed when the small fraction of hot electrons are produced and modify the ambipolar potential. Cold ions are confined in the plasma for the time it takes for the ions to scatter up in energy and escape from the potential well. Hence the confinement time of the low energy particles is increased over its value in the absence of the local potential well, viz., the ion transit time. This stabilizing technique suggests that suppression of the DCLC mode can be achieved with a much lower cold ion throughput and consequently a smaller energy drain.

The stabilizing effect of cold ions trapped locally in ECR heating zones both on the plasma midplane and off the midplane has been studied numerically using the formulation of Pearlstein et al.^{8, 9} in which the dispersion relation for the DCLC mode is solved for finite sized plasmas containing hot and cold ion components. The electrons and cold ions are modeled by Maxwellian velocity distributions. The hot ions are modeled by a bimodal loss cone distribution which reduces to the collisional distribution of Berk, Pearlstein, and Cordey⁸ when the ion anisotropy index ($WGAM = \frac{E_{\perp}}{E_{\parallel}}$) is taken as unity.

The self-consistent spatial distributions of the electrons and ions using these model distributions have been calculated by Fader¹⁰ along with the predicted plasma potential profile as a function of the position of the ECR heating zone and the anisotropy of the electron and ion distributions. These calculated cold ion spatial distributions have been modeled for the calculations reported here, as a Gaussian peaked at the resonance position z_r

$$n = n_0 \exp \left\{ -(z - z_r)^2 / L^2 \right\} \quad (1)$$

As shown by the stability boundaries in Fig. 1 for cold ions peaked about the midplane, cold ions uniformly distributed throughout the plasma, with cold to hot ion temperature ratios (T_W/T_H) as indicated, would stabilize the DCLC mode with a cold ion density fraction consistent with that predicted by the criteria of Baldwin, Berk, and Pearlstein⁵ and of Berk and Gerver¹¹.

Narrow cold ion distributions ($L \approx .1 L_m$), where L_m is the scale length for magnetic fields modeled by

$$B(z) = B_0 \left(1 + z^2/L_m^2 - \frac{z^4/L_m^4}{4(R-1)} \right) \quad (2)$$

and R is the mirror ratio, were found not to stabilize the DCLC mode. Broader ($L \approx .4 L_m$) distributions were found to stabilize with peak cold ion fractions which were larger than the uniform density requirements.

The combination of a locally peaked cold ion distribution with a low level uniform distribution was considered. For plasmas typical of those produced by injection of 10 keV neutrals into a warm target plasma in the LITE mirror facility,¹² the ratio of cool to hot plasma in the uniform portion of the cold ion distribution, (n_u/n_H) , is plotted against the cool to hot ratio in the peak, (n_o/n_H) , in Fig. 2. For peaks which have scale lengths $L_S \geq .4$ significant decrease in the uniform component can be achieved with values of the peak ratio less than or comparable to the uniform ratio required in the absence of a peak. Similar behavior was found in calculations with the cool plasma parameters typical of Ioffe's experiment.² In Fig. 3 is shown a plot of the uniform ratio versus the peak ratio for Ioffe's parameters for the case $L_S = .4$. Large decreases in the required uniform cold plasma component are evident for small peak values of the cold plasma.

In Fig. 4, the stability boundary for the hot plasma is shown as a function of the resonance position, Z_R , for the scale length $L_S = .4$. The

plot indicates that trapping the cold ions closer to the midplane is more effective in suppressing the DCLC mode. The resonance positions shown are for the axial positions $Z_R = 0., .44$ and $.88$ cm.

The preceding plots all employed the collisional distribution corresponding to that used by Berk, Pearlstein and Cordey.⁸ Increasing the value of the ion anisotropy index ($WGAM = \frac{E_{\perp}}{E_{\parallel}}$) from unity to three produces the stability boundaries shown in Fig. 5. The uniform fraction of cold ions for a particular peak value is lower for $WGAM = 3$ than for the corresponding $WGAM = 1$ case. Since the hot ions are bunched more in the midplane when $WGAM$ is larger, it is expected that the peaked cold ion component would be more effective in suppressing the DCLC mode in this case, so that less uniformly distributed cold plasma would be necessary. While this suggests that highly anisotropic hot ion plasmas would be more easily stabilized by local trapping of cold ions, calculations by Fader¹⁰ indicate that it may be more difficult to product a potential well by ECR heating of electrons when the ions are highly anisotropic. Consequently, while the trapped cold ions would be more effective in suppressing the DCLC fluctuations, it remains to be demonstrated that the potential modifications necessary to trap these ions could be achieved.

These calculations indicate that the trapping of small amounts of cold ions near the midplane of a mirror confined plasma can reduce the amount of uniformly distributed cold plasma required to stabilize the plasma against the

DCLC instability. Consequently, less externally supplied cold plasma would be needed to maintain stability and the cold ion throughput and energy drag could be reduced from that found in present stream stabilized mirrors. Furthermore, cold ions trapped on the midplane are more effective in producing stability than cold ions trapped off of the midplane.

References

1. Y. V. Gott, M. S. Ioffe, B. I. Kanaev, A. G. Motlich, V. P. Pastukhov, and R. I. Sobolev: Nuclear Fusion Supplement, p. 55 (1975).
2. M. S. Ioffe, B. I. Kanaev, V. P. Pastukhov, and E. E. Yushmanov: Sov. Phys. - JETP, 40, 1064 (1975).
3. F. H. Coensgen, W. F. Cummins, C. Gormezano, B. G. Logan, A. W. Molvik, W. E. Nexsen, T. C. Simonen, B. W. Stallard, and W. C. Turner: Phys. Rev. Lett., 37, 143 (1976).
4. B. G. Logan, J. F. Clauser, F. H. Coensgen, D. L. Correll, W. F. Cummins, C. Gormezano, A. W. Molvik, W. E. Nexsen, T. C. Simonen, B. W. Stallard, and W. C. Turner: Phys. Rev. Lett., 37, 1468 (1976).
5. D. E. Baldwin, H. L. Berk, and L. D. Pearlstein: Phys. Rev. Lett., 36, 1051, (1976).
6. R. E. Aamodt: Science Applications, Inc. Report LAPS-6B, SAI-76-727-LJ, August 1976.
7. B. I. Kanaev, V. P. Pastukhov, and E. E. Yushmanov, JETP Letters, 18, 204 (1973).
8. H. L. Berk, L. D. Pearlstein, and J. G. Cordey: Phys. Fluids, 15, 891 (1972).
9. L. D. Pearlstein and N. Maron: Proceedings of the Annual Meeting on Theoretical Aspects of Controlled Thermonuclear Research, May 4-6, 1977, San Diego, California.
10. W. J. Fader: Bulletin of the American Physical Society, 22, 1093 (1977). (Paper 3G5).

References (Cont.)

11. H. L. Berk and M. J. Gerver: Phys. Fluids, 19, 1646 (1976).
12. A. F. Haught, W. B. Ard, W. J. Fader, R. A. Jong, A. E. Mensing, D. H. Polk, R. G. Tomlinson, and J. T. Woo: Plasma Physics and Controlled Nuclear Fusion Research. Vol. I, p. 391, IAEA, Vienna, (1974).

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PEAK DENSITY VS SCALE LENGTH

$$n = n_0 \exp - (Z_s/L_s)^2$$

$$T_W/T_H = \begin{cases} 0.003 & \text{---} \\ 0.03 & \text{---} \end{cases} \left. \vphantom{\begin{matrix} 0.003 \\ 0.03 \end{matrix}} \right\} E_i = 15 \text{ keV}$$

$$T_W/T_H = 0.1 \quad \text{---} \quad E_i = 0.3 \text{ keV}$$

HOT PLASMA STABILITY BOUNDARY

$$n = n_0 \exp(-Z^2/L^2)$$

$$T_i = 10 \text{ keV}$$

$$L_s = L/L_m$$

$$L_m = 11 \text{ cm}$$

$$\text{WGAM} = 1$$

FIG. 2

COOL PLASMA STABILITY BOUNDARY

$$n/n_0 = \exp(-Z^2/L^2)$$

$$L_s = L/L_m$$

$$T_i = 200 \text{ eV}$$

$$L_m = 30 \text{ cm}$$

DEPENDENCE OF HOT PLASMA STABILITY BOUNDARY ON COLD PLASMA LOCATION

$$T_i = 10 \text{ keV}$$

$$L_s = 0.4$$

FIG. 4

FIG. 5