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GLOBAL STRUCTURES OF ALFVÉN-BALLOONING MODES IN MAGNETOSPHERIC PLASMAS

BY

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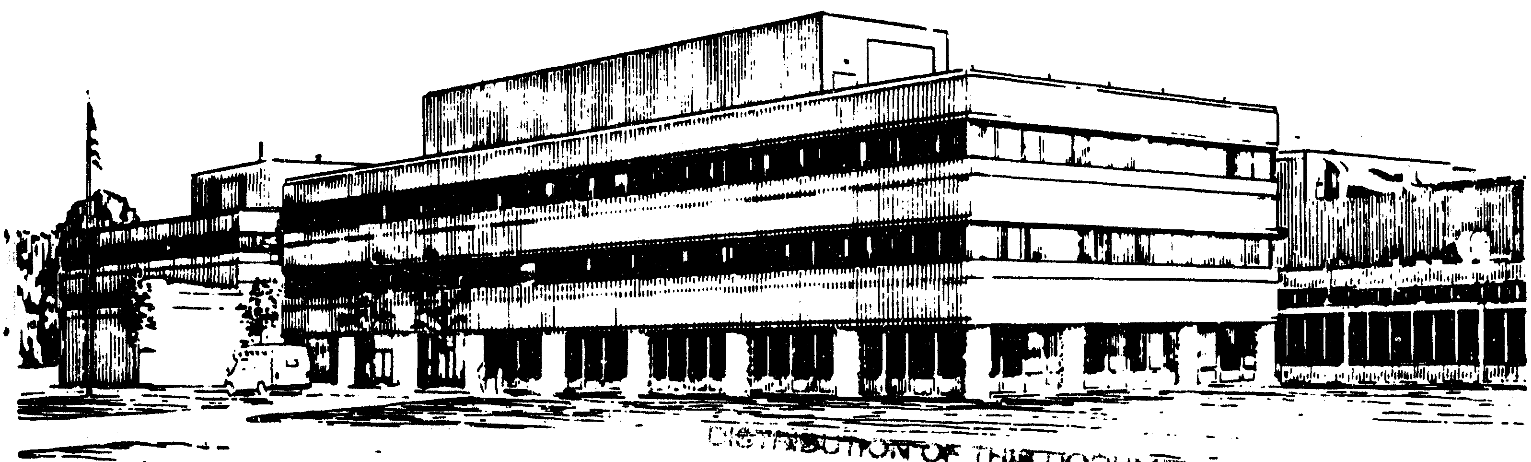

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Global structures of Alfvén-ballooning modes in magnetospheric plasmas

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Abstract

We show that a steep plasma pressure gradient can lead to radially localized Alfvén modes, which are damped through coupling to filed line resonances. These have been called drift Alfvén balloning modes (DABM) by [3] and are the prime candidates to explain Pc4-Pc5 geomagnetic pulsations observed during storms. A strong dependence of the damping rate on the azimuthal wave number m is established, as well as on the equilibrium profile. A minimum azimuthal mode number can be found for the DABM to be radially trapped. We find that higher m DABMs are better localized, which is consistent with high-m observations.

Introduction

Magnetospheric pulsations of the Pc4 and Pc5 types have periods on the order of 100 sec, and are often observed during a geomagnetic storm. Theoretical studies such as those by Chen and Hasegawa [3] etc. have identified the pulsations as drift Alfvén balloning mode instabilities excited via wave-particle interactions with resonant highly energetic protons. It was found that their growth rate is maximized for azimuthal numbers $m \approx 100$ [1] and for a field-aligned structure that is odd with respect to the equator. These results fit quite well with satellite observations such as those of Takahashi [4]. However, to the authors' knowledge, previous studies of these instabilities have been restricted to the one dimensional eigenvalue problem along the field line, i.e. no attempt was made to characterize the radial structure of these modes. It is obvious, that for the DABM to be a realistic instability candidate, it must be radially localized in the magnetosphere. This paper,

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to the best of our knowledge, is the first that deals with these matters. As it will be shown, a plasma pressure gradient, like the one suggested by data in [5], could create a potential well in the radial direction which localizes the global Alfvén-ballooning mode. In the picture developed here, the global mode suffers a small but finite damping due to absorption of the wave energy at the flux surface whose file line resonant frequency matches the global mode eigenfrequency. Both radial localization and damping depend on m as well as the steepness of the pressure gradient. The damping rate is negligible for moderately high m.

Theoretical framework

In this paper we will be using a model equilibrium that is axisymmetric. The magnetohydrodynamic (MHD) momentum equations for the components η_ϕ,η_ψ of the perturbed perpendicular vector potential, were expressed in magnetic coordinates (ϕ,χ,ψ) for the cold plasma case in [2, 6]. For the hot plasma case they are the following:

$$D_\phi\eta_\phi - \frac{\kappa_\psi}{B_0^2}(p_\perp + p_\parallel) = (\frac{\partial}{\partial\psi} - \beta\prime)(B_\chi + p_\perp/B_0^2), \quad (1)$$

$$D_\psi\eta_\psi = -\frac{\partial(B_\chi + p_\perp/B_0^2)}{\partial\phi}, \quad (2)$$

and

$$B_\chi = \frac{\partial\eta_\psi}{\partial\phi} - \frac{\partial\eta_\phi}{\partial\psi}, \quad (3)$$

where $\vec{B}_0 = \vec{\nabla}\phi \times \vec{\nabla}\psi$ is the equilibrium magnetic field, $\beta\prime = 2\frac{\partial P_0}{\partial\psi}/B_0^2$, the perturbed electric field is $\vec{E} = -\partial_t\eta_\phi\vec{\nabla}\phi - \partial_t\eta_\psi\vec{\nabla}\psi$, ϕ is the azimuthal angle, ψ the equilibrium flux function, χ is the coordinate along B_0 ($d\chi = B_0 dl$), l is the arc length on a flux surface, B_χ is the perturbed parallel magnetic field, $p_\perp$ and $p_\parallel$ perturbed perpendicular and parallel pressures, $\kappa_\psi|\vec{\nabla}\psi|^2 = \vec{\kappa}\cdot\vec{\nabla}\psi$ with $\vec{\kappa}$ the curvature of $\vec{B}_0$, P_0 is the equilibrium pressure which is assumed to be isotropic in this analysis, and

$$D_\phi = \frac{1}{B_0}\partial_l \frac{|\vec{\nabla}\phi|^2}{B_0}\partial_l + \frac{\omega^2|\vec{\nabla}\phi|^2}{v_A^2 B_0^2}, \tag{4}$$

$$D_\psi = \frac{1}{B_0}\partial_l \frac{|\vec{\nabla}\psi|^2}{B_0}\partial_l + \frac{\omega^2|\vec{\nabla}\psi|^2}{v_A^2 B_0^2}. \tag{5}$$

In general, $p_\perp$ and $p_\parallel$ can contain contributions from trapped particles, wave-particle interactions and hydrodynamic terms as in [3]. We will include wave-particle interactions in a future and more detailed article, but for now we will assume an isotropic and incompressible MHD fluid, which leads to $p_\perp = p_\parallel = -\eta_\phi \partial_\psi P_0$.

Also we define the following operator:

$$E_\phi = D_\phi + 2\frac{\partial P_0}{\partial \psi}\frac{\kappa_\psi}{B_0^2}. \tag{6}$$

D_ψ as well as E_ϕ act only on the l-dependence of their arguments, and they depend only parametrically on ψ. For each and every ψ, D_ψ and E_ϕ have a complete set of eigenfunctions $\hat{\eta_\phi}(n,l,\psi)$, $\hat{\eta_\psi}(n,l,\psi)$ along l with eigenvalues $\omega_\psi^2(n,\psi)$ and $\omega_\phi^2(n,\psi)$ respectively, corresponding to the n^{th} mode of each operator. In the present work we will consider only the case $n = 1$, which corresponds to the lowest odd mode with respect to the equator, and suppress n from now on. As expected, the solution to system Eqs. (1-3) depends critically on the ω_ϕ^2 and ω_ψ^2 profiles which themselves depend on the equilibrium. The plasma equilibrium used in this paper was constructed in [1], by solving perturbatively the Grad-Shafranov equation with the plasma pressure profile $P(\psi) = P_0(1 + (\frac{\psi-\psi_0}{\epsilon})^{2\alpha})^{-1}$.Thus the self-consistent effect of a pressure gradient is only modeled locally and not globally. The equilibrium pressure profile employed is shown in Fig. 1, and the resulting ω_ϕ^2 and ω_ψ^2 in Fig. 2.

From Fig. 2 it is seen that $\omega_\phi^2(\psi)$ drops monotonically, as one moves away from the planet (i.e. ψ decreases) until some local minimum $\omega_{\min}^2$, then rises until a local maximum $\omega_{\max}^2$ and subsequently drops monotonically again, as radial distance increases, tending to zero. If β_l is sufficiently high

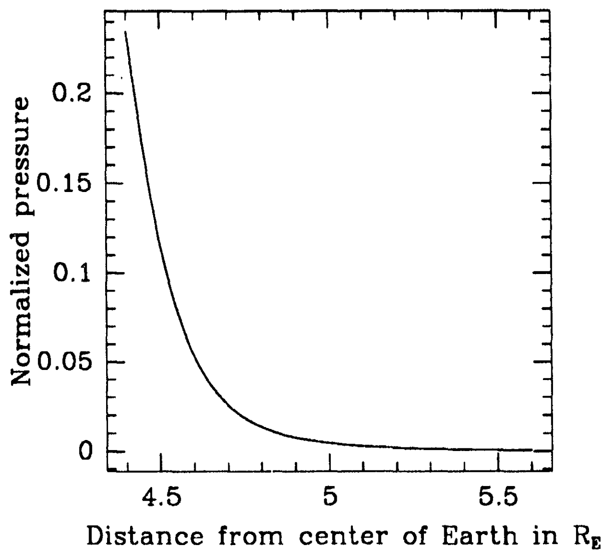

Figure 1: Pressure profile

ω^2_{min} can become negative. $\omega^2_\psi(\psi)$ is always above $\omega^2_\phi(\psi)$ and monotonically dropping with increasing radial distance.

The operators E_ϕ and D_ψ are self-adjoint because of the exclusion of wave-particle interactions, in the sense that $\int_S^N dl B_0 \eta E_\phi \zeta = \int_S^N dl B_0 \zeta E_\phi \eta$, with boundary conditions η, $\zeta = 0$ at the boundary points, and similarly for D_ψ. Equations 1, 2, 3 are not generally tractable in an analytic fashion. We adopt the ordering $|\frac{\partial}{\partial\phi}| = |m| \approx O(10^2) >> 1$, which is consistent with observations, and we will be using $1/m$ as an expansion parameter. Then we can employ the WKB approximation in the radial dimension to reduce the two-dimensional problem to two nested 1D problems.

WKB calculation

We introduce the following WKB *ansatz*:

$$\eta_\phi = \hat{\eta_\phi}(\psi, l) \exp mS(\psi), \tag{7}$$

$$\eta_\psi = \hat{\eta_\psi}(\psi, l) \exp mS(\psi), \tag{8}$$

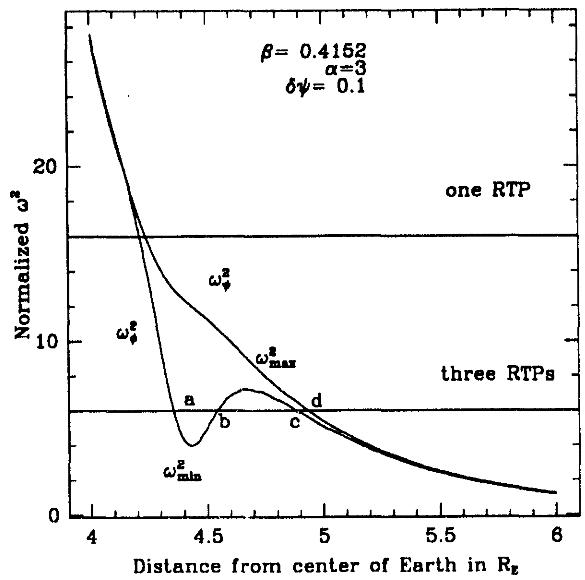

Figure 2: $\omega_\phi^2(\psi)$,$\omega_\psi^2(\psi)$ profiles

$$p_\perp \quad = \quad \hat{p_\perp}(\psi, l) \exp mS(\psi). \tag{9}$$

Define $S_{0\psi} \equiv \partial S_0/\partial\psi$. Introduce the expansion:

$$S \quad = \quad S_0 + \frac{1}{m}S_1 + \frac{1}{m^2}S_2 + \cdots, \tag{10}$$

$$\hat{\eta_\phi} \quad = \quad \hat{\eta_{0\phi}} + \frac{1}{m}\hat{\eta_{1\phi}} + \frac{1}{m^2}\hat{\eta_{2\phi}} + \cdots, \tag{11}$$

$$\hat{\eta_\psi} \quad = \quad \hat{\eta_{0\psi}} + \frac{1}{m}\hat{\eta_{1\psi}} + \frac{1}{m^2}\hat{\eta_{2\psi}} + \cdots. \tag{12}$$

With the WKB ordering that $S_{0\psi}^2 >> \frac{\partial S_{0\psi}}{\partial\psi}$ and $S_\psi\hat{\eta} >> \frac{\partial\hat{\eta}}{\partial\psi}$, we have to lowest order:

$$(S_{0\psi}^2 D_\psi - E_\phi)\hat{\eta_0} = 0. \tag{13}$$

This is equivalent to $S_{0\psi}^2 = \hat{E_\phi}/\hat{D_\psi}$ with the definitions $\hat{E_\phi} = \langle\hat{\eta_0}, E_\phi\hat{\eta_0}\rangle$, $\hat{D_\psi} = \langle\hat{\eta_0}, D_\psi\hat{\eta_0}\rangle$. Here $\langle f_1, f_2\rangle = \int_S^N dl f_1(l) f_2(l) B_0$, that is, we integrate along the field line between its southern and northern boundary points. Equation (13) was derived using the boundary conditions that $\hat{\eta_0} = 0$ at $l = l_N, l_S$, which are the points where the field line intersects the ionosphere. The ionosphere is assummed to be perfectly conducting. $S_{0\psi}$ can be found as the eigenvalue of Eq. (13). In the next order exploiting the self-adjointness of E_ϕ and D_ψ it can be shown that

$$(\partial_\psi S_{0\psi} + 2S_{0\psi}S_{1\psi})\hat{D_\psi} + S_{0\psi}\hat{R} = 0, \tag{14}$$

where:

$$\hat{R} = \langle \hat{\eta_{0\phi}}, (D_\psi(\partial_\psi \hat{\eta_{0\phi}} + \frac{\hat{p_{\perp 0}}}{B_0^2}) + (\partial_\psi - \beta\prime) D_\psi \hat{\eta_{0\phi}})\rangle. \tag{15}$$

Solutions near turning points

The WKB ansatz breaks down close to a ψ where $S_{0\psi}(\psi)$ becomes zero. This ψ is termed a regular turning point (RTP). By inspection of Eq. (13), this happens when $E_\phi(l, \psi_{RTP}; \omega^2)\hat{\eta_0}(l, \psi_{RTP}; \omega^2) = 0$ or equivalently, when $\hat{\eta_0}(l, \psi_{RTP}) = \hat{\eta_\phi}(l, \psi_{RTP})$ and $\omega^2 = \omega_\phi^2(\psi_{RTP})$. Another type of turning point is the singular one, encountered if for some ψ there is no finite eigenvalue $S_{0\psi}$ for Eq. (13). From Eq. (13) this singular turning point (STP), occurs when $\hat{D_\psi} = 0$ or equivalently when $\hat{\eta_0}(l, \psi_{STP}) = \eta_\psi(l, \hat{\psi}_{STP})$ and $\omega^2 = \omega_\psi^2(\psi_{STP})$. The singular turning point corresponds physically to the resonant layer where the wave energy could be dissipated.

To derive the connection formulas near ψ_{RTP} we use a multiscale expansion of the system Eqs. (1-3), which involves a fast and a slow ψ dependence, and an expansion in powers of $m^{-2/3}$. Define $\delta = \psi - \psi_{RTP}$ and $E_\phi^{(0)} = E_\phi(\psi_{RTP})$. Then to lowest order:

$$E_\phi^{(0)} \eta_0 = 0, \tag{16}$$

which is the regular turning point condition. Eq. (16) indicates that η_0 is separable in l and δ, namely, $\eta_0 = \hat{\eta_\phi}(l, \psi_{RTP})h(\delta)$. Using the self adjointness of $E_\phi^{(0)}$ we get from the next order:

$$\frac{\partial^2 h}{\partial \delta^2} = h\delta \frac{\langle \hat{\eta_\phi}, E_\phi^{(1)} \hat{\eta_\phi}\rangle}{\langle \hat{\eta_\phi}, D_\psi^{(0)} \hat{\eta_\phi}\rangle}, \tag{17}$$

where $E_\phi^{(1)}$, and $D_\psi^{(0)}$ correspond to $\frac{\partial E_\phi}{\partial \psi}$, and D_ψ respectively, evaluated at ψ_{RTP}. Eq. (17) is exactly the Airy equation for the ψ dependence of the global mode near ψ_{RTP}. The WKB connection formulae can be derived by asymptotically matching the solutions of Eq. (17) to the WKB solutions given in the preceding section. Thus we have shown that by employing the WKB method a two dimensional eigenvalue problem can be reduced to two nested one dimensional problems. We could carry out a similar analysis near the STP but this is not necessary as will be discussed later on.

The Dispersion Relation

By referring to Fig. 2 and remembering that $S_{0\psi} = 0$ when $\omega^2 = \omega_\phi^2$, we can identify the possibilities that exist for the regular turning points, and similarly for the singular turning point. In the figure, we see that depending on ω^2 we have the following cases:

1. $\omega^2 > \omega_{\max}^2$ gives one regular turning point,

2. $\omega_{\max}^2 > \omega^2 > \max(0, \omega_{\min}^2)$ gives three regular turning points,

3. For large $\beta\prime$ we have the possibility $\omega_{\min}^2 < \omega^2 < 0$ (i.e. local balooning unstable mode) and we obtain two RTP's,

4. $0 < \omega^2 < \omega_{\min}^2$ we get one RTP.

There is always a singular turning point at $\omega^2 = \omega_\psi^2(\psi_{STP})$.

We recall from Eq. (16) ($S_{0\psi}^2 = \frac{\hat{E_\phi}}{\hat{D_\psi}}$) that for $\omega^2 \neq \omega_\phi^2(\psi)$, $\hat{E_\phi}(\psi)$ will be nonzero. The same can be said for ω_ψ^2 and $\hat{D_\psi}$. Writing out explicitly $\hat{E_\phi}$ we get :

$$\hat{E_\phi} = \int_S^N dl \frac{1}{B_0^2} \left((\kappa_\psi \beta\prime + \frac{\omega^2 |\vec{\nabla}\phi|^2}{v_A^2}) \hat{\eta_0}^2 - |\partial_l \hat{\eta_0}|^2 |\vec{\nabla}\phi|^2 \right) . \tag{18}$$

The above equation is the variational problem from which $E_\phi \hat{\eta_\phi} = 0$ can be derived, and thus it is clear that $\hat{E_\phi}(\psi) < 0$ for $\omega^2 < \omega_\phi^2(\psi)$. Similarly $\hat{D_\psi} < 0$ for $\omega^2 < \omega_\psi^2$. Assuming that $\frac{\partial \omega_\phi^2}{\partial \psi} \neq 0$ at $\psi = \psi_{RTP}$ and noting that $\omega_\phi^2 < \omega_\psi^2$ we can conclude that:

1. $S_{0\psi}^2(\psi) > 0$ for $\omega^2 < \omega_\phi^2(\psi) < \omega_\psi^2(\psi)$,

2. $S_{0\psi}^2(\psi) < 0$ for $\omega_\phi^2(\psi) < \omega^2 < \omega_\psi^2(\psi)$,

3. $S_{0\psi}^2(\psi) > 0$ for $\omega_\phi^2(\psi) < \omega_\psi^2(\psi) < \omega^2$.

To have possible radial localization we need to have negative $S_{0\psi}^2$ between two RTP's, and according to the above arguments this can happen only if $\max(\omega_{\rm min}^2, 0) < \omega^2 < \omega_{\rm max}^2$. The case of $\omega_{\rm min}^2 < \omega^2 < 0$ could also fulfill this condition, but it corresponds to ideal MHD instabilities, which require higher drive and is beyond the scope of the present analysis. Thus, we have to look for the global eigenvalue ω^2 between these two limits. It is possible, in principle, that the ω_ϕ^2-well could become so shallow that there will be only one regular turning point in all cases and thus no mode could be localized there. That could happen if $\beta\prime$ is below a certain threshold.

To completely specify the problem, we need to impose boundary conditions. Referring to the previous results and Fig. 2, we see that, to the left of the RTP a the WKB solution will be a sum of an exponentially growing and an exponentially decaying part. The appropriate boundary condition in this region is that the solution be spatially decaying. The other boundary condition is supplied by the existence of the resonant layer (STP) at d. In the MHD description the mode energy will be completely absorbed there, and so we know that there cannot be a reflected wave in the region between c and d. This corresponds to an outgoing wave boundary condition.

Proceeding as in the usual one-dimensional case, we connect the various WKB solutions through the turning points and impose the relevant boundary conditions. We then obtain the following WKB eigenvalue condition or dispersion relation:

$$\sin\left(mI_{ab}(\omega) + \frac{\pi}{2}\right) + \frac{i}{4}\exp(-2mT(\omega))\cos\left(mI_{ab}(\omega) + \frac{\pi}{2}\right) = 0, \tag{19}$$

where $I_{ab} = \int_a^b d\psi\sqrt{-S_{0\psi}^2}$ and $T = \int_b^c d\psi S_{0\psi}$, with a,b being the two RTP's defining the localization area, and b,c defining the tunnelling region. Note that, assumming total wave absorption, the exact location of the resonant layer does not enter in our results.

Assuming that $\omega = \omega_r + i\gamma$ and that $\gamma \ll \omega_r$, we get

$$I_{ab}(\omega_r) = \frac{(r+\frac{1}{2})}{m}\pi, \tag{20}$$

$$\gamma = -\frac{\exp(-2mT(\omega_r))}{4m\dfrac{\partial I_{ab}(\omega_r)}{\partial \omega_r}} \tag{21}$$

In Eq. (20) r is the radial 'quantum' number, the number of radial nodes of the global mode between the turning points a and b. Equation (21) shows that the damping rate is controlled by the exponential tunnelling factor. That means that the least damped modes are those that have large T, or, equivalently, those localized close to the bottom of the well. However we can see from the quantization condition that the larger m is, the smaller I_{ab} has to be, so that we have to have the corresponding eigenfrequencies closer to $\omega^2_{\rm min}$. At the same time $\exp(-2mT)$ becomes very small too, justifying the assumption $\gamma \ll \omega_r$ for $m \gg 1$.

Physically, higher m modes have to tunnel their way for a longer distance, and thus by the time they couple to the field line resonance, there is very little energy to be dissipated. Note that in the present work, we have assummed first harmonic antisymmetric modes; i.e. modes with only one node. For higher longitudinal mode numbers, we would simply have to find a higher eigenmode of Eq. (13) with a corresponding $S^2_{0\psi}$ eigenvalue.

We can predict which is the lowest m number a mode can have and still be expected to be radially localized and not leak out readily towards the shear Alfvén continuous spectrum. We can calculate $I_{ab}(\omega_r = \omega_{max})$ and use the fact that $(r+\frac{1}{2})/m < I_{ab}(\omega_{max})$ to get the lower limit on m. An approximate upper limit for the observable m can be found by looking at the growth rates due to the resonant particle interactions, which show a clear peak [1]. Figure 3 shows ω and γ for various m numbers, assuming one radial node. Note that as m becomes smaller, the damping rate increases sharply due to the decreasing tunnelling loss. From Eq. (20) we can also see that higher r numbers require, higher m_{min} modes, which however are not easily excited by wave-particle interactions.

This means that effectively only $r = 0, 1$ (radial nodes) are viable candidates. Also modes with higher longitudinal mode numbers (n), are less likely to be localized. The reason for this is that, for high n the effect of the ballooning term in E_ϕ is very diminished and the well will become shallow or disappear altogether. Interestingly enough the lowest n is favored by the excitation mechanism too.

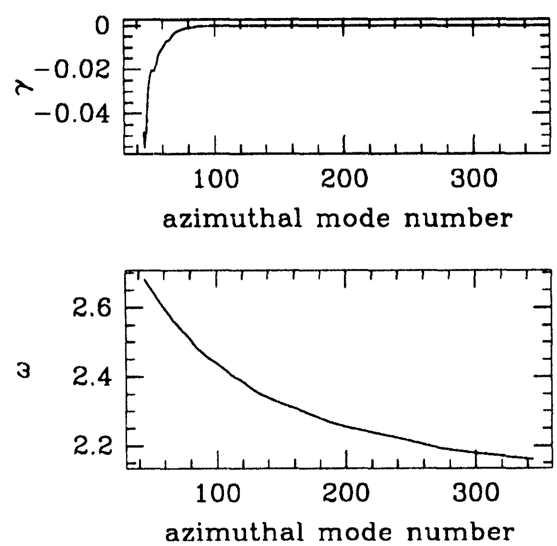

Figure 3: ω and γ versus azimuthal mode number

Summary and Discussion

We have shown that the Drift Alfvén Ballooning modes, previously studied locally along the field line, can be radially localized in the presence of a steep enough pressure gradient which almost vanishes after a point. The radial extent of localization is defined essentially by the locations of the maximum and the minimums of the pressure gradient as a function of radial distance. These modes, however, suffer finite damping through their coupling to field line resonances. It is important to note that the damping is negligible for $m \approx 100$ and longitudinal "quantum number" $n = 1$, (odd symmetry with respect to the equator). These numbers correspond to modes with peak growth rates, as well as to those observed by satellites. The present results thus further support the DABM as a viable instability candidate. Observing the radial structure of Pc4 and Pc5 waves in the terrestial magnetosphere has always been difficult given that most satellites are in geosynchronous orbits. The predictions of this work thus, might not be easily testable even though they are consistent with satellite observations. Finally, we note that the present approach can be readily extended to include effects of wave-particle interactions as well as pressure anisotropy. The results will be reported in a more detailed future publication.

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