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TOROIDAL ALFVÉN WAVE STABILITY IN IGNITED TOKAMAKS

By

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Abstract

The effects of fusion-product alpha particles on the stability of global-type shear Alfvén waves in an ignited tokamak plasma are investigated in toroidal geometry. Finite toroidicity can lead to stabilization of the global Alfvén eigenmodes, but it induces a new global shear Alfvén eigenmode, which is strongly destabilized via transit resonance with alpha particles.

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1. Introduction

Ignition tokamak experiments, which are being proposed for the near future, will contain alpha particles as fusion reaction products. These particles will have very high energies (e.g., the alpha particles in a D-T reaction are born at 3.5 MeV), and in neither coordinate nor velocity space will they be in thermodynamic equilibrium. Consequently, "thermonuclear" instabilities could be excited in a burning plasma by alpha particles. An example would be instabilities driven by the expansion free energy associated with the spatial gradient of the alpha particle pressure, ∇P_α , which enters the dynamical equations that describe the plasma behavior through the alpha particle diamagnetic frequency ω_{α} . It is expected that in a reacting tokamak plasma, the birth distribution of the alpha particles will be sharply peaked near the magnetic axis: typically, $L_\alpha/a = |a d(\ln P_\alpha)/dr|^{-1} \lesssim 0.4$ for ignition parameters, where L_α is the alpha particle pressure gradient scale length and a is the plasma minor radius.

In the work reported here, we will focus on the destabilization of shear Alfvén waves by alpha particles. For typical ignition parameters the alpha particle velocity $v_\alpha = (E_\alpha/M_\alpha)^{1/2} = 9 \times 10^8 \text{ cm/sec}$ for an energy of 3.5 MeV is comparable to the Alfvén speed $v_A = B/(N_i M_i)^{1/2}$. Thus, the transiting alpha particles could destabilize shear Alfvén waves via inverse Landau damping through the $\omega = k_\parallel v_\alpha$ wave-particle resonance, where $\omega \lesssim \omega_A$ and the shear Alfvén mode frequency is $\omega_A = k_\parallel v_A$. Here, $k_\parallel = (m-nq)/qR$ is the parallel wavenumber for linearized waves that are Fourier decomposed as $\exp[i(m\theta - n\phi - \omega t)]$, with (n,m) and (ϕ,θ) the toroidal/poloidal mode numbers and angles, respectively; also, $q(r)$ is the tokamak safety factor, and R is the major radius.

Briefly, let us describe the various types of shear Alfvén waves in a tokamak plasma. The ideal MHD equation of motion, written in cylindrical geometry (the Hain-Lüst equation), shows that the coefficient of the highest-order radial derivative term, $d^2\xi_r/dr^2$, where ξ_r is the radial fluid displacement, vanishes at radial locations where $\omega^2 = k_\parallel^2 v_A^2$. This corresponds to the shear Alfvén resonance condition, and frequencies ω that satisfy $\min[\omega_A^2(r)] \leq \omega^2 \leq \max[\omega_A^2(r)]$ lie in the shear Alfvén continuum. This resonance leads to a singular mode structure; however, if electron parallel dynamics and ion finite Larmor radius effects are included, one obtains a nonsingular solution known as the kinetic Alfvén wave (KAW). Its mode structure is fairly localized, and hence it experiences strong electron Landau damping.

In this work, we will investigate two global types of shear Alfvén waves that have radially extended mode structure. Both types have low mode numbers n and m . The first type of global shear Alfvén wave is a regular, spatially nonresonant wave whose frequency lies just below the minimum of the continuum, i.e. $\omega < k_\parallel v_A$ and $k_\perp \neq 0$. This wave is called the Global Alfvén Eigenmode (GAE). Previous theoretical analysis of this mode was limited to cylindrical geometry, where it was found that transit wave-particle resonant interaction with super-Alfvénic alpha particles could destabilize it, although with weak growth rates [1]. When finite toroidicity is included, GAE modes with different poloidal mode numbers will become coupled. We will examine such toroidal mode coupling and show that it tends to stabilize the GAE modes completely.

We will also study another type of global shear Alfvén wave, one that exists only in toroidal geometry. Its frequency lies within "gaps" in the shear Alfvén continuum that are created due to toroidal coupling. The

existence of this so-called Toroidicity-Induced Alfvén Eigenmode (TAE) was previously shown in the ideal MHD limit without alpha particles [2]. In the present work, we find that this TAE mode can be strongly destabilized by alpha particles in a burning tokamak plasma.

Section 2 of this paper will discuss the GAE modes and how they are stabilized by finite toroidicity. Section 3 will describe alpha particle excitation of the TAE global modes. A brief theoretical interpretation of the computational results will be given in Sec. 4, along with some concluding comments in Sec. 5.

2. Stability of Global Alfvén Eigenmodes

To begin with, we recall that in the cylindrical limit ($a/R \rightarrow 0$), the GAE mode has been found to be destabilized by inverse Landau resonance with transiting alpha particles [1,3]. The alpha particles are more appropriately modeled with a slowing-down distribution than with a Maxwellian, and at $L_a/a = 0.4$ the growth rate is 10^{-3} of the mode frequency. We also recall that the inverse Landau resonance in the alpha particle kinetic response is $\omega \approx k_{\parallel}^{(m \pm 1)} v_{\parallel \alpha} \gg k_{\parallel} v_{d\alpha}$, since the magnetic curvature drift velocity ($v_{d\alpha} = \cos \theta$) introduces coupling to the $m \pm 1$ sidebands, which then couple back to the main mode m .

In toroidal geometry, we find the new result that finite toroidicity can stabilize the GAE modes via coupling to electron Landau damped sideband modes that resonate with the shear Alfvén continuum [4]. The theoretical framework is contained in the general wave equation, written in the form

$$\nabla \times \nabla \times \underline{E} - \frac{\omega^2}{c^2} \underline{\chi}_F \cdot \underline{E} = \frac{\omega^2}{c^2} (\underline{\chi}_k + \underline{\chi}_\alpha) \cdot \underline{E} \quad (1)$$

Here, $\underline{E}$ is the perturbed electric field, and the susceptibility tensor has been separated into three parts: $\chi_f = (c/V_A)^2 (\hat{I} - \hat{b}\hat{b})$ for the plasma fluid response, χ_k for the core plasma kinetic response (including ion FLR and parallel electron dynamics, but in the low frequency limit $\omega/\omega_{ci} < 1$), and χ_α for the drift-kinetic alpha particle response. It turns out that the dominant finite toroidicity effects arise from the left-hand side of Eq. (1), i.e., the $\nabla \times \nabla \times$ operator and the χ_f tensor, and these were derived for concentric magnetic surfaces in toroidal geometry. The resulting system of coupled differential equations was truncated to a six-by-six system for the m and $m \pm 1$ harmonics of the two transverse electric field components; this truncation is equivalent to an expansion to first order in the inverse aspect ratio, a/R .

Numerical solution of these coupled equations produced the following results. Shown in Fig. 1 is the growth rate of the $n=1$, $m=-2$ GAE mode coupled to the sidebands $(1,-1)$ and $(1,0)$, as a function of the inverse aspect ratio. At $a/R=0$ (i.e., in the cylindrical limit), alpha particles excite this mode to be unstable. However, as toroidicity is introduced, the mode becomes stable - here, as shown in Fig. 1, for aspect ratios $R/a < 10$. Figure 1 was computed for $L_\alpha/a = 0.25$, for which the growth rate of this GAE mode at $a/R=0$ is maximum. The real frequency of the mode is $\omega_r = 7.95 \times 10^6 \text{ s}^{-1}$.

A surprising result is that the $n=0$ GAE mode is found to be stable in toroidal geometry. Previous calculations in cylindrical geometry had indicated that the $n=0$ mode is more unstable in the presence of fusion alpha particles than the modes with $n \neq 0$. In contrast, fully toroidal numerical calculations, performed with a nonvariational kinetic-MHD code to be described in the next section, show that the $n=0$ mode is stable. Briefly stated, the reason for this result is that for $n=0$, the $\pm m$ sidebands have very nearly the same amplitude since the alpha particles are perturbative (i.e., $\beta_\alpha \ll 1$);

hence, the alpha particle excitation that enters through $\omega_{\alpha 0}$, which is proportional to the poloidal mode number m , tends to cancel. This cancellation will be discussed again during the theoretical interpretation given in Sec. 4.

3. Stability of Toroidicity-Induced Alfvén Eigenmodes

Finite toroidicity introduces a new global Alfvén mode, the TAE, which is strongly destabilized by alpha particles. This effect was studied with a model profile in Ref. 4 by treating alpha particle effects perturbatively. Here, we self-consistently investigate the alpha particle destabilization of the TAE mode by employing numerically computed finite β equilibria with noncircular flux surfaces. The existence of the TAE mode was previously demonstrated in the ideal MHD limit [2], in the absence of alpha particles. The high-mode-number version had been explored earlier [5]. This mode exists inside gaps, due to toroidal coupling, in the shear Alfvén continuum spectrum. For example, modes (n, m) and $(n, m+1)$ couple at radial location r_0 , where $q(r_0) = (m + 1/2)/n$, to form a gap. For $n=m=1$, this gap is bounded by [2]

$$\omega_{\pm}^2 = \omega_0^2 + \left\langle \frac{\gamma_1 p B^2 k_s^2}{|v\psi|^2} \right\rangle \pm 2\omega_0^2 \left(\frac{r_0}{R} + \Delta'(r_0) \right), \quad (2)$$

where $\omega_0^2 = (v_A/2qR)^2$, $\Delta(r)$ is the Shafranov shift of the nonconcentric flux surfaces and $\Delta' > 0$, k_s is the geodesic component of the magnetic curvature, and ψ is the poloidal flux.

In order to examine this mode, we adopt an ideal MHD fluid description for the core plasma (labeled c) and a drift kinetic description for the alpha particles, which are assumed to have low beta but be highly energetic:

$\beta_\alpha < \beta_c, T_\alpha \gg T_c$. The alpha particles interact with the core plasma through their perturbed pressure tensor $\delta \underline{P}_\alpha$ in the momentum balance equation for the fluid displacement $\underline{\xi}$:

$$\omega^2 \rho \underline{\xi} = \nabla \delta P_c + \nabla \cdot (\delta \underline{P}_\alpha) + \delta \underline{B} \times \underline{J} + \underline{B} \times \delta \underline{J}, \quad (3)$$

where $\delta \underline{B}$ and $\delta \underline{J}$ are the perturbed magnetic field and current and where $\delta \underline{P}_\alpha$ is derived from the solution for the alpha particle perturbed distribution function, which incorporates both transit and bounce resonance effects. Core plasma kinetic effects on this mode are neglected here.

The resulting integro-differential eigenmode equations were solved in a general toroidal flux coordinate system for realistic numerical equilibria by means of a nonvariational kinetic-MHD code (NOVA-K); the details concerning this code have been described at a previous Varenna workshop [6]. The alpha particle equilibrium distribution function, $F_{\alpha\alpha}$, was taken to be isotropic in pitch angle Λ and slowing down in energy ϵ , as follows:

$$F_{\alpha\alpha} = \begin{cases} c(\psi) \epsilon^{-3/2} & \text{for } \epsilon < \epsilon_\alpha \\ 0 & \text{for } \epsilon > \epsilon_\alpha \end{cases}. \quad (4)$$

Here ϵ_α is the alpha particle birth energy, and $c(\psi)$ is proportional to the alpha particle density, which was taken to be functionally related as $c(\psi) = [P(\psi)]^{7/2}$ to the total plasma pressure $P(\psi)$.

Figure 2 shows the growth rate for the $n=1$ fixed boundary TAE mode as a function of the alpha particle diamagnetic drift frequency $\omega_{\alpha 0}$ (evaluated at $y^{1/2} = 0.5$ and for $m=1$) normalized to $\omega_A = V_A(0)/q(a)R$ for a numerical equilibrium with a circular plasma boundary. The pressure profile is $P(y) = (1 - y^2)^2$, with $y = \psi/\Delta\psi$ the normalized poloidal flux. The other parameters for Fig.

2 are $q(0) = 1.05$, $q(a) = 2.3$, $R/a = 4$, $\langle \delta_{\text{total}} \rangle = 1.89\%$, and $\langle \delta_{\alpha} \rangle = 0.4\%$. The real frequency of the mode is $\omega_r/\omega_A = -0.705$, which lies within the continuum gap formed by the toroidal coupling of the $m=1$ and $m=2$ modes at $q=1.5$. The wave functions for the $m=1$ and $m=2$ modes show that these modes peak near $y^{1/2} = 0.5$; hence $\omega_{* \alpha}$ is a good measure of the alpha particle free energy that can be tapped via inverse Landau damping. Figure 2 shows that $\omega_{* \alpha}$ is large enough to overcome the usual Landau damping when $\omega_{* \alpha}/\omega_A \geq 1.5$ and that beyond this threshold the growth rate γ is approximately linearly proportional to $\omega_{* \alpha}$. From these results we find that for typical ignition parameters, the growth rate of the $n=1$ TAE mode can be of the order of 10^{-2} of the real frequency - which is about one order of magnitude larger than the GAE mode growth rates (evaluated at small a/R).

4. Theoretical Interpretation

In this section we present some interpretations of the numerical results concerning the two types of global shear Alfvén waves. From the eigenmode equations [cf. Eq. (3)], one can construct the quadratic form

$$\omega^2 \delta K = \delta W_f + \delta W_{\alpha} \quad (5)$$

Here, δK is the inertial energy, δW_f is the fluid-like potential energy, and δW_{α} is the potential energy of the kinetic alpha particles (including their transit and bounce resonances). For the mode frequency, we write $\omega = \omega_r + i\gamma$ and assume that the growth rate is small ($|\gamma| \ll |\omega_r|$). Then Eq. (5) yields

$$\gamma \approx -\text{sgn}(\omega_r) \int d\psi d\Lambda \left(\frac{9\pi^2 \hat{r}_t}{4 \delta K} \right) P(\psi) \sum_{m, m', p=-\infty}^{\infty} \frac{[\omega_r - \omega_{* \alpha}^{(m)}] \left(\frac{\omega_r/\omega_t}{p-nq} \right)^2}{(p-nq)^2 \omega_t^2} \text{Re} \{ \langle G_{m,p} \rangle^* \langle G_{m',p} \rangle \}, \quad (6)$$

where $G_{m,p} = F_m(\lambda, \theta) \exp[i[(m-nq)\theta - (p-nq)\omega_t \hat{t}]]$, $F_m(\lambda, \theta)$ depends on the pitch angle, the equilibrium θ -variation and the mode amplitude, $\hat{t}$ is the time-like variable measuring alpha particle position along a field line, ω_t and $\hat{t}$ are the alpha particle transit frequency and transit time, respectively, and $\langle G_{m,p} \rangle$ is the transit average of $G_{m,p}$. The transit harmonic number is p , where p values close to m contribute the most. We define $\omega_{*a}^{(m)} = m\omega_{*a}$, where, as before, ω_{*a} is the diamagnetic frequency for $m=1$.

Equation (6) shows that the instability condition is $\omega_{*a}^{(m)} \geq \omega_r$; here the alpha particle free energy drive overcomes usual Landau damping. This marginal stability condition agrees with the numerical result, shown in Fig. 2, for which $\omega_{*a}/\omega_A = 1.5$ is marginal. Above this threshold, Eq. (6) indicates that the growth rate γ will tend to scale linearly with ω_{*a} , which again agrees with Fig. 2.

The form of Eq. (6) also shows why the $n=0$ GAE mode is stable in toroidal geometry. Since the alpha particle effects are perturbatively small, to a good approximation the mode structure is given by the ideal MHD toroidal wavefunctions. Being self-adjoint, these wavefunctions have the symmetry property $|G_{m,p}|^2 = |G_{-m,-p}|^2$ for $n=0$. Thus, the alpha particle excitation terms $\omega_{*a}^{(m)}$ and $\omega_{*a}^{(-m)}$ are seen to tend to cancel in Eq. (6), so that the $n=0$ GAE mode is stable.

More precisely, we note that this result for the $n=0$ mode holds when the toroidicity is not too small, since we know that in the cylindrical limit, the $n=0$ mode with single m can be destabilized by transiting alpha particles. Moreover, strictly speaking, the symmetry between the $\pm m$ mode amplitudes is broken by the presence of the ω_{*a} terms. A simple zero-dimensional mode coupling argument [7] allows an estimate of how large the toroidicity should be to give the stable result that is observed numerically. In general, this

estimate indicates that the $n=0$ mode could be unstable only if $(\gamma_m/\omega_m) > 0(\langle \epsilon \rangle^{m+1})$, where ω_m and γ_m are the real frequency and the cylindrical growth rate due to the alpha particle contribution for mode $(0,m)$, and $\langle \epsilon \rangle$ measures the effect of inverse aspect ratio. However, this condition is not satisfied for the typical ignition parameters used in the numerical studies of this paper. In particular, for $m=1$ and $m=2$, which are the most unstable modes in the cylindrical limit, we have $\gamma_m/\omega_m \leq 10^{-3}$, whereas $\langle \epsilon \rangle = 0.25$. Therefore we may conclude that the result of the $n=0$ GAE mode being stable in toroidal geometry is indeed valid.

5. Concluding Comments

The results described herein seem to indicate that the global Alfvén eigenmodes (GAE) will not be problematic in ignition tokamaks. Toroidicity tends to stabilize the $n=0$ GAE modes for $a/R > 0.1$, and the $n=0$ GAE mode is stable when the $\pm m$ mode coupling is taken into account.

Therefore, primary attention - especially experimental - should be focused on the toroidicity-induced shear Alfvén eigenmodes (TAE), which can be strongly destabilized by alpha particles. It should be noted that if the edge plasma density value is reduced, the toroidicity-induced mode may possibly resonate with the shear Alfvén continuum near the plasma periphery, an effect which could be stabilizing and should be investigated. Also, we note that an effort is underway to explore the nonlinear saturation of the linearly unstable TAE modes [8].

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Figure Captions

Figure 1. Growth rate γ of the $n=1$, $m=-2$ global Alfvén eigenmode coupled to the sidebands $(1,-1)$ and $(1,0)$, as a function of the inverse aspect ratio a/R .

Figure 2. Growth rate for γ for the $n=1$ toroidicity-induced shear Alfvén eigenmode as a function of the alpha particle diamagnetic drift frequency ω_{α} (normalized to the shear Alfvén frequency ω_A).

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Fig. 1

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Fig. 2

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