AXISYMMETRIC, WALL-STABILIZED TANDEM MIRRORS

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ABSTRACT. The possibility of an axisymmetric tandem mirror in which stability accrues from wall stabilization is discussed. It is found that the stability requirements are compatible with thermal barrier requirements so that the thermal barrier plug cell can also provide stabilization. Thus, a single axisymmetric end cell can plug and stabilize a high-beta plasma solenoid. Self-stabilization of the central cell and other magnetic configurations are also discussed.

1. INTRODUCTION

A tandem mirror is a linear device in which confinement in a central solenoid results from electrostatic potential 'plugs' located at either ends of the device. The efficiency of creating the potentials can be improved by the interposition of a ‘thermal barrier’ between the central solenoid and the plugs [1]. A thermal barrier is an along-the-field-line potential depression that serves to thermally insulate the central-cell ‘thermal’ electron population from a suprathermal population in the end plug. The potential depression results from the interposition between the central cell and the plug of a dense, localized hot-electron population.

MHD stability in a tandem mirror usually derives from the presence of a quadrupole, ‘minimum-B’ cell containing high plasma pressure. This ‘anchor’ is required primarily for stabilizing modes with low azimuthal mode number (m), since higher-order modes are stabilized by central-cell Finite Larmor Radius (FLR) effects. However, the presence of a quadrupole cell necessarily leads to a more complicated coil set and a flux tube distortion away from axisymmetry. The non-axisymmetry of this system constrains equilibrium and can produce radial transport. Furthermore, the large curvature present in the transition region of these devices enhances the drive for MHD and electrostatic (trapped-particle) ballooning instabilities.

One approach to obtaining enhanced axisymmetry is the ‘axicell’ arrangement [2]. In this geometry, the thermal barrier and plug can be produced in the same mirror cell, the ‘axicell’, through the use of sloshing ions. The quadrupole anchor can then be located outside of the confinement region. This arrangement provides axisymmetric ion confinement, but is more susceptible to trapped particle modes than to other schemes [3].

Recent work by Berk et al. [4] suggests the possibility of the use of a wall stabilization mechanism for m = 1 curvature driven modes. The source of stability is the image currents generated by placing the wall (or properly shaped conductors) in close proximity of a high-beta axially localized plasma. This work contains a low-beta approximation. Wall stabilization has also been studied in the MHD limit for arbitrary beta by D'Ippolito and Myra [5] and by Pearlstein and Kaiser [6].

In this paper, we discuss schemes for obtaining a totally axisymmetric tandem mirror in which stability derives from the aforementioned wall stabilization. It will be shown that the stabilization criteria are compatible with the requirements for the hot-electron population of a thermal barrier. The hot-electron population present in the axicell could both create the thermal barrier and provide MHD stability for m = 1 curvature driven instabilities. Higher m-modes can be stabilized by FLR stabilization deriving from both the axicell sloshing ions and the central-cell thermal ions. Thus, the need for quadrupole anchors could be eliminated. The maximum beta that can be confined in the central cell would then be determined by MHD ballooning and trapped-particle stability requirements.

Additionally, we consider the possibility of the central cell being self-stabilized by the same wall effect. The isotropic nature of the central-cell plasma will weaken the wall response and very high beta is required for stabilization. Nevertheless, this scenario raises the possibility of a high-beta, linear confinement device.
containing simple axisymmetric plugs. Such a device would clearly produce a very desirable arrangement for a fusion reactor.

Section 2 discusses the axicell stabilization requirements. Section 3 discusses central-cell wall stabilization and in Section 4 other interesting anchor arrangements based on a high-beta ion population are discussed. Section 5 gives the conclusions.

2. WALL-STABILIZED AXICELL REQUIREMENTS

We consider an axicell thermal barrier arrangement [2] in which the magnetic field and potential are as shown schematically in Fig. 1. The thermal barrier is formed by a disc-shaped high-beta hot-electron plasma that is mirror confined near the axicell field minimum (point b).

From the point of view of thermal barrier formation, the axicell midplane field should be low in order to minimize the trapping rate of central-cell ions that traverse the barrier region. Furthermore, the electron temperature must be high in order to eliminate the expulsion of hot electrons from the barrier region and to reduce electron collisionality. Thus, high beta is desirable. Additionally, the hot-electron power balance requires a minimization of the hot-electron volume, which requires maximum hot-electron anisotropy. For example, in the MARS tandem mirror reactor design study [7], the thermal barrier was formed by electrons with a mean energy of 820 keV, a beta of 50% and an anisotropy ($A = P_p / P_t$, the ratio of perpendicular to parallel pressures) of $A = 4$. The self-consistent magnetic field at the thermal barrier was $B_b = 1.2$ T.

The stability requirement from MHD for an isolated mirror cell obtains a simple form for a sharp boundary pressure model [6] and is given by

$$\omega^2 = \omega_{\text{vac}}^2 = \frac{2p}{B_{\text{vac}}^2} \frac{r_{\text{vac}}''}{r_{\text{vac}}} + \frac{1}{8} \left( \frac{\beta_{\text{vac}}}{1 - \beta_{\text{vac}}} \right)^2 (1 - p/B_{\text{vac}}^2)$$

where $\rho$ is the mass density, $\beta_{\text{vac}} = 2P_t/B_{\text{vac}}^2$, $p = (P_t + P_r)/2$ and $A = (R_p^2 + R_r^2)/(R_r^2 - R_p^2)$, with $R_p$ and $R_r$ the respective plasma and wall radii. Primes represent along-the-field-line derivatives. This expression is valid at arbitrary beta and contains the long-thin approximation ($\partial / \partial z \ll \partial / \partial r$) and a sharp-boundary pressure model. The first term in the integral will be recognized as the MHD drive due to the vacuum curvature and the second term represents the effect of wall stabilization. This expression is valid for large $\Lambda$, i.e. for the wall close to the plasma edge. The third term is small and can be shown to be destabilizing. For the wall right at the plasma edge, $\Lambda \to \infty$, and this term is zero; the first two terms of Eq. (1) can be viewed as providing a necessary condition for stability. This condition can thus be written as $S > 1$, with $S$ defined as

$$S = 1/16 \int_{-L}^{L} ds (1 - p/B_{\text{vac}}^2)(\beta_{\text{vac}}/(1 - \beta_{\text{vac}}))^2$$

$$\times \left[ \int_{-L}^{L} ds \frac{p}{B_{\text{vac}}^2} \frac{r_{\text{vac}}''}{r_{\text{vac}}} \right]^{-1}$$

Stabilization clearly requires a high-beta axially localized disc. We note that at high beta the plasma disc will dig a diamagnetic well and axial localization can be obtained with a more isotropic plasma, whereas
at low beta axial localization requires high plasma anisotropy.

In order to continue further, we must impose a dependence for \( P_1 \), i.e. \( P_1(B) \). (The assumption of a sharp boundary eliminates radial dependence.) The field-line derivatives can then be evaluated using the long-thin equilibrium condition

\[
2P_1 + B^2 = B_{\text{vac}}^2
\]

(3)
given a vacuum field profile. Additionally, the dependence \( P(B) \) follows from the axial pressure balance:

\[
\frac{\partial}{\partial B} \left( \frac{P(B)}{B} \right) = -\frac{P_x}{B^2}
\]

(4)

A simple pressure model that may be used to evaluate the stability requirements is

\[
P_1(B) = P_0 \left( 1 - B^2/B_1^2 \right) \quad B < B_1
\]

\[
= 0 \quad B > B_1
\]

(5)

with

\[
P_0 = \frac{\beta_0}{2} \frac{B_{\text{vac}}^2 B_1^2}{B_1^2 - B_{\text{vac}}^2 (1 - \beta_0)}
\]

\( \beta_0 \) is the midplane vacuum beta (\( \beta_0 \equiv 2P_1(s = 0)/B_{\text{vac}}^2 \)) and \( B_{\text{vac}} \) and \( B_1 \) are respectively the vacuum midplane field and the field at which the pressure goes to zero. This distribution function has been called an ideal distribution, since it does not contain a mirror-mode limit. In this model, the pressure goes to zero at \( B = B_1 \), which is less than the peak mirror field \( B_m \). (We define a vacuum mirror ratio \( R_1 \equiv B_1/B_{\text{vac}} \), which is less than the axicell mirror ratio \( R_T \equiv B_m/B_{\text{vac}} \).) Thus, \( R_1 \) determines the localization of the hot-electron disc. From Eq. (3), we can obtain \( P_1(B) \) and the midplane anisotropy

\[
P_\alpha/P_0 = \frac{B_1 + B_0}{B_1 - B_0}
\]

(6)

with \( B_0 \) the midplane (beta depressed) field. Notice that, as beta increases, \( B_0 \) will decrease and the anisotropy will approach one.

Assuming a parabolic vacuum field, we can now evaluate Eq.(2). Figure 2 displays the stabilization factor, \( S \), as a function of midplane beta. Distributions with a small mirror ratio, \( R_1 \), have a larger anisotropy and exhibit stronger stabilization. Since these profiles can have the same spatial extent, it is clearly anisotropy (not axial localization) that determines stability.

A second model pressure distribution that allows pressure to extend out to the mirror peak is given by

\[
P_1(B) = nP_0(B/B_m)^2(1 - B/B_m)^{n-1}
\]

(7)

with

\[
P_0 = \frac{\beta_0 B_m^2}{2n(1 - \beta_0)(1 - \sqrt{1 - \beta_0}/R_{\text{vac}})^{n-1}}
\]

\( R_{\text{vac}} = B_m/(B_0 \sqrt{1 - \beta_0}) \). This distribution has an anisotropy given by

\[
P_\alpha/P_0 = n/(B_m/B_0 - 1)
\]

(8)

Again, assuming a parabolic vacuum field, we evaluate the stability factor. Curves for \( S = 1 \) and \( S = 2 \) are shown in Fig. 3 as a function of midplane anisotropy. The circles represent the results for the ideal distribution (shown in Fig. 2). The very close agreement between these two pressure models implies that the stabilization depends on the midplane anisotropy and is not sensitive to the exact \( P_1(B) \) dependence. Some deviation between the pressure models is seen for \( S = 2 \).

The pressure model of Eq. (7) contains a mirror-mode limitation on \( \beta_0 \) given by

\[
\frac{\beta_0}{1 - \beta_0} < \frac{2(B_m/B - 1)}{(n + 1 - 2B_m/B)}
\]
The mirror-mode limit, which is a boundary imposed by equilibrium, is indicated by the dashed curve in Fig. 3. Thus, we observe a stability window between the wall stabilization requirements and the requirements of equilibrium at high beta.

The MARS thermal barrier operating point, $\beta_0 = 50\%$, $A = 4$, is also indicated in Fig. 3. We see that this operating point falls at $S \approx 2.5$.

A last point of interest that can be gleaned from this distribution is the effect of a 'sloshing' hot population. From Eq. (7), we can determine that the pressure will peak at $B_{\text{slosh}} = 2B_m/(1 + n)$. For large $n$ and large mirror ratio, we then find

$$R_{\text{slosh}} \equiv \frac{B_{\text{slosh}}}{B_0} = 2P_1/P_\parallel$$

For $P_1/P_\parallel < 2$, the pressure peak moves off the midplane and Fig. 3 indicates an increase in the required midplane beta for stabilization. Thus, if the stability of anisotropy-driven modes requires a sloshing character to the high-beta component, we can still obtain a stable regime.

In summary, we observe the existence of a high-beta stability window which may be bounded by equilibrium requirements. We note that there is a decreased stabilization as the wall moves back from the plasma. Berk et al. [3] showed that at low beta the stabilization is proportional to $(r_p/r_w)^2$ and therefore is expected to be strong only when the plasma edge is near to the wall. Additionally, in a tandem mirror the added instability drive that comes from the central cell would further increase the required $S$ value. An appropriate value of the stabilization factor, $S$, might therefore be $S \approx 2$.

Up to now, we have considered the MHD result of Ref. [6] and one could question the applicability of the MHD formalism. The energy of the disc forming hot electrons is expected to be in the range $500–900$ keV, and at high beta their drift frequency $\omega_D \approx \beta_h \omega_{*h}$ ($\omega_{*h}$ is the hot-species diamagnetic drift frequency) will greatly exceed the central-cell diamagnetic drift frequency $\omega_{*c}$ and the MHD growth rate that characterizes MHD modes. Thus, the electrons form a 'hot' species in the accepted EBT terminology.

One can show, however, that the failure to satisfy a 'decoupling' condition will lead to an MHD-like response of the hot ions (Ref. [8], Eq. (73)) even if the core beta is negligible. This decoupling condition sets the requirement

$$\gamma_{\text{MHD}}/\omega_{kh} > 0.5$$

with

$$\gamma_{\text{MHD}}^2 = \gamma_a^2 (1 - S) + \gamma_c^2$$

where $\omega_{kh}$ is the hot-electron curvature drift frequency and $\gamma_a$, $\gamma_c$ are respectively the axicell and central-cell MHD growth rates. For stability, $S > 1$ and $|\gamma_a^2 (1 - S)| > \gamma_c^2$, so that $\gamma_{\text{MHD}}^2 < 0$. We can estimate

$$\gamma_{\text{MHD}}^2/\omega_{kh}^2 \approx \frac{T_c}{T_a} \left( \frac{L_a}{r_a} \right)^2 \frac{V_a}{V_c} \frac{n_{eh}}{n_c} \frac{1}{(k_i \rho_i)^2}$$

with the subscripts $c$ and $a$ representing central-cell and axicell quantities; $V$ is the volume, $k_i$ is the wave number $(m/r_c)$, $L_a$ is the axicell half-length and $\rho_i$ is the ion gyroradius. We have approximated the curvature as $k_a \approx r_a/L_a^2$. Estimating $T_a/T_c \approx 25$, $L_a/r_a \approx 5$, $V_c/V_a \approx 20$, $n_c/n_{eh} \approx 10$, $k_i \rho_i \approx 0.013$, $B_c/B_a \approx 2$, we obtain $\gamma_{\text{MHD}}/\omega_{kh} \approx 6$.

A critical issue in the above scheme is startup. The hot electrons in the axicell thermal barrier may be started up in the fashion of the EBT experiment [9]. This means that the electrons could be heated up to sufficient energies in order to decouple from the core plasma while maintaining a low core beta (below the Van Dam-Lee limit) and a sufficiently low core density preventing the hot-electron interchange mode (the Berk-Dominguez mode). Once the hot-electron beta exceeds the beta required for wall stabilization, the sloshing-ion beams could be turned on since they will
be stabilized by the electrons for the \( m = 1 \) mode and the core electron density could drop.

Higher azimuthal mode numbers \( (m \gg 2) \) can be stabilized by ion FLR effects deriving either from the sloshing ions that are present in the axicell or from the central-cell ion population. Considering the sloshing ions, Berk et al. [4] suggest that the high-beta FLR stabilization term should enter the dispersion relation in the following manner:

\[
\omega^2 = \gamma_{\text{MHD}}^2 \left( 1 - S_m \frac{\beta_i k_i^2 \rho_i^2}{2\kappa \Delta} \right)
\]  

(12)

with

\[
S_m = S \left( (R_p/R_w)^2 m + 1 - |m| / |m| \right)
\]

where \( \beta_i \) is the ion beta, \( k_i \) is the perpendicular wave vector \((k_i = (m^2 - 1)/r_i^2)\), \( \kappa \) is the curvature, \( \Delta \) is the pressure gradient scale length, and \( g \) is a geometric factor representing the fact that the ions which supply the FLR are sloshing and fill a larger volume than the electrons (which dominate the drive and wall term).

If wall stabilization is significant only for \( m = 1 \) modes (for example if the wall is replaced by an \( m = 1 \) conductor), we can set \( S = 0 \) and obtain approximate requirements for FLR stabilization:

\[
\left( \rho_i/r_a \right)^2 > \frac{1}{(m^2 - 1)} \frac{2\kappa \Delta}{g \theta_i}
\]  

(13)

We will take \( \kappa \approx r_a/L_a^2 \), with \( L_a \) the axicell half-length. Then, for \( m = 2 \) and typical parameters, \( r_a/L_a = 0.1 \), \( \Delta = r_a/5, g = 2, \beta_i = 0.1 \), we find \( \rho_i/r_a > 0.08 \) for barrier parameters of \( B = 1.5 \) T, \( r_a = 0.7 \) m. This would require a midplane perpendicular energy of \( E_{ll} < 226 \) keV for tritium and \( E_{ll} > 340 \) keV for deuterium. The total ion energy at the sloshing-beam injection point would be approximately double this value. In the MARS study the upper limit on energy based on adiabaticity was about 700 keV [7]. We note that the energetic 'sloshing ions' form a second 'hot' species and for MHD coupling must satisfy the same criterion (Eq. (10)) as the electrons.

In addition to the high-beta FLR term that enters the instability drive (Eq. (12)), there is the usual low-beta FLR term that enters the dispersion relation as a charge separation term (and can create stable negative energy waves). Taking account of this term will further reduce the stringency of the FLR stability requirement.

With the thermal barrier electrons providing stabilization for \( m = 1 \) modes and the sloshing ions providing stability for \( m > 1 \), the axicell would be stable to all instability. The central cell could then be started up, which would add some MHD instability drive as well as additional FLR stabilization. It is important to note that when the axicell dominates (and stabilizes the MHD drive), the central cell enters the decoupling condition as a result of inertia. From Eq. (11) we observe that the central-cell temperature does not enter the decoupling condition. This means that during startup the axicell can stabilize the relatively cool central cell (when Eq. (10) is satisfied). Of course, if the central cell is self-stabilized, as discussed below, we need not consider the decoupling condition.

3. CENTRAL-CELL WALL STABILIZATION

Tandem mirror economics favours high-beta central-cell operation. High beta permits designs to lower the field, down to a limit set by alpha adiabaticity, and to raise the plasma pressure, up to a limit set by neutron wall loading [10]. Additionally, at high beta with a relatively sharp boundary pressure profile, as is desired for wall stabilization, classical conduction within the edge region can limit the desired edge gradient. These effects are examined in Ref. [10] and a central-cell beta in the range of 80–90% is found to be favourable.

The strength of the stabilization depends on the spatial localization of beta within the mirror. This localization is weak for the isotopic central-cell plasma but can become significant at high beta. The only non-isotropic component in the central cell is the hot alpha particles produced by D-T fusions; these are only mildly anisotropic and account for less than 20% beta. (The anisotropy of the alphas can be enhanced through the use of polarized nuclei, but this possibility will not be analysed here.)

Since the central-cell plasma is isotropic, the pressure becomes independent of \( B \) and axial position, which will greatly simplify the along-the-field-line integrals in Eq. (1). In this limit, \( P_i(B) = P_i(B) = P_0 \).

To evaluate Eq. (2), we choose a parabolic field shape, \( B(z) = B_{w0} + (B_m - B_{w0})(z^2/L^2) \), with \( L \) the central-cell ramp length. Since the central-cell mirror ratio is large, typically greater than 8, we can allow the limit on integration to extend to infinity in order to obtain an analytic estimate of \( S \). Additionally, to obtain this estimate, we will estimate the factor \( (1 - \rho/B_i(s)) \to (1 - \beta_i/2) \) and thus remove this factor from the integral.
FIG. 4. Stabilization factor versus beta for an isotropic plasma confined in a high-mirror-ratio cell (such as the tandem mirror central cell). Both the analytic and the numerical solutions are shown.

The integral along the axis is performed with the aid of the residue theorem and yields the stabilization factor $S$,

$$S = \frac{16}{3} \frac{(1 - \beta_0/2)}{\beta_0} \cdot \left[ 1 + \frac{(1.5 \beta_0^{-1/2} - 1.25)}{(1 - \beta_0^{-1/2})^{1/2}} - \frac{(1.5 \beta_0^{-1/2} + 1.25)}{(1 + \beta_0^{1/2})^{1/2}} \right]$$

Figure 4 shows $S$ versus $\beta_0$ from Eq. (14) as well as the result obtained from a numerical integration (including the $(1 - p/B^2_0)$ term). The $S = 1$ boundary is seen to occur at $\beta_0 = 83\%$, and above this beta value the curve rises steeply. Notice that the parabolic field-scale length has dropped out of the ratio, indicating an insensitivity to the axial field profile.

Recently, studies have been performed by Potok [10] to model a tandem mirror with a central-cell beta near one; it is shown that the optimum beta is $80-90\%$. As indicated in Ref. [10], at high beta the minimum allowable field is set by alpha adiabaticity requirements and the maximum field is set by neutron wall loading.

In addition, classical confinement (of particles and energy) serves to limit the edge pressure gradient (recall that we desire $r_p/r_w \leq 1$). For $\beta_{cc} = 90\%$, $n_c = 1.2 \times 10^{14} \text{cm}^{-3}$, a plasma radius of 1.7 m, an ion temperature of 30 keV and a solenoid vacuum field of 2 T, Potok [10] finds a resulting neutron wall loading of 3.5 MW-m$^{-2}$. Alternatively, he finds for a vacuum field of 2.5 T and a plasma radius of 1.3 m a wall loading of 6.9 MW-m$^{-2}$ at the same electron temperature. The MARS design [7] has a central-cell field of 4.7 T and a neutron wall loading of 4.2 MW-m$^{-2}$. Thus, high beta can significantly reduce the central-cell field requirements and raise neutron wall loading.

4. OTHER ANCHOR CONFIGURATIONS

A high-beta hot-electron or ion disc is subject to high-beta anisotropy driven instability, such as the relativistic Whistler [11] and cyclotron maser instability for electrons [11, 12]. At high beta, relativistic effects will tend to stabilize the Whistler modes [11]. For ions the Alfvén ion cyclotron (AIC) instability [13] will present a limitation on beta. AIC modes can be stabilized by axial localization [13], although this process is limited by the requirements on ion adiabaticity.

Utilization of ion discs would be desirable since hot ions produce strong FLR stabilization. Two possible...
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Ion disc magnetic configurations are shown schematically in Fig. 5. In these configurations, the ion disc forms an inside anchor (preferable to an outside anchor because of superior stabilization properties for trapped-particle modes).

In Fig. 5(a) the anchor cell is located beyond the central-cell choke coil. The high-beta anchor ion population would require an axial extent of 10 to 20 ion gyroradii for adiabaticity. AIC stability would also require an axial extent of not more than this [14].

The power requirements can be minimized by pumping the passing ions so as to prevent local trapping. For example, for \( B_{\text{vac}} = 2 \, \text{T} \), \( n_a = 7 \times 10^{13} \, \text{cm}^{-3} \), \( T_i = 500 \, \text{keV} \) (\( \beta = 0.6 \)), the power required (drag and pitch-angle scatter) is about 3 MW-m\(^{-3}\). The pumping requirements are increased in proportion to the effective anchor volume, which might be about 5 m\(^3\). This arrangement will increase the length of the end-cell region, increasing the trapping current of the passing ions and thereby the minimum central-cell length required for ignition.

The pumping requirements can be diminished by placing the anchor on the inside of the choke coil (Fig. 5(b)). A small throttle coil creates a mirror (mirror ratio about 2–3) between the central cell and the anchor so as to confine most of the energetic central-cell alphas, which would otherwise suffer from loss of adiabaticity. In this case, however, the density rises to the central-cell level, which increases the power requirements for the hot ions. For a hot-ion density of \( 7 \times 10^{13} \, \text{cm}^{-3} \) and a core-ion density of \( 2 \times 10^{14} \, \text{cm}^{-3} \), the required power for the anchor is about 9 MW-m\(^{-3}\) of volume.

5. CONCLUSIONS

The use of wall stabilization for \( m = 1 \) curvature driven modes is seen to present the possibility of a completely axisymmetric tandem mirror reactor that can operate at high betas and can have compact and relatively simple end plugs. The requirements are seen to be compatible with a thermal barrier. Furthermore, there is a possibility for the central cell to be self-stabilized, which would permit beta near one and eliminate the possibility of trapped-particle modes.

Wall stabilization could be obtained through the use of properly shaped and positioned conductors. This possibility would cut down on sputtering and wall loading, and would increase the access to the plasma for neutral beams.

Finally, Fleischmann et al. [15] pointed out that the finite conductivity of the walls can give rise to a residual instability (on the resistive decay time-scale) that would require some feedback stabilization.

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