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GLOBAL ALFVÉN EIGENMODES IN DIII–D*

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1. INTRODUCTION

Global Alfvén modes, such as the Toroidicity-Induced Alfvén Eigenmode (TAE), pose a serious threat for strongly-heated tokamaks since they can result in saturation of the achievable beam β at moderate levels and they may also cause serious α-particle losses in future ignited devices. The DIII–D tokamak has a unique capability for study of the resonant excitation of these instabilities by energetic beam ions. TAE modes have now been observed in DIII–D over a wide range of operating conditions, including both circular cross-section and elongated ($\kappa \sim 1.8$) discharges. Equilibrium reconstructions of several representative discharges, using all available external magnetic and internal profile data, have been done and analyzed in detail. The computed real mode frequencies of the TAE modes are in good agreement with the experimentally observed mode frequencies and differ significantly from the estimated kinetic ballooning mode frequencies. The TAE calculations include coupling to the Alfvén and acoustic continuum branches of the MHD spectrum and generally indicate that the simplified circular cross-section, large aspect-ratio assumptions made in analytic calculations are poor approximations to the actual TAE mode structures. In particular, the global TAE modes are almost always coupled to one or more continuum branches by toroidicity, poloidal shaping, and finite β effects. Estimates of the various resonant excitation and damping mechanisms, including continuum damping, have been made and the total is found to be in reasonable agreement with the experimental threshold.

2. COMPARISON: THEORETICAL AND EXPERIMENTAL FREQUENCIES

In low field, beam-heated DIII–D discharges, MHD modes are observed which rotate with respect to the plasma at a real finite frequency. Typically, modes with toroidal mode numbers n in the range $3 \lesssim n \lesssim 10$ are observed. The real frequency, when corrected for Doppler shift due to the plasma rotation, is in approximate agreement with the frequency of TAE modes as predicted by the simple formula [1] $\omega \sim V_A/2qR$, where V_A is the Alfvén speed, R is the major radius, and q is the value of the safety factor at which the modes peak; typically $q \sim 1.5$ is taken. The results from a typical sequence of similar shots are shown, for example, in Fig. 1 (solid points). Here, the Alfvén speed was varied by systematically changing the toroidal field, keeping the total current set at 0.6 MA, density 3×10^{19} m^{-3} and beam power constant at 5 MW. The straight line corresponds to $\omega = V_A/2qR$ with $q = 1.5$. The diamonds are the predictions for kinetically destabilized ballooning modes [2], which are also predicted to have a real nonzero frequency in the plasma frame. We note in passing that these are not in agreement with the observed frequency dependence, thus eliminating those modes as being responsible for the observations.

The equilibria were reconstructed using all available profile and magnetic data, as described in Ref. 3, for shots 71515 and 71524. They were then analyzed using the two ideal MHD stability codes CONT [4] and GATO [5]. The CONT code computes the singular eigenvalues of the ideal MHD equations, providing both the real frequencies and spatial locations of the stable Alfvén

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and acoustic continuum modes. The result is a global picture of the MHD continua, including the continuum gaps induced by both toroidicity [4] and shaping [4,6]. The GATO code has been modified to compute the real part of the eigenvalues and eigenvectors of the ideally stable modes. The continuum is discretized and the logarithmic singularities of the continuum modes are truncated by the finite mesh. Global modes, however, are well represented and appear in gaps within the appropriate continua. The gap extrema predicted by the CONT and GATO codes are in extremely good agreement.

The spectrum predicted from the CONT code for shot 71515 is shown in Fig. 2 for toroidal mode number $n = 3$. The solid curves are the continuum branches (labeled with the dominant poloidal mode number for $m \leqslant 9$), plotted as the normalized square of the frequency against the radial position at which the mode is resonant. The dotted curves are the envelopes of the gaps for all n. The experimentally-measured density profile was used in both the CONT and GATO calculations. For clarity, acoustic continua are not shown in Fig. 2. The lowest gap near $\omega^2 \gtrsim 0$ is due to finite β [4]. The next gap is the toroidicity-induced gap and the third gap is created by the coupling of poloidal harmonics m and $m+2$ ($m = 3, 4, \ldots$) induced mostly by the elongation of the discharge [6] ($\kappa \sim 1.7$). The GATO calculations, which in this case, included full coupling between the Alfvén and acoustic branches (*i.e.*, compressibility included) as well as the presence of a finite vacuum and the correct DIII–D wall location, revealed two groups of closely-spaced TAE-like modes with frequencies lying within the TAE gap near $\omega^2 = 0.101$ and 0.133. Within each group, the TAE-like eigenmodes comprise a global component, coupled to singular Alfvén and acoustic continuum harmonics. The Fourier analysis of the perpendicular displacement is shown in Fig. 3 (a) and (b) for two eigenmodes typical of the respective groups. For the lower frequency mode at $\omega^2 = 0.101$ in Fig. 3 (a), the $m = 4, 5$ TAE-like structure is coupled to several higher $m, m+1$ TAE-like components and a singular $m = 16, 17$ continuum component near the edge. There is also a small $m = 3$ continuum component near the center, as expected from Fig. 2. On the other hand, the higher frequency mode ($\omega^2 = 0.133$) in Fig. 3 (b) is dominated by the $m = 3, 4$ TAE component with smaller higher m TAE structures and several high m ($m \sim 12$) continua inside the edge. The small $m = 3$ singular branch is also present near the center. The multiple modes in each group are a result of coupling to the higher harmonics — both global harmonics from the multiple local gaps ($m, m+1$ for $m > 4$) and continuum harmonics where the TAE gap is intersected by the continua [7] as in Fig. 2.

The frequencies computed for these two modes are also shown as circled crosses in Fig. 1. The agreement with the observed frequency is extremely good. We have also looked for experimental evidence of multiple modes in shot 71515, as predicted by the GATO code. However, TAE activity in this particular shot was relatively weak and, although there is some suggestion from the data that two or more modes are present at slightly different frequencies, the evidence is inconclusive. Nevertheless, more detailed analysis of the similar, but lower field shot 71524, where the TAE activity was much stronger, does reveal the existence of at least two distinct TAE modes separated in frequency by about 20%. These are also indicated in Fig. 1 by the circled pluses. The stability analysis for this shot, however, was limited to the incompressible case, in which acoustic coupling is eliminated. A conducting wall was also imposed on the plasma surface. Incompressibility introduces a downward frequency shift of about 25% and, together with the change in boundary conditions, modifies the coupling between the various poloidal harmonics. Nevertheless, the results do indicate that there are probably at least two modes in this shot, analogous to those shown in Fig. 3. For reference, the eigenvalue of the highest frequency eigenmode from that analysis is shown in Fig. 1 (open circle); the correct frequency is probably about 25% higher than that shown.

3. STABILITY THRESHOLDS

Resonance between energetic particle drift motion and the rotating TAE wave introduces an imaginary part to the otherwise real ideal MHD eigenvalue which, in the absence of dissipation, can result in an exponential growth rate [1] $\gamma_\alpha > 0$; in DIII–D, the energetic particles are the injected neutral beams. The mechanism is essentially an inverse Landau interaction with energy transferred from the fast particles to the wave. The condition for resonance is that the parallel beam velocity $V_\parallel$ be near V_A, with $\gamma_\alpha > 0$ for $V_\parallel > V_\mathrm{A}$. The growth rate γ_α is proportional to the fast-ion beta β_b.

Additionally, however, it has been shown that a secondary resonance exists [8] at $V_\parallel \gtrsim V_\mathrm{A}/3$ as a result of the coupling of poloidal modes m and $m+1$. Several damping mechanisms also compete with the fast-ion drive. For a TAE mode localized near a particular surface, the overall growth rate normalized to the real frequency can be written $\gamma \sim \gamma_\alpha - (\gamma_\mathrm{CD} + \gamma_\mathrm{IL} + \gamma_\mathrm{EL})$, where γ_α is the destabilization from the fast particles. The other three terms represent the various damping contributions. Specifically, γ_CD is from continuum damping, which results from phase mixing of the wave with nearby continuum modes. γ_IL is the Landau damping from resonance between the wave and background ions ($V_{\parallel\,\mathrm{bg}} \lesssim V_\mathrm{A}$), including a secondary resonance when $V_{\parallel\,\mathrm{bg}} \lesssim V_\mathrm{A}/3$ analogous to the fast-ion inverse-Landau driving mechanism. Finally, γ_EL is the dissipation due to energy lost from the wave to the electrons. For passing electrons, this is from collisionless Landau damping and is essentially negligible. Inclusion of finite conductivity and parallel electron dynamics for trapped electrons, on the other hand, has been shown to result in much larger damping rates [9] and this latter contribution is the dominant damping mechanism in our case; the overall growth rate is largely a balance between this and the fast particle drive, though the other two terms do contribute. The condition $\gamma > 0$ for instability then becomes an expression for a threshold β_b, above which the TAE modes are destabilized.

Expressions for the various growth rate terms have only been derived in the specific case of large aspect ratio, circular cross-section, and are meant to be applied at the location of the TAE mode [8–10]; a general theory for realistic cross-sections is still unavailable. Nevertheless, we can apply the local formulae at the $q = 1.5$ location, which is roughly where the GATO calculations predict the modes peak. This naive application of the simplified theories, while not entirely applicable in DIII–D, is still probably valid to within a factor of order one. We find for shot 71524 that the TAE mode should be unstable over a range of toroidal mode numbers $3 \leqslant n \leqslant 9$, which is in complete agreement with the observations. Analysis of the other shots is still in progress.

In conclusion, TAE modes have been observed in DIII–D and the observed frequencies are in extremely good agreement with the results of full MHD calculations. Moreover, multiple TAE modes with the same toroidal mode number have been identified in both the experimental observations and the calculations; these are separated in frequency by about 20% in either case. The threshold beam β_b, as calculated from simplified formulae, is also in good agreement with the observations.

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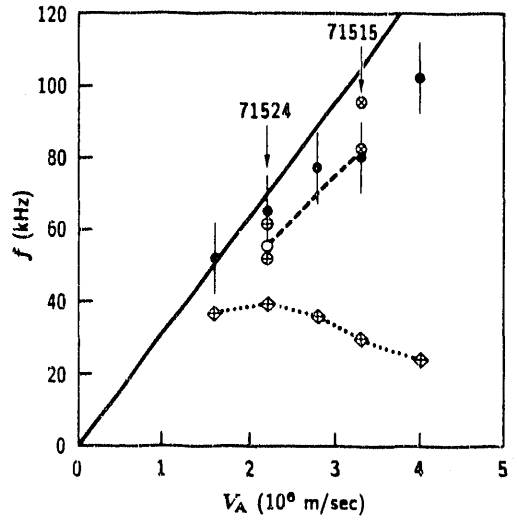

FIG. 1. Frequencies of TAE modes in a sequence of DIII–D shots with varying toroidal field.

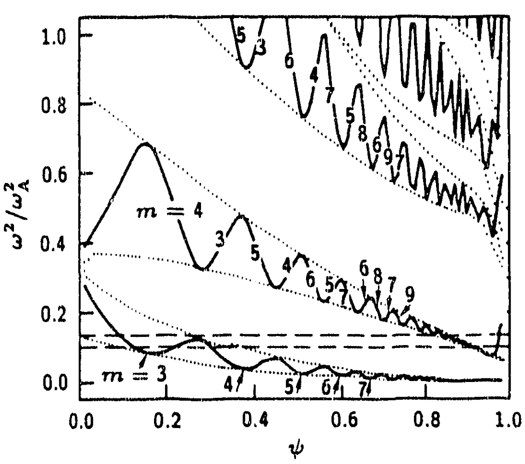

FIG. 2. Continuum shear Alfvén spectrum for shot 71515 with $n = 3$. The TAE mode frequencies computed with GATO are indicated by the broken lines at $\omega^2 = 0.101$ and 0.133.

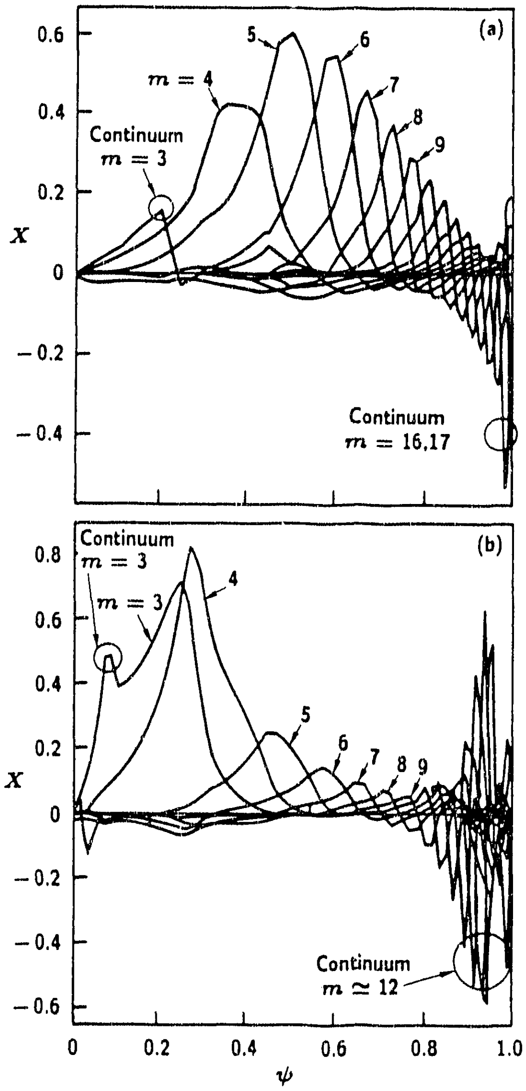

FIG. 3. Fourier decomposition of $X = \xi \cdot \nabla\Psi/|\nabla\Psi|$ for two TAE modes at (a) $\omega^2 = 0.101$ and (b) $\omega^2 = 0.133$.

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