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ANOMALOUS TRANSPORT IN MIRROR SYSTEMS

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ABSTRACT

As now being explored for fusion applications confinement systems based on the mirror principle embody two kinds of plasma regimes. These two regimes are: a) High-beta plasmas, stabilized against MHD and other low frequency plasma instabilities by magnetic-well fields, but characterized by non-maxwellian ion distributions. b) near-maxwellian plasmas, confined electrostatically (as in the tandem mirror) or in a field-reversed region within the mirror cell. Common to both situations are the questions of anomalous transport owing to high frequency instabilities in the non-maxwellian portions of the plasmas. This report will summarize the status of theory and of experimental data bearing on these questions, with particular reference to the high temperature regimes of interest for fusion power.

Key words: Anomalous transport; mirror machine; ion-cyclotron instabilities; loss-cone; plasma stabilization.

1) INTRODUCTION

Progress in mirror research toward the fusion goal has required the solution of two interrelated problems. The first of these problems concerned the need to develop technological methods for creating and sustaining fusion plasmas in a mirror system. The second problem has been to achieve the control of anomalous transport processes peculiar to mirror-confined plasmas.

The solution to the first problem has been found through the use of high intensity neutral beams. Through this technique it is possible to create, and then to sustain in a dc magnetic field, plasmas at the densities and temperatures needed for a fusion power system.

With respect to the second problem, while encouraging experimental and theoretical progress has been made, final proof that anomalous transport processes can be adequately suppressed in a full-sized mirror fusion power plant must await tests in scaled-up systems. For this reason the discussion of plasma

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of anomalous transport in large-scale mirror experiments is based on theoretical extrapolations from a data base of small-scale experiments. Nevertheless, there are some areas in which it is clear that the relevant problem has already been brought under control so that scale-up will not change the picture in an unfavorable way. An example is the use of the magnetic well principle to suppress MHD instabilities, especially high values of the plasma pressure parameter, β .

Nevertheless, however, it is necessary to distinguish different embodiments of mirror systems, among which the problems of anomalous transport differ.

Among these embodiments that are characterizable as "open" systems, there are three:

1. The conventional single cell mirror system utilizing a magnetic well field. Most mirror experiments have utilized this configuration.
2. The Tandem Mirror (Dimov, Zakaidakov, Kishinevsky, 1976; Fowler, Logan, 1977) in which two conventional magnetic well mirror cells are used as electrostatic plugs to contain a maxwellian plasma in a central solenoid.
3. The Field-Reversed Mirror, in which mirror fields are used to confine a field-reversed plasma entity.
4. The multiple mirror, in which many mirror cells are placed in series in order to increase the confinement time by inhibiting the rate of escape of plasma out the ends.

As mentioned earlier, owing to the use of magnetic well fields in the conventional mirror, MHD instabilities are suppressed and only high frequency instabilities are possibly present to cause anomalous transport. This remark also applies to the end plugs of the TM. However the TM central cell, containing as it does a near-maxwellian plasma, will not be subject to h.f. instabilities. Depending on the field configuration employed and on the plasma β MHD fluting or ballooning modes may be present. In the FRM, again the main issue is MHD stability, rather than high frequency instabilities. Finally, in the multiple mirror system the issues are probably two-fold: MHD stability of the plasma column (as dependent on the effectiveness of the, necessarily, average minimum-B fields that must be employed), and h.f. instabilities that might occur in the outer cells, where loss-cone-like distribution functions may exist.

In summary, therefore, we see that mirror systems may be subject to anomalous transport from two general causes: High Frequency instabilities driven by loss cone distributions in conventional mirror cells, or else MHD-like instabilities occurring in the central cell of a TM or in the plasma of a FRM, or in the plasma column of a multiple mirror system. In this report we will concentrate on describing the theoretical and experimental work that has dealt with the first of these problems, high frequency instabilities and their effects. Leaving aside early work with conventional mirrors, such as the experiments of Ioffe (Ioffe and Yushmanov, 1962), there exists relatively little experimental work on MHD effects in mirror systems of the FRM and multiple mirror types. Also there is as yet virtually no experimental information on the TM, the principles of which are only now beginning to be investigated. For these latter systems it is therefore necessary to depend mainly on theory in order to make predictions as to anomalous transport. Problems of transport in the central cell of a TM will be discussed at this conference by Ryutov.

1. HIGH FREQUENCY INSTABILITY IN MIRROR SYSTEMS - GENERAL CONSIDERATIONS

The driving source for high frequency instabilities in plasma confined in a conventional mirror cell is the non-maxwellian nature of the plasma that is a prerequisite for confinement. This requirement can be understood from the conditions that describe the conditions for pressure equilibrium along the magnetic field lines:

$$\hat{b} \cdot \nabla p = -(\rho_i + \rho_e) \frac{\hat{b} \cdot \mathbf{B}}{B}, \quad (1)$$

where $\hat{b}$ is a unit vector directed along $\mathbf{B}$ and ρ_i and ρ_e are respectively the perpendicular and parallel components of the plasma pressure tensor. If the plasma pressure is isotropic ($\rho_i = \rho_e = p$) there can be no gradient in pressure along $\mathbf{B}$, thus no longitudinal containment. Anisotropy of the total plasma pressure is thus a minimum requirement for mirror confinement. However the nature of the situation imposes further requirements on the particle distribution functions. Particle loss through the mirrors results from the scattering of particles in velocity space into the forbidden pitch angles or "loss-cone" region where the mirror effect is no longer strong enough to prevent escape. Since the scattering rate for electrons, owing to their higher velocities, is much more rapid than that for ions, their loss rate would be expected to be higher. In actual fact, however, this differential loss rate leads immediately and automatically to the buildup of a positive ambipolar plasma potential that brings the electron loss rate to equality with that for the ions. The magnitude of this potential, as determined by classical collisional transport rates, is typically 4 to 5 kT_e. Thus, except for a truncation at energies greater than the ambipolar potential, the electron distribution functions in mirror-confined plasmas are maxwellian. In $v_\perp, v_\parallel$ space (evaluated at the minimum point $B = B_{\min}$ between the mirrors) they therefore fill a region bounded from either side by two hyperboloids of revolution, as indicated schematically in Fig. 1. In addition to the large ambipolar potential drop between the center of the plasma and the mirrors there is a smaller potential drop ($\phi_m - \phi_w$) occurring in a sheath near the wall, set up in usual manner to maintain quasi-neutrality in the streaming plasma between the mirror and the wall. This potential and its effect on the electron confinement is also shown on the figure.

The ions, on the other hand, cannot be confined unless their kinetic energy exceeds a minimum value given by

$$W_{\min} = \frac{Ze\phi_m}{R_m - 1}, \quad (2)$$

where ϕ_m is the ambipolar potential and R_m is the mirror ratio (ratio of mirror field to $B_{\min}$). The loss cone boundaries for the ions can also be described in $v_\perp, v_\parallel$ space (as evaluated at $B_{\min}$) by a hyperboloid of revolution as shown in Fig. 2. However in this case no low energy particles or particles remain trapped whose $B_{\min}$ pitch angles $\theta = \tan^{-1} v_\perp/v_\parallel$ are smaller than the critical pitch angle, $\theta_c = \sin^{-1} [(1 + e\phi_m/kT_e)/R_m]^{1/2}$.

From the above general considerations and from theory some general conclusions can be drawn concerning the characteristics of plasma instabilities in plasmas confined in a magnetic well mirror cell. These are:

1. Low frequency instabilities ($\omega \approx \omega_{ci}$) are forbidden on energetic grounds by the magnetic well.
2. Being quasi-maxwellian, the electrons do not represent a source of free energy to drive instabilities at electron frequencies ($\omega \approx \omega_{ce} \approx \omega_{pe}$).
3. The above constraints limit the frequency spectrum of instabilities to ion characteristic frequencies, i.e. $\omega_{ci} \leq \omega \leq \omega_{pi}$.

In addition to the above general constraints, instabilities in mirror plasmas at temperatures and densities appropriate for fusion applications are subject to additional constraints that still further restrict the number, the type, and the effects of plasma instabilities. These limitations will be discussed below.

Fig. 1 Loss hyperbola for electrons (Symmetric with respect to vertical axis).

Fig. 2 Loss hyperbola for ions (symmetric with respect to vertical axis)

III) INSTABILITY MODES IN MIRROR-CONFINED FUSION PLASMAS

Early mirror experiments were often carried out at low plasma densities ($\omega_{pi} \approx \omega_{ci}$) and with poorly randomized ion distribution functions and under poor vacuum conditions. In these experiments a large number of unstable high frequency modes were detected, almost all of which have been shown to be suppressed under plasma conditions more representative of those required for fusion power. Among these early-found modes were the following:

- **Negative mass instabilities:** Suppressed by magnetic well fields, by broadening ion distribution functions and by increasing plasma density.
- **Dory-Guest-Harris mode** (Dory, Guest and Harris, 1965): Suppressed by broadening ion distribution functions.
- **Harris anisotropy mode** (Harris, 1961; Hall, Heckrotte and Kammash, 1964): Eliminated by broadening ion angular distributions and by increasing plasma density.
- **Double-hump instability** (Pearlstein, Rosenbluth and Chang, 1966): Avoided by eliminating double-peaking in ion energy distribution functions.

It was therefore found, as predicted by theory and as confirmed by experiment, that the effect of pushing the plasma conditions in the experiments closer to

fusion parameters was to eliminate or suppress many of the high frequency instability modes that plagued earlier mirror experiments. A partial explanation for this happy circumstance can be found by examining the approximate dispersion equation describing unstable waves in homogeneous mirror plasmas. This equation can be written in the form

$$D(\omega, k) = 0 = 1 + \left(\frac{\omega_{pe}^2}{\omega^2} \right) - \left(\frac{\omega_{pe}^2}{\omega^2} \right) \cdot \left(\frac{k_{\parallel}^2}{k^2} \right) - \left(\frac{\omega_{pi}^2}{\omega_{ci}^2} \right) G_i. \quad (3)$$

Here ω_{pe} , ω_{ce} and ω_{pi} , ω_{ci} are the electron and ion plasma and cyclotron frequencies respectively and G_i is the ion term, consisting of a sum over resonant terms involving the ion cyclotron frequency and its harmonics. These terms contain Bessel-weighted integrals of the ion distribution functions and their derivatives.

Since

$$1 \ll (\omega_{pi}^2/\omega_{ci}^2) \ll (\omega_{pe}^2/\omega_{ce}^2),$$

and since $\omega \sim \omega_{ci} \ll \omega_{ce}$ the electron terms dominate in determining $k_{\parallel}$, leading to the result that

$$k_{\parallel}^2/k^2 \approx \frac{\omega^2}{\omega_{ce}^2} \ll 1.$$

In other words the instabilities are necessarily highly flute-like, i.e. either have $k_{\parallel} = 0$ or at least have long wavelengths along the field lines compared to their perpendicular wavelength. In addition to this restriction, electron Landau damping effects can impose the further restriction that the parallel phase velocities of the unstable wave must exceed the maximum parallel velocity of the electrons, i.e. $\omega/k_{\parallel} > (e\phi/2m_e)^{1/2}$. There is also a third condition, namely that in order to couple to the ion motions the perpendicular wavelength of the unstable waves should be comparable to the ion gyroradii, i.e. $k_{\perp} a_i \gtrsim 1$. Still further restrictions, ones which can be exploited in minimizing the tendency to instability, result from the presence of the magnetic field gradients associated with the mirrors (detunes cyclotron resonances and limits the axial wavelengths of the modes). Finite plasma beta has also been shown to exert a stabilizing effect on some modes.

As a result of the above constraints not only is the number of unstable modes in mirror fusion plasmas sharply delimited but also the potentially unstable regime in ω, k space is restricted. When these constraints and restrictions are imposed there appear to remain only three unstable modes that must be dealt with specifically. These are:

- The Drift Cyclotron Loss Cone (DCLC) mode, an azimuthal electron drift wave arising in the presence of a density gradient and driven unstable by the positive slope region of the ion distribution function ($\partial f/\partial v_i > 0$).
- The Alfvén Ion Cyclotron (AIC) mode, a longitudinally propagating Alfvén wave driven unstable at the ion cyclotron frequency by anisotropy of the ion distribution functions.

- The High Frequency Convective Loss Cone (HFCLC) mode, a high frequency short perpendicular wavelength ($k_{\perp} a_i \gg 1$) longitudinally propagating electron plasma wave, driven unstable by the positive $\partial f / \partial v_i$ regions of the ion distribution function.

The main issue determining the extent of anomalous transport in a conventional mirror cell therefore comes down to the question of the presence or absence of the above three instability modes. To anticipate what is to come, thus far only the DCLC mode has been evident in mirror experiments carried out in the fusion regime. For this mode a stabilizing means (warm plasma stabilization) has been demonstrated. In theory the other two modes, the AIC and the HFCLC modes, are also subject to control, by distribution function controls such as warm plasma stabilization or by limiting the plasma length, these two techniques either being used separately or in combination. In fact, since these other modes are less robust than the DCLC mode their absence can possibly be explained by the fact that up till now the circumstances acting to suppress the DCLC mode have at the same time eliminated these two modes.

IV) THE NATURE OF ANOMALOUS TRANSPORT EFFECTS IN MIRROR-CONFINED PLASMAS

The foregoing considerations concerning the characteristics of the unstable modes in mirror-confined plasmas strongly influence the nature of the anomalous transport that can be expected to occur as a result of their presence. As noted previously, electron mobility and electron Landau damping enforce the condition $k_{\parallel} / k_{\perp} \ll 1$, and these constraints, together with the requirement for coupling to the ion motions impose the further constraint that $\omega_{ci} \leq \omega \leq \omega_{pi}$ and that $k_{\perp} a_i \geq 1$ (more typically $k_{\perp} a_i \gg 1$ under fusion plasma conditions).

It follows from the condition $k_{\parallel} \ll k_{\perp}$ and from the influence of electron Landau damping that under all circumstances the parallel component of the fluctuating electric field of the instabilities is small compared to the perpendicular component, i.e. $E_{\parallel} / E_{\perp} \ll 1$. It furthermore follows from this circumstance and from the frequency spectrum involved that under the influence of the wave fields the electrons will respond adiabatically - i.e. will oscillate transversely with the $E \times B$ drift velocity, while the ions will behave non-adiabatically, i.e. will be subject to a two-dimensional diffusion in velocity space, the primary effect of which can be expected to be a stochastic one, i.e. a diffusion in $v_{\perp}$ space toward broader distributions and higher perpendicular temperatures. The two-dimensional nature of this velocity-space diffusion process is to be contrasted with the 3-d velocity-space diffusion arising from coulomb collisions. In collisional diffusion the ions and the electrons separately diffuse toward isotropic maxwellian distribution functions, as modified and constrained by the mirror loss-cone velocity-space boundary conditions previously discussed. In addition, the equilibrium electron temperature always being lower than the ion temperatures ($T_e / T_i \leq 0.1$), the ions are subjected to a dynamical friction drag which tends to slow them down, with relatively little accompanying diffusion in pitch angle. This slowing down competes with the energy-broadening effects of collisions and the ion heating effect of neutral beam injection to establish a steady state ion distribution function in energy and pitch angle. The general nature of these diffusion processes in $v_{\perp}$, $v_{\parallel}$ space is shown schematically in Fig. 3.

If we now superpose on this picture the two-dimensional diffusion (diffusion in $v_{\perp}$ only) caused by ion-cyclotron instabilities such as the DCLC mode, the main effect will be to diffuse the ion energies toward higher and lower values, i.e. to enhance the high energy end of the distribution in competition with electron

drag, while at the same time causing ions to diffuse down to the loss boundary, i.e. toward the minimum energy for trapping, $Z\phi/(R-1)$.

- Ions, confined by mirrors, cool by heating electrons

Fig. 3 Schematic representation of particle diffusion in a mirror system.

The theory of both collisional and instability-induced diffusion processes in mirror systems has been reviewed by Baldwin (Baldwin, 1977). In the next three sections we will address the specific case of the DCLC mode, the anomalous transport that it causes and the effect of warm plasma stabilization.

V) THE DCLC MODE AND THE CONCEPT OF WARM PLASMA STABILIZATION

The DCLC mode was predicted theoretically in 1966 (Post and Rosenbluth, 1966). As mentioned previously it was expected to occur in all mirror-confined plasmas for which the plasma radius (actually the radial density gradient characteristic length) was smaller than a critical value. This critical value, as measured in dimensionless units (ratio of plasma radius R_p to mean ion gyroradius, a_i), was fairly large, $R_p/a_i \approx 100$. Thus for the relatively small ($R_p/a_i \approx 3$) plasmas typically under study experimentally the DCLC mode was expected to be strongly excited, possessing a relatively rapid growth rate ($\gamma \sim \omega_{ci}$). Under such conditions early theoretical estimates (Galeev, 1966) of the rate of anomalous transport that the DCLC mode was expected to cause gave the prediction that it would result in loss of plasma confinement in times of order the transit time of the ions between the mirrors, i.e. $\tau \sim L_p/\bar{v}_i$, plasma length divided by mean ion thermal velocity. In fact no such rapid losses were observed. In the plasmas then under study (for example PR-6 in Moscow or 2X in Livermore) which had ion temperatures of under 500 eV to 1 keV and densities of order several times 10^{13} cm^{-3} loss times were seen to be of order of the classical (collisional) mirror loss rates (~ 0.5 millisecond in 2X), i.e. nearly three orders of magnitude slower than the anomalous transport rate first predicted for the DCLC mode. Not only were greatly enhanced end losses not seen, but there was also no clear evidence of even the existence of the azimuthal plasma wave, with its phase velocity in the ion diamagnetic direction, that would have been the hallmark of the DCLC mode.

Apart from some early evidence, from PR-6 (Baiborodov, Gott and Ioffe, 1969; Baiborodov, and co-workers, 1969) and PR-7 (Gott and co-workers, 1975), of the possible existence of a DCLC-like mode the paradox of its apparent unimportance compared to the theoretical predictions was not fully resolved until high power

neutral beams were added to the Livermore 2X mirror experiment, by that time called 2XIIB. In this experiment it was possible to drive the plasma ion mean energies up to 13 keV, at which point it became very clear that the confinement parameter, $n\tau$, for the plasma, rather than scaling up as $E^{3/2}$, as predicted by classical theory, was seen to exhibit an approximately constant value of τ , still at a value of order 0.5 ms, roughly the same as earlier observed at much lower mean ion energies where near-classical loss rates prevailed. At the same time there was clear evidence of r.f. activity at or near the midplane ion cyclotron frequency, associated with an azimuthal wave of a phase velocity and azimuthal wavelength consistent with the DCLC mode. There still remained the mystery of why this instability caused only relatively mildly enhanced losses, as compared to the early theoretical estimates.

The final resolution of this mystery came only after the experimental demonstration of a means for controlling the DCLC mode, the "warm plasma stabilization" technique mentioned earlier. In this technique a low density stream of plasma is sent through one of the mirrors, penetrating into the confined plasma. As a result of stochastic diffusion in the instability fields ions from this plasma and ions diffused down from the contained distribution tend to fill in the low energy end of the ion distribution function, thereby driving the instability toward a point of marginal stability. Depending upon the value of R_p/a_i the fractional amount of warm plasma required to achieve stability can be relatively small; it is not necessary to fill the loss cone completely to suppress the mode.

Tried out on 2XIIB, warm plasma stabilization was found to be highly effective, reducing ion cyclotron fluctuations to low levels and essentially restoring the $E^{3/2}$ scaling of $n\tau$ predicted by classical theory (Coensgen and co-workers, 1975). But now, in studying the experimentally observed scaling laws of warm plasma stabilization and their theoretical interpretation it became clear that the puzzling absence of strong losses from the DCLC mode could be explained. If no warm plasma was supplied the mode would grow only to that amplitude needed to generate its own "warm plasma" by v_1 diffusion, i.e. it will always drive itself to the point of marginal stability. At plasma ion temperatures of order 1 keV and densities of 10^{13} or higher the fluctuation level needed to maintain marginal stability was too low to significantly increase the plasma loss rate over its classical value. But at 10 keV and comparable densities the classical loss rate dropped well below that set by the level of DCLC-induced fluctuations needed to provide the self-generated warm plasma. Under these circumstances adding warm plasma dramatically reduced the fluctuation levels and correspondingly increased the confinement time. Furthermore, from the scaling laws for the DCLC mode it would be predicted that for R_p/a_i values larger than the values of ~ 3 characteristic of 2XIIB smaller and smaller fractional densities of warm plasma would be required to stabilize the mode.

The theory of warm plasma stabilization, to be outlined in the next section, resulted in a completely satisfactory explanation of the 2XIIB results and in a basis for extrapolation to scaled-up experiments.

VI) OUTLINE OF THEORY OF WARM PLASMA STABILIZATION OF THE DCLC MODE

On the basis of the new understanding of the DCLC mode achieved in the 2XIIB and PR-6, PR-7 experiments Baldwin, Berk, and Pearlstein (1976) have developed a detailed theory of the processes involved, including both the theory of the marginal stability conditions for the DCLC mode (Berk and Gerver, 1976) and the treatment, by quasilinear theory, of the instability-induced diffusion processes involved (Berk and Stewart, 1977). We will here outline the essential elements

of this theory in the context of the issue of anomalous transport. Our starting point is the dispersion equation for azimuthally- propagating flute-like ($k_{\parallel} = 0$) waves in a plasma with a radial density gradient of characteristic length R_p . Compared to the approximate dispersion equation that we discussed previously, this dispersion equation sets $k_{\parallel} = 0$ and introduces new terms to take into account the density gradient and the effects of finite beta. The equation is

$$D(\omega, k) = 1 + \frac{\omega_{pe}^2}{\omega_{ce}^2} + \frac{\omega_{pi}^2}{\omega_{ci}^2} \left\{ \frac{\omega_{pi}^2}{k_{\perp}^2 c^2} - \frac{\omega_{ci}}{k_{\perp} R_p \omega} \left(1 + \frac{\beta}{2} \right) + G_i \right\}, \quad (4)$$

where the ion term, G_i , is

$$G_i = \sum_{n=-\infty}^{\infty} \frac{\omega_{ci}^2}{k_{\perp}^2} \int d^3v \left(\frac{\omega}{v_{\parallel}} \frac{\partial f}{\partial v_{\parallel}} - \frac{k_{\perp}}{\omega_{ci}} \left| \frac{\partial f}{\partial r} \right| \right) \cdot J_n^2 \left(\frac{k_{\perp} v_{\perp}}{\omega_{ci}} \right) \left[\frac{1}{\omega - n\omega_{ci}} \right]. \quad (5)$$

The above equations in the main are to be treated for the case of marginal stability of the lowest harmonic ($n = 1$) of the ion cyclotron frequency. Under this condition $\omega \approx \omega_{ci}$, all higher harmonics are stable, and simplifications result.

Given an ion distribution function $f(v_{\perp}, v_{\parallel})$, obtained, for example, from Fokker-Planck solutions, the ion term G_i may be evaluated and the wave stability properties determined. As a model which reproduces the essential features of the problem we shall adopt a loss-cone-like distribution function that consists of the product of a maxwellian distribution function in $v_{\parallel}$ and the difference of two maxwellians in $v_{\perp}$, thereby modeling the typical loss-cone "hole" in $v_{\perp}$ space. To this primary ion distribution function we will then add a fractional component comprised of a lower temperature maxwellian that will approximately model the warm plasma component. Since $k_{\parallel} = 0$ the $v_{\parallel}$ integration may be performed immediately, to yield distribution functions $F(v_{\perp})$. Thus we have

$$F(v_{\perp}) = \frac{1}{v_H^2 - v_h^2} \left[\exp \left(\frac{-v_{\perp}^2}{v_H^2} \right) - \exp \left(\frac{-v_{\perp}^2}{v_h^2} \right) \right] + \frac{\Delta}{v_w^2} \exp \left(\frac{-v_{\perp}^2}{v_w^2} \right). \quad (6)$$

Here v_H^2 is the mean squared velocity of the "hot" ion distribution, v_h^2 is that for the "hole", and v_w^2 is that for the warm plasma, whose fractional density $n_w/n_H = \Delta$.

In evaluating G_i integrals are repeatedly encountered which are of the form

$$A_n = \int_0^{\infty} dv_{\perp}^2 \left[\exp \left(\frac{-v_{\perp}^2}{v_0^2} \right) \right] J_n^2 \left(\frac{k_{\perp} v_{\perp}}{\omega_{ci}} \right), \quad (7)$$

where $v_0^2 = v_H^2, v_h^2, v_w^2$.

These may be evaluated by G.R. 6.633-2

$$\int_0^\infty e^{-\rho^2 x^2} J_p(\alpha x) J_p(\beta x) x dx = \frac{1}{2\rho^2} \exp\left(-\frac{\alpha^2 + \beta^2}{4\rho^2}\right) I_p\left(\frac{\alpha\beta}{2\rho^2}\right). \quad (8)$$

We will also utilize the following identity to evaluate sums over the ion cyclotron harmonics $n \neq +1$.

$$\sum_{n \neq 1} \frac{I_n(b)e^{-b}}{1-n} = \frac{1}{b} \left[1 - I_0(b)e^{-b} \right]. \quad (9)$$

Further simplifications result from recognizing that the term $\omega_{pi}^2/k_1^2 c^2$ in the dispersion equation reduces to $\beta/k_1^2 a_1^2$ and that $\beta \approx 1$ and $k_1^2 a_1^2 \gg 1$ for the cases of fusion interest. Here we have taken $a_1 \equiv \omega_{ci} v_H$

With the above reductions the dispersion equation takes the form

$$\begin{aligned} D(\omega, k) = & 1 + \frac{\omega_{pe}^2}{\omega_{ce}^2} + \frac{\omega_{pi}^2}{\omega_{ci}^2} \left\{ \left(\frac{\beta}{2} \right) \frac{1}{x^2} - \frac{1}{\sqrt{2}} \left(\frac{\omega_{ci}}{\omega} \right) \left(\frac{a_1}{R_p} \right) \left(1 + \frac{\beta}{2} \right) \frac{1}{x} \right. \\ & + \frac{1}{x^4} H_0(\alpha^2, x^2) + \frac{1}{4x^3} \left(\frac{a_1}{R_p} \right) G_0(\alpha^2, x^2) \\ & \left. + \frac{1}{x^2} G(\alpha^2, x^2) \left[\frac{\frac{\omega_{ci}}{\omega}}{\frac{\omega_{ci}}{\omega} - 1} \right] + \frac{1}{\sqrt{2}} \frac{1}{x} \left(\frac{a_1}{R_p} \right) H(\alpha^2, x^2) \left[\frac{1}{\frac{\omega_{ci}}{\omega} - 1} \right] \right\}. \quad (10) \end{aligned}$$

Here $x^2 = (k_1^2 a_1^2)/2$, $\alpha^2 = v_H^2/v_H^2$ and, for the case $\Delta = 0$:

$$G(\alpha^2, x^2) = \frac{1}{1-\alpha^2} \left[I_1(\alpha^2 x^2) \exp(-\alpha^2 x^2) - I_1(x^2) \exp(-x^2) \right], \quad (11)$$

$$H(\alpha^2, x^2) = \frac{1}{1-\alpha^2} \left[I_1(x^2) \exp(-x^2) - \alpha^2 I_1(\alpha^2 x^2) \exp(-\alpha^2 x^2) \right], \quad (12)$$

$$\begin{aligned} H_0(\alpha^2, x^2) = & \frac{1}{1-\alpha^2} \left\{ \frac{1}{\alpha^2} \left[1 - I_0(\alpha^2 x^2) \exp(-\alpha^2 x^2) \right] \right. \\ & \left. - \left[1 - I_0(x^2) \exp(-x^2) \right] \right\}, \quad (13) \end{aligned}$$

$$G_0(\alpha^2, x^2) = \frac{1}{1-\alpha^2} \left[I_0(\alpha^2 x^2) \exp(-\alpha^2 x^2) - I_0(x^2) \exp(-x^2) \right]. \quad (14)$$

For $\Delta \neq 0$ there are added terms of similar form.

Equation (10) may now be used, in the limit $\omega \approx \omega_{ci}$, to find the values of $k_{\perp a_i}$, the growth rates in the absence of warm plasma, and, finally, the values of Δ required to achieve marginal stability.

Baldwin, Berk and Pearlstein (1920) have shown that depending on the values of R_p/a_i and β , there are three regimes, corresponding to increasingly large values of Δ required for stabilization, that result from solving Eq. (10). These regimes are tabulated in Table I.

TABLE I Definition of Plasma Parameters for Warm Plasma Stabilization of DCLC Mode

<u>Regime</u>	Range of $\epsilon = a_i/R_p$	$\omega/k_{\perp} v_h$	$\Delta_{\min}$
I	$2\epsilon^{\frac{1}{2}} \left(\frac{m}{M} + \frac{\omega_{ci}^2}{\omega_{pi}^2} \right)^{\frac{1}{2}} < \epsilon \leq \beta^2 \left(\frac{v_h}{v_H} \right)$	$\epsilon \left(\frac{v_H}{v_h} \right)$	$2 \left(\frac{v_H}{v_h} \right) \left(\frac{\epsilon}{\beta} \right)^3$
II	$\beta^2 \left(\frac{v_h}{v_H} \right) \leq \epsilon \leq \left(\frac{v_h}{v_H} \right)$	$\left(\epsilon \frac{v_H}{v_h} \right)^{\frac{1}{2}}$	$2 \left(\frac{v_h}{v_H} \right)^{\frac{1}{2}} \epsilon^{3/2}$
III	$\left(\frac{v_h}{v_H} \right) \leq \epsilon$	1	$\frac{1}{2} \left(\frac{v_h}{v_H} \right)^2$

Δ is evaluated at $\omega/k = v_w$

Regime III corresponds to those cases of small diameter plasmas, such as in 2XII B, where the density gradients are steepest and the largest values of Δ are required. In this limit, solving Eq. (10) to lowest order we find $k_{\perp a_i}$ in terms of R_p/a_i . In this limit, in Eq. (10) the term

$$\frac{1}{x} H_0(\alpha^2, x^2) \rightarrow \frac{1}{x^2},$$

so that we find, to this order:

$$\left(\frac{\beta}{2} \right) \frac{1}{x^2} + \frac{1}{x^2} = \frac{1}{\sqrt{2}} \left(\frac{\omega_{ci}}{\omega} \right) \left(\frac{a_i}{R_p} \right) \left(1 + \frac{\beta}{2} \right) \frac{1}{x}, \quad (14)$$

i.e. at $\omega \approx \omega_{ci}$

$$x = \sqrt{2} \left(\frac{R_p}{a_i} \right),$$

$$k_1 a_i = 2 \left(\frac{R_p}{a_i} \right).$$

(15)

Thus even in this limit at marginal stability the instability has been forced to fairly short perpendicular wavelengths ($k_1 a_i \approx 6$; $\lambda_1 \approx a_i$ in 2XIIB). As the density gradient weakens (larger plasma radii) this trend toward larger $k_1 a_i$ will continue, modified somewhat by beta corrections. For example, in regime II

$$k_1 a_i \approx 2 \left(\frac{\beta}{2 + \beta} \right) \left(\frac{R_p}{a_i} \right) \approx \beta \left(\frac{R_p}{a_i} \right).$$

(16)

Utilizing (15) and (16) we may determine all of the parameters defining the marginally stable distribution functions appropriate to these two regimes. Examples of marginally stable distribution functions for these two regimes, for the case $(R_p/a_i) = 3$ (Regime III) and $(R_p/a_i) = 25$ (Regime II) and the beta and T_e/T_i ratios listed are shown in Figs. 4 and 5. Note that in neither case is the loss cone completely filled at the marginal stability point. In fact for large R_p/a_i in this model the marginally stable distribution is actually double-humped (although fortunately Δ is too small to stimulate the double-humped instability mode).

Fig. 4 Example of ion distribution function $F(v_\perp)$ marginally stable against DCLC mode in a plasma of small R_p/a_i .

Fig. 5 Example of ion distribution function $F(v_\perp)$ marginally stable against DCLC mode in a plasma of large R_p/a_i .

curves of R_p/a_i vs Δ with beta as a parameter for marginal stability have been calculated by Berk and Gerber (1976). These are shown in Fig. 6. The strongly stabilizing effects of finite beta for $R_p/a_i > 10$ are evident. These effects can be expected to play an important role as the plasma radius is increased toward the values needed for fusion applications.

Fig. 6 Minimum fraction of warm plasma, $\Delta_{\min}$, needed for stability against DCC mode as a function of R_p/a_i , for various values of β and v_H/v_n .

An additional scaling law can be deduced from the marginal stability condition. This has to do with the relationship between the electron temperature and the value of v_n , in turn coupled to the value of the ambipolar potential, Φ . Recognizing that the source of electron heating is the energy transferred in the course of electron-drag cooling of the ions we may write for the electron heating the equation

$$\frac{d \left(\frac{3}{2} n k T_e \right)}{dt} = \tau_d^{-1} n_H T_H - n k T_e J, \quad (17)$$

where the current $J = v_n n_w$ is the flow of lost particles carrying with them the electron energy drain due to this end loss. The electrons in effect carry out an energy $n k T_e$ per electron. In an isolated plasma Fokker-Planck calculations give $\eta \approx 5.5$. However in 2XIIIB additional effects, such as electrons returning from the end walls increase the value of η to approximately 8.

Equation (17) may be solved in steady-state to yield a relationship between T_e and T_i , given a value for Δ (which determines n_w). Taking $v_L = v_w/L_p$, i.e. the mean transit time of the warm ions, we find for regime III (highest values of Δ):

$$T_e = 3.5 \times 10^{-6} (n_H L_p)^{1/4} T_H^{1/2} \text{ keV} \quad (18)$$

Putting in values representative of 2XIIIB, $L_p = 25 \text{ cm}$, $n_H = 2 \times 10^{14} \text{ cm}^{-3}$ and $T_H = 10 \text{ keV}$, we find $T_e = 150 \text{ eV}$, in approximate agreement with 2XIIIB experimental results. This value, $T_e/T_i = 1.015$, is of course much lower than the value $T_e/T_i = 0.1$ that would be expected for an isolated mirror plasma in the absence of warm plasma stabilization. The low value of T_e/T_i reflects the need for a high value of β for plasmas with steep density gradients ($R_p/a_i = 3$ in 2XIIIB). As R_p/a_i is increased one would expect T_e/T_i to increase correspondingly, as the result of being able to decrease β while still maintaining stability. Preliminary experimental evidence showing that this is indeed the case comes again from 2XIIIB. (Correll, 1979) By changing the plasma injection parameters (neutral beam aiming) and by shifting to hydrogen from deuterium (smaller a_i) the range $2 \leq R_p/a_i \leq 6$ could be covered. Fig. 7 illustrates the increases in T_e and n that resulted, increases which agree closely with theoretical predictions.

Fig. 7 Electron temperature and n in 2XIIIB as functions of R_p/a_i at nearly constant beta.

VI) QUASILINEAR CALCULATION OF THE DCLC MODE AND ITS TRANSPORT CHARACTERISTICS

The foregoing discussion has concerned itself mainly with the requirements for warm plasma stabilization and scaling laws deduced from these requirements. To deal theoretically with the actual processes going on in the plasma as it attains a state of marginal stability Berk and Stewart (1977) have applied

the problem. Each of the computed sets which reproduces the most important aspects of the PXLII experimental data. Their corresponding λ and D diffusion coefficient (Eq. 1) associated with the instability (Eq. 2) of spectral width $\Delta\omega$. $D(v_i)$ is given by

$$D(v_i) = \frac{1}{n_i} \sum_n \left(\frac{1}{v_{ci}} \right)^2 \int \frac{dk_{\parallel} dk_{\perp}}{(2\pi)^3} \left(\frac{v_i}{v_{ci}} \right)^2 \left(\frac{k_{\parallel} v_i}{v_{ci}} \right) \frac{1}{(v_k - n v_{ci})^2 (v_k)^2} \quad (19)$$

Since the fundamental mode is the most unstable, the frequency $\omega = v_{ci}$ is chosen, together with a growth rate derived from the linear dispersion equation (Eq. 10) in evaluating the coefficients. Since the unstable waves must have $v_k > v_{ci}$, it then follows that for $v_i > v_{ci}$ and for fluctuations centered at $\omega = v_{ci}$

$$D(v_i) \propto (v_i/v_{ci})^3 \quad (20)$$

hot particles have scattering rates that are of order $(v_h/v_{ci})^5$ times slower than those of the hole particles.

Using the theory above outlined Berk and Stewart (1977) have calculated the time evolution of important plasma parameters in PXLII, obtaining satisfying agreement with the experimental results. Fig. 8 illustrates such a comparison.

From the quasilinear calculations it is also possible to determine the value of λ required for given plasma parameters and to compare the predicted values with experimental values (determined by measuring the escaping current of warm ions). Fig. 9 illustrates such a comparison, where the data points represent various specific shots where data indicating quiescent plasma conditions were obtained. It can be seen that there is general agreement between the measured and theoretical λ values, although there seems to be a tendency for λ_{exp} to be somewhat lower than that indicated by theory.

Further confirmation of the general quasilinear picture of r.f.-induced velocity-space transport has been obtained by direct measurements of the buildup of a high energy tail in the distribution and comparison of these measurements with the predictions of the quasilinear code. (Turner, 1977) By performing these measurements both during the time that a warm plasma stream was being injected and during the time following the turn-off of that stream a comparison could be made of the rates of turbulent diffusion pertaining under both conditions. In this way it was possible to show directly the reduced turbulence level associated with warm plasma stabilization. The results of this comparison, shown as changes in the energy distribution (changes in amplitude of signals in the channels of the energy analyzer) are shown in Fig. 10. Evident is the depletion of the lower energy channels, coupled with increases in the higher energy channels, as upward diffusion in energy took place.

Having established the correlation between the quasilinear code and the experiment it was possible to fit the data before and after stream turnoff with a diffusion coefficient. Expressed in terms of E_i these values were

$$D(E_i) = \frac{6.6 \times 10^5}{E_i^2} \text{ keV}^2/\text{sec.} \quad (21)$$

Fig. 8 Comparison between ion and electron temperature and plasma density derived from quasilinear calculation and experiment.

during stream stabilization, and

$$D(E_1) = \frac{23.0 \times 10^5}{E_1^2} \text{ keV}^2/\text{sec. following} \quad (22)$$

turn-off of the stream. To be sure the above values are particular to the 2XIIB plasma and much smaller values for $D(E_1)$ under stabilized conditions would be expected for plasmas with the larger R_p/a_i values appropriate for scaled-up systems. Nevertheless the above data are useful in showing the marked reduction in anomalous transport achieved even for a plasma of small R_p/a_i .

Fig. 9 Comparison between theoretical and experimental values of warm plasma density required to stabilize DCLC mode.

Fig. 10 Left hand side: Experimental data on changes in ion energy distribution function in 2XIIB following stream turnoff. Right hand side: Comparison between experiment and 2-D simulation of rates of anomalous diffusion following stream turnoff, as a function of ion energy.

In addition to the above confirmations of the theoretical picture of the DCLC mode and its effects, additional confirmation comes from an analysis of data concerning the properties of the wave itself. As observed in 2XIIIB, (Turner, Powers, and Simonen, 1977) the mode has a sharply defined frequency close to ω_{ci} at the midplane (as corrected for the finite beta of the plasma), has a phase velocity in the direction of the ion diamagnetic drift (correct), and exhibits $k_{\perp} a_i$ values of ~ 6 , in agreement with the theoretical value (Eq. 15), $k_{\perp} a_i = 2(R_p/a_i)$.

Finally, we note that although the dominant observable effect of the DCLC mode will be the transport in velocity space that it causes, there should also be spatial transport associated with the velocity-space transport. Futch (1979) has incorporated quasi-linear effects into a Fokker-Planck radial transport code that also takes into account radial transport effects associated with charge exchange between trapped ions and the injected neutral beams. The result of his incorporation of quasilinear effects into the code is to greatly improve the agreement between both the radial profiles (broadened) and the electron temperature (lowered), as compared to the "classical" Fokker-Planck code in which these effects were not included.

VII) THE HFCLC MODE AND ITS STABILIZATION

For reasons that are only partly understood the HFCLC mode has not been clearly identified as being present in high temperature plasmas such as those studied in 2XIIIB. One reason for this may be that it, too, is subject to stabilization by warm plasma. Thus those circumstances that create a stable state for the DCLC mode, as discussed in the last section, may at the same time weaken or suppress the HFCLC mode.

The theory of the HFCLC mode, put forth in 1965-1966 (Rosenbluth and Post, 1965; Post and Rosenbluth, 1966), shows that it is by nature a convective mode in which transverse ion motion drives a longitudinally propagating electron wave (see Section III). As such there then exist three mechanisms that can act to stabilize it or to reduce its growth. These are:

1. Landau damping effects from finite electron temperature.
2. Limitation of plasma length and minimization of reflection at the ends of the plasma.
3. Warm plasma stabilization.

These mechanisms will be discussed in turn following a brief outline of the theory of this mode.

In the limit of low T_e , i.e. not including Landau damping effects, the dispersion equation for the HFCLC mode can be written

$$0 = 1 + \left(\frac{\omega_{pe}^2}{\omega^2 - \omega_{ce}^2} \right) - \left(\frac{\omega_{pe}^2}{\omega^2} \right) \left(\frac{k_{\parallel}^2}{k^2} \right) - \frac{\omega_{pi}^2}{k^2 v_i^2} F_i \left(\frac{\omega}{k v_i} \right) \quad (23)$$

where the ion driving term $F_i(y)$ is given by the integral

$$F(y) = -2 \int_0^{\infty} dx \left(\frac{\partial \psi}{\partial x} \right) \frac{1}{(1-x/y)^{1/2}} \quad (24)$$

where $f(v) = n(v_1^2/v_1^2)$ is the normalized number of ions with a given perpendicular energy

$$\psi(x) = \hat{\psi} \int_{-\infty}^{\infty} f_0(v_1^2, v_{||}^2) dv_{||}, \text{ normalized}$$

to unity, i.e., $\int_{-\infty}^{\infty} \psi(x) dx = 1$. Instability occurs only for those values of $x = \omega/kv_1$ where $\text{Im}(y) > 0$.

The HFCLC mode has a relatively weak spatial growth for well spread loss-cone distributions, with an e-folding length given by

$$L = 2a_1 \left[1 + \frac{\omega_{ce}^2}{\omega_{pe}^2} \right]^{\frac{1}{2}} \left[|y \text{Im} F(y)| \right]^{-1} \text{ cm} \quad (25)$$

Typical values of $|y \text{Im} F(y)|_{\text{max neg}}^{-1}$ are ~ 10 , giving minimum e-folding lengths of order $20 a_1$, uncorrected for electron Landau damping.

Horton (1967) has shown that finite T_e and the accompanying electron Landau damping effects restrict the available range of y and increases the e-folding lengths for this mode. In his calculations using an ion distribution function modeling a $R_m = 2$ mirror loss-cone distribution, for example, he found that increasing T_e/T_i to 0.1, compared to the $T_e = 0$ case increased the e-folding length by a factor of 2 while decreasing the area of instability in ω, k space by roughly a factor 4.

Similarly, adding a modest amount of warm plasma to a loss cone distribution will stabilize the mode, even in the absence of finite T_e effects. Using the same model ion distribution function as employed by Horton, for which $\psi(x) = x e^{-x}$, and adding a maxwellian component $\psi(x) = e^{-x}$ of lower mean energy Post (1976) found stabilization at relatively small values of Λ . Table II, reproduced from this work illustrates the values obtained for Λ . In terms of the relative mean ion energies of the warm maxwellian and the loss-cone distribution functions, the stability criterion is

$$\sum_j f_j (Z_j/A_j) y_j^2 \text{Im} F(y_j) > 0 \text{ stable}, \quad (26)$$

where f_j is the fractional density of the j th species, $\sum_j f_j = 1$. Note that by contrast with the DCLC mode the warm plasma components need not be of the same ion type as the main distribution to provide a stabilizing influence. For maxwellian warm plasmas

$$\text{Im} F(y) = 2\pi^{\frac{1}{2}} y \exp(-y^2) > 0.$$

Under conditions where warm plasma stabilization is complete there is of course no further reason for concern about this mode. It is conceivable, however, that situations could arise where stability against the DCLC mode was achieved (for example at large R_p/a_1) without there being enough warm plasma present to stabilize the HFCLC mode. In such cases the critical issue becomes the matter

of the growth lengths of this mode and whether or not it undergoes reflections at the ends of the plasma of sufficient magnitude to permit it to become an absolute instability. Early calculations (Aamodt and Book, 1966) of reflections seemed to indicate that provided steep density gradients were avoided the mode would remain convective, i.e. innocuous for plasma lengths less than of order 10 e-folding lengths. More recent calculations by Gerver (1979), however, in investigating slower growing but longer wavelength parts of the unstable spectrum in y-space indicate that the mode can become absolute. However, whether this will actually occur in practice, and what its actual consequences in terms of anomalous transport effects will be is not clear at this time. Certainly if the plasma length is restricted while at the same time finite T_e and some warm plasma stabilization is employed one might expect to find operating regimes where the HFCLC mode does not interfere appreciably with confinement. Future tests in scaled-up mirror systems such as MFTF at Livermore should be able to answer these questions.

TABLE II Relative Temperatures, Δ , and Relative Energy Density for Warm Plasma Stabilization of HFCLC Mode

T_W/T_H	Δ	Relative Energy Density
1.0	0.33	0.25
0.5	0.15	0.056
0.3	0.077	0.017
0.2	0.057	0.0077 Min Δ
0.18	0.057	0.0070
0.16	0.058	0.0057 Min. Energy Density
0.12	0.076	0.0068
0.10	0.104	0.0078

VIII) THE ALFVEN ION CYCLOTRON MODE

The last of the three modes that we have earlier singled out for special consideration is the AIC mode (Davidson and Ogden, 1975), a mode feeding on the anisotropy of the plasma in distribution functions rather than on the existence of the loss-cone hole. Present indications are that as long as well randomized ion distributions are used, contained in mirror cells of reasonable mirror ratio (for example $R_m = 5$ in the presence of the plasma), the AIC mode should not pose a serious threat to the achievement of fusion plasmas. The effect of this mode is to impose a beta limit on the plasma above which the AIC mode becomes absolute. For model ion distribution function (bi-maxwellian) approximating that in a high beta mirror plasma it is found that the stability boundary of the AIC mode is virtually the same as the limit imposed on beta by magnetostatic equilibrium (i.e. the mirror mode). Fig. 11 summarizes the results for such a case, for a vacuum $R_m = 2$. Further theoretical study of the AIC mode and of stabilizing effects not yet included in the theory is underway at Livermore. Since the mode has not yet been observed experimentally it is not possible to draw firm conclusions relative to its significance at this time, but present indications are that it will not be of major importance.

Fig. 11 Critical stability boundaries in $\beta_{\perp}, \beta_{\parallel}$ for mirror mode and AIC mode.

IX) NEGATIVE ENERGY MODE

For completeness we will mention another theoretically predicted mode, the negative energy mode, that can in principle appear under certain conditions. So-called negative energy waves can propagate in a plasma in regions near cyclotron resonances where the slope of the dispersion curve as it crosses the axis is opposite to that for a maxwellian. Loss cone distributions without a warm plasma component have such slope characteristics. Negative energy waves are not of themselves necessarily unstable, but can be rendered unstable by dissipation, for example by loss of wave energy through the mirrors. An alternative picture is to recognize that these waves are convective instabilities (lower frequency relatives of the HFCLC) rendered absolute by reflections at the turning points (high-order resonances) at the ends of the plasma.

Typically slow growing, the negative energy modes are not a priori expected to be as important in their effects on confinement as the DCLC mode. They are sensitively dependent for their existence on details of the plasma distribution functions, on scale lengths and on the presence or absence of warm plasma, generally being stabilized at lower values of Δ than those required to stabilize the DCLC, except at very large R_p/a_i . They are also stabilized by limiting the magnetic scale length of the plasma, which prevents the formation of an unstable normal mode. Critical scale lengths for stabilizing the negative energy modes without the need for warm plasma have been calculated by Berk, Pearlstein and Cordey (1972), including beta-dependent stabilizing effects. Fig. 12 shows a plot of these lengths for the two most unstable modes of this instability, as a function of $\omega_{pi}^2/\omega_{ci}^2$. The scale lengths found are seen to be of order 10 to 100 a_i , depending on plasma parameters, i.e. of order of the critical lengths for the unstabilized HFCLC mode.

Fig. 12 Critical lengths for negative energy modes as a function of plasma density.

Because of their relative fragility and since certain stabilizing effects have been left out of theoretical treatments of the negative energy modes, it is not clear whether they will ever be of importance in mirror confinement. So far as is known such modes have not been detected in mirror experiments. However, until scaled-up experiments are available in which these modes might be more prone to appear, their possible importance to mirror systems must remain an unanswered question.

X) THE TANDEM MIRROR AND ITS EFFECT ON HIGH FREQUENCY INSTABILITIES

We will not here specifically discuss the aspects of possibly anomalous transport associated with MHD activity or particle drifts in the central cell of the Tandem Mirror. There is however one aspect of the TM idea that is of particular importance to the issue of high frequency instabilities in the mirror plugs of the TM. Recall that a TM system consists of a central solenoid confinement region bounded on each end by one or more mirror cells containing a high temperature plasma the ambipolar potentials of which act to confine maxwellian central plasma. The ions of the central plasma, diffusing upward in energy as a result of collisions, are finally able to escape through the barriers. This escaping flux of ions thus represent an intrinsic source of "warm" plasma ions that can contribute to the stability of the mirror end plugs. There is already preliminary evidence that this process is operating in the Livermore TMX Tandem Mirror Experiment. Whether this stream is alone sufficient to achieve stability in a full-sized Tandem Mirror system will clearly depend on the machine parameters of that system. The plugs of TM systems with sufficiently long central cells, or ones operating in the so-called two-component mode (Post, and co-workers, 1973) would be expected to be adequately stabilized by the warm plasma streams emerging from the central solenoid. In this latter case the stream intensity should be high enough to suppress all the high frequency modes that have been here considered. In the central cell itself both the maxwellian nature of the plasma and the absence of any net current along the field should mean that the only concerns relating to

anomalous transport are those arising from low-frequency MHD modes or from drift-induced radial transport arising from the lack of symmetry of the confining fields. These latter problems are under study and believed to be solvable by care in the magnet design.

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