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ON RESONANT EXCITATIONS OF HIGH-N MAGNETOHYDRODYNAMIC MODES BY
ENERGETIC/ALPHA PARTICLES IN TOKAMAKS

By

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ON RESONANT EXCITATIONS OF HIGH-N MAGNETOHYDRODYNAMIC MODES BY
ENERGETIC/ALPHA PARTICLES IN TOKAMAKS

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ABSTRACT

Analytical theories for the excitations in tokamaks of magnetohydrodynamic (MHD) modes with large toroidal mode numbers ($n \gg 1$) are presented. Specifically, only instability mechanisms due to resonances with energetic ions/alpha particles are considered. It is noted that, while trapped energetic particles contribute to the ideal region, circulating energetic particles contribute mainly to the singular layer dynamics. A unified dispersion relation manifesting both fishbone-like modes and beam transit-resonance modes is then driven. Finally, we also analyze the stability property of toroidicity-induced shear Alfvén waves excited via transit resonances with alpha particles in ignited tokamaks.

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I. INTRODUCTION

In recent years, observations of MHD oscillations associated with high power ICRF/neutral beam tokamak heating experiments [1-3] have stimulated theoretical investigations on the stability properties of MHD modes as affected by energetic particles [4-6]. Another independent motivation for theoretical research in this area is to understand the α -particle effects on MHD stabilities in future ignited tokamaks [7].

While there has been significant progress in understanding both resonant and nonresonant effects of energetic trapped particles in MHD modes, those due to energetic circulating particles are relatively less known. Both the experimental observations of MHD modes during tangential beam injection [2] and the existence of circulating α particles, however, indicate that stability effects due to energetic circulating particles need to be analyzed in more detail: this is the prime motivation of the present work.

In order to illustrate the underlying physics more clearly, only high- n MHD modes in a relatively simple model equilibrium are considered here. Section II presents the theoretical model and the relevant eigenmode equations. Theoretical analyses are then divided into two classes of modes. Section III deals with low-frequency modes [6,8]; i.e., the mode frequency (ω) is much smaller than the toroidal Alfvén frequency ($\omega_A = v_A/qR_0$; v_A the Alfvén velocity, q the safety factor, and R_0 the major radius). Resonant excitations of toroidicity-induced shear Alfvén waves (TIAW) with $\omega \sim \omega_A$ [9] by circulating α particles with $v_\alpha \sim v_A$ [10] are then considered in Sec. IV. A brief summary and final discussions are given in Sec. V.

II. THEORETICAL MODEL AND EIGENMODE EQUATIONS

We consider a large-aspect ratio ($a/R_0 \ll 1$) tokamak plasma consisting of

a warm core (c) component and an energetic hot (h) component. In terms of the smallness parameter ϵ , we adopt the following formal orderings: $\beta_c \sim \epsilon$, $\beta_h \sim \epsilon^2$, $T_c/T_h \sim \epsilon^2$, $n_h/n_c \sim \epsilon^3$, $k_{\perp c} \sim \epsilon^{1/2}$, and $k_{\perp h} \sim \epsilon^{3/2}$. Furthermore, to make analytical progress possible, we also employ the Connor-Hastie-Taylor analytical model equilibrium of circular magnetic surfaces [11].

Within the present model, the ideal MHD approximation can be shown to remain valid; i.e. $\delta E_{\parallel} = 0$, or

$$i\omega B\delta A_{\parallel}/c = \underline{B} \cdot \nabla \delta\phi; \quad (1)$$

where $\delta\phi$ and $\delta A_{\parallel}$ are, respectively, the scalar and the parallel (to $\underline{B}$) vector potentials. Meanwhile, perpendicular pressure balance gives

$$B\delta B_{\perp} = 4\pi n_i e \omega_{pi} \delta\phi / \omega, \quad (2)$$

where $n_i e \omega_{pi} \equiv (c/B) (\underline{k} \times \underline{e}_{\parallel}) \cdot \nabla P$. The eigenmode equation then corresponds to the following vorticity equation

$$\left[\frac{d}{d\theta} p \frac{d}{d\theta} + \Omega^2 p (1 + 2\epsilon \cos\theta) + \Delta g \right] \delta\phi - \sum_j \eta_j \langle \omega_d \delta G \rangle_j = 0, \quad (3)$$

where we have used the well-known ballooning-mode representation. θ is, thus, the extended coordinate along $\underline{B}$, $\Omega^2 \equiv \omega(\omega - \omega_{pi})/\omega_A^2$, $p(\theta) = 1 + (s\theta - \Delta \sin\theta)^2$, $s = rq'/q$, $\Delta = \beta' q^2 R_0 \sim \epsilon \beta_p$, $g(\theta) = \cos\theta + (s\theta - \Delta \sin\theta) \sin\theta$, $\eta_j = 4\pi \omega e_j / c^2$, $\omega_{dj}(\theta) = \Omega_{dj} g(\theta)$, $\Omega_{dj} = (nq/r) (\sqrt{v_{\perp}^2/2 + v_{\parallel}^2}) / \omega_{cj} R_0$, and $\langle \dots \rangle \equiv \int d^3 \underline{v} (\dots)_j$ with j being the relevant species. δG_j , on the otherhand, is given by the following gyrokinetic equation

$$[\omega_t \frac{d}{d\theta} - i(\omega - \omega_d)]_j \delta G_j = i(\frac{e}{m})_j Q F_{0j} (\omega_d \delta \phi / \omega)_j, \quad (4)$$

where $\omega_t \equiv v_{||}/QR_0$, $Q F_{0j} \equiv (\omega \partial/\partial E + \hat{\omega}_*) F_{0j}$, $E \equiv v^2/2$, and $\hat{\omega}_* F_{0j} \equiv (\mathbf{k} \times \mathbf{e}_{||}/\omega_{cj}) \cdot \nabla F_{0j}$. We note also that, in the present analyses, it is more convenient to let $\delta \hat{\phi} = p^{1/2} \delta \phi$ and rewrite Eq. (3) in the following standard form

$$\left\{ \frac{d^2}{d\theta^2} + \Omega^2(1+2\epsilon \cos\theta) + \frac{\Delta g}{p} - \frac{(s-\Delta \cos\theta)^2}{p^2} \right\} \delta \hat{\phi} - p^{-1/2} \sum_j n_j \langle \omega_d \delta G \rangle_j = 0. \quad (5)$$

Equations (4) and (5) correspond to coupled differential-integral eigenmode equations and are generally difficult to solve analytically. Analytical progress and, hence, physical insights can, however, be obtained if the basic mode structures contain two spatial scales. In general, the two spatial scales do exist; corresponding to, respectively, the inertial $\theta_1 \sim 0$ ($\omega_A/\omega \gg 1$) and the ideal $\theta_0 \sim 0(1)$ regions. For the present purpose, we shall further simplify the analyses by considering the weak-shear $s < 1$ limit and leave the general case for future considerations. That is, in this work, we have $\theta_1 \sim 0$ ($1/s \gg 1$) $\gg \theta_0 \sim 0(1)$. The analytical approach is then to (surface) average Eq. (5) over θ_0 ; the short connection length. The reduced (cylindrical) Eq. (5) in the long scale θ_1 can then be solved via asymptotic matching between the inertial $s\theta_1 \gg 1$ and the ideal $s\theta_1 \sim 0(1)$ regions. Specific analyses differ between the low-frequency ($|\Omega^2| \ll 1$) and TIAW ($|\Omega| \sim 1$) modes and will be presented separately in Secs. III and IV.

III. LOW-FREQUENCY MODES

In this section, we consider MHD modes with $|\omega/\omega_A| \sim |\Omega| < 1$. Since the

modes are resonantly excited, we have, correspondingly, assumed $|\bar{\omega}_{dh}|, |\omega_{th}| < \omega_A$.

Noting that $|\Omega^2|, \Delta, s^2 \ll 1$, Eq. (5) then indicates that $\hat{\delta\phi}$ can be expanded as

$$\hat{\delta\phi} = \hat{\delta\phi}_0(\theta_1) + \hat{\delta\phi}_1(\theta_0, \theta_1) + \dots \quad (6)$$

Substituting Eq. (6) into Eq. (4), we then have for trapped energetic particles

$$\delta G_{h,t} = \left(\frac{e}{m}\right)_h \left\{ \frac{Q F_0}{\bar{\omega}_d - \omega_{h,t}} \right\} \frac{\bar{\omega}_{dh}}{\omega} p_0^{-1/2} \hat{\delta\phi}_0, \quad (7)$$

where $p_0 = (1 + s^2 \theta_1^2)$. In deriving Eq. (7), we have noted that $|\omega| \sim |\bar{\omega}_{dh}| \ll \omega_{bh}$ with ω_{bh} being the bounce frequency. Let us also remark that the effects due to bounce resonances can be shown to be relatively negligible. As to the circulating (untrapped) energetic particles, Eq. (4) yields, noting $|\omega| \sim |\omega_{th}| > |\omega_{dh}|$,

$$\begin{aligned} \delta G_{h,u} = i \left(\frac{e}{m}\right)_h \left\{ -\frac{Q F_0}{\omega^2 - \omega_t^2} \right\} \frac{\omega_{dh}}{\omega} p_0^{-1/2} \hat{\delta\phi}_0 \\ \times [(s\theta_1 \omega_t + i\omega) \cos\theta_0 + (i\omega - s\theta_1) \sin\theta_0] \quad (8) \end{aligned}$$

Substituting Eqs. (7) and (8) into Eq. (5), we then have, after averaging over θ_0 , the following surface-averaged eigenmode equation

$$\left[\frac{d^2}{d\theta_1^2} + \Omega^2 + \frac{u_f}{p_0^2} - \frac{s^2}{p_0^2} - \frac{u_{h,t}}{p_0} - \delta W_{h,u} \right] \hat{\delta\phi}_0 = 0, \quad (9)$$

where $u_f = s\Delta^2$ represents the MHD fluid destabilizing term,

$$u_{h,t} = 4\pi q^2 \frac{m_h}{B^2} \left\langle \left(\frac{v_\perp^2}{2} \right) \frac{QF_0}{\bar{\omega}_d - \omega_{h,t}} \right\rangle \quad (10)$$

which corresponds to the kinetic effects of trapped energetic particles, and

$$\delta W_{h,u} = 2\pi q^2 \frac{m_h}{B^2} \left\langle \frac{(v_\perp^2/2 + v_\parallel^2) \omega Q F_0}{\omega_t^2 - \omega^2} \right\rangle_{h,u} , \quad (11)$$

which corresponds to those of circulating particles.

Equation (9) can be readily solved via asymptotic matching. Noting that in the inertial region $|S\theta_1|^2 \gg 1$, the inertia (Ω^2) and circulating-particle ($\delta W_{h,u}$) terms make the dominant contributions. On the other hand, u_f , shear, and trapped particles contribute mainly in the ideal region $|S\theta_1| \sim 1$. Matching $d\ln\hat{\phi}_0/d\theta_1$ between these two regions, we then derive the following dispersion relation

$$i \left[\frac{\omega(\omega - \omega_{pi})}{\omega_A^2} - \delta W_{h,u} \right]^{1/2} = \delta W_f + \delta W_{h,t} , \quad (12)$$

where

$$\delta W_f = \pi (s - \Delta^2)/4 , \quad (13)$$

$$\delta W_{h,t} = \pi u_{h,t}/2s , \quad (14)$$

and causality constraint requires

$$\text{Re} (\delta W_f + \delta W_{h,t}) < 0 . \quad (15)$$

Equation (12) represents a unified dispersion relation generalizing that of fishbone-like modes [5,6]. To illustrate mode excitations via transit resonances, let us assume no trapped particles and a slowing-down distribution for the untrapped energetic particles, i.e., $F_{oh} = C(r)E^{-3/2}\delta(\lambda - \lambda_0)$ where $\lambda = \mu/E$, $\lambda_0 B = v_{\perp}^2/v_A^2 \ll 1$, $E \leq E_m$, and $C = P_h/(2^{5/2}\pi m_i B_0 E_m)$. Equation (12) then reduces to

$$\omega(\omega - \omega_{pi}) + (\delta W_F \omega_A)^2 + \omega^2 q^2 (P_h / 4n_{oi} E_m) (\omega \omega_{*h} / \omega_{tm}^2 - 1) \times \ln(1 - \omega_{tm}^2 / \omega^2) = 0 \quad (16)$$

Here, $\omega_{tm} = (2E_m)^{1/2} / qR_0$. Taking $|\omega_{*h}| > |\omega_{tm}| > |\omega_{pi}|$, $|\delta W_F \omega_A|$, we have $\omega \approx \omega_{tm} + i\gamma$, and

$$\frac{\gamma}{\omega_{tm}} = \frac{\pi}{8} \frac{q^2 P_h}{n_{oi} E_m} \left(\frac{\omega_{*h}}{\omega_{tm}} - 1 \right)^2 \frac{\omega_{tm}}{\omega_{*h}} \quad (17)$$

Noting the additional constraint $k_{\perp} \rho_h = k_{\theta} \rho_h s_{\theta 1} < 1$, the above analyses are thus valid for

$$|\delta W_F + \delta W_{h,t}|/s \approx |\delta W_F/s| \geq k_{\theta} \rho_h \geq L_{ph}/qR_0 \quad (18)$$

L_{ph} being the scale length of P_h and

$$1 > \sqrt{2E_m}/v_A > |\delta W_F| \quad (19)$$

Note that Eqs. (15) and (19) indicate that instabilities tend to occur fairly close to the marginal stability of ideal MHD ballooning modes.

IV. TOROIDICITY-INDUCED SHEAR ALFVEN WAVES.

In this section, we consider excitations of toroidicity-induced shear Alfvén waves (TIAW) via transit resonances with circulating energetic/α particles in tokamaks such as the proposed CIT ignition experiments. For TIAW, we have $n^2 \approx 1/4$ and, hence, it is suggestive to let

$$\hat{\phi} = A(\theta_1) \cos(\theta_0/2) + B(\theta_1) \sin(\theta_0/2) \quad (20)$$

Equation (5) then yields, for circulation particles,

$$\begin{aligned} \delta G_{j,u} = & -\frac{i}{2} \left(\frac{e}{m} \right)_j \left(\frac{QF_0}{\omega_t - \omega} \right)_j p_0^{-1/2} \sum_{l=1/2, 3/2} (l^2 - \omega^2 / \omega_{tj}^2)^{-1} \\ & \times [(i\omega C_{2j} / \omega_{tj} + 2S_{2j}) \cos l\theta_0 - (lC_{2j} - i\omega S_{2j} / \omega_{tj}) \sin l\theta_0] \quad (21) \end{aligned}$$

where $C_{1/2} = A + Bs\theta_1$, $S_{1/2} = -B + As\theta_1$, $C_{3/2} = A - Bs\theta_1$, and $S_{3/2} = B + As\theta_1$.

Substituting Eqs. (20) and (21) into Eq. (5) and neglecting terms of $d^2A/d\theta_1^2$ and $d^2B/d\theta_1^2$, we then have the following coupled set of equations between A and B:

$$dA/d\theta_1 - (is_2 + 4s\theta_1/2p_0)A = f_- B \quad (22)$$

$$dB/d\theta_1 - (is_2 - 4s\theta_1/2p_0)B = -f_+ A \quad (23)$$

where

$$f_{\pm} = \Gamma_{\pm} - s^2/p_0^2 \pm 4/2p_0 \quad (24)$$

$$\Gamma_{\pm} = (1 \mp \epsilon) \alpha^2 - 1/4 - \beta_1, \quad (25)$$

$$\beta_1 = \pi \frac{\omega q^2}{B^2} \sum_{j=\alpha, e} m_j < (\frac{v_{\perp}^2}{2} + v_{\parallel}^2) \frac{QF_0}{\omega_c^2/4 - \omega^2} >_j, \quad (26)$$

and

$$\beta_2 = \pi \frac{\omega q^2}{B^2} \sum_{j=\alpha, e} m_j \frac{\omega_{tj}}{\omega_c} < (\frac{v_{\perp}^2}{2} + v_{\parallel}^2) \frac{QF_0}{\omega_c^2/4 - \omega^2} >_j. \quad (27)$$

Here, we have kept wave-particle interactions of electrons and α particles; providing, respectively, the dominant stabilizing and destabilizing effects.

Equations (22) and (23) can, again, be solved via asymptotic matching. The resultant dispersion relation is as follows:

$$[-\Gamma_{-}\Gamma_{+}]^{1/2} = -\Gamma_{-}\pi(s-\Delta/s)/4, \quad (28)$$

with the causality constraint

$$\text{Re}[-\Gamma_{-}(s-\Delta/s)] \propto (s-\Delta/s) > 0. \quad (29)$$

We then have, approximately, for $\omega = \omega_r + i\gamma$

$$\omega_r = \omega_A/2, \quad (30)$$

and

$$\frac{\gamma}{\omega_A} = f_{\alpha} k_{\theta} \rho_{\alpha} \beta_{\alpha} \frac{R_0}{L_{p\alpha}} - (\beta_e \frac{m_e}{m_i})^{1/2}. \quad (31)$$

Here, $f_{ra} \sim 0(1)$ is the fraction of resonant α particles. Again, the validity conditions of the above analyses require

$$\frac{v_A}{v_\alpha} \gtrsim k_\theta \rho_\alpha \gtrsim L_{Pa}/qR_0 \quad (32)$$

The corresponding maximum growth rate and critical value of β_α are then given by

$$\frac{\gamma_{\max}}{\omega_A} = \beta_\alpha \frac{v_A}{v_\alpha} \frac{r}{L_{Pa}} \sim 0(10^{-2}) \quad (33)$$

and

$$\beta_{\alpha, \text{crit}} = (v_\alpha/v_A f_{ra})(\beta_e m_e/m_i)^{1/2} \sim 0(10^{-3} - 10^{-2}) \quad (34)$$

Note that $\beta_{\alpha, \text{crit}}$ could readily be exceeded for CIT parameters.

V. SUMMARY AND DISCUSSIONS

In the present work, we have derived analytic dispersion relations for high- n MHD modes resonantly excited by energetic ions/alpha particles in tokamaks. We demonstrate that the dominant contributions of trapped and circulating energetic particles are, respectively, in the ideal and inertial regions. For low-frequency ($\omega \sim \bar{\omega}_{dh}, \omega_{th} \ll \omega_A$) modes, the resultant dispersion relation unifies both fishbone-like modes and modes due to transit resonances. Specific instabilities with $\omega \sim \omega_{th}$ are derived for a model slowing-down untrapped beam ions and are shown to occur near the marginal stability of high- n ideal MHD ballooning modes. In ignited plasmas,

meanwhile, we have $v_g \sim v_A$ and excitations of toroidicity-induced shear Alfvén waves (TIAW) via transit resonances with α particles are analytically demonstrated. Note that high- n TIAW exists as a discrete mode even below the $\epsilon\beta_p$ limit of the ideal MHD ballooning modes and only suffers from weak electron Landau damping.

Several approximations, such as the model equilibrium and weak shear, are made here to simplify the analyses and, thereby, illuminate the physics more clearly. While we can expect straightforward extensions of the present work to general tokamak equilibria, the qualitative features of the stability properties should, however, be similar to those obtained here.

The most crucial extension of the present work, we believe, is to examine the $1/n$ corrections to TIAW. In this respect, TIAW may suffer finite damping due to shear Alfvén resonant absorption induced by finite- n toroidal coupling and the realistic ω_A profile. Details of this analysis will be reported in a future publication.

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