Energy channeling due to energetic-ion-driven instabilities in tokamak

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A kinetic formalism has been developed to study the emission of the destabilized ion Bernstein wave by the energetic α particles in the core region of the tokamak, with very peaked radial distribution. The destabilized waves are driven in the central region of the tokamak but are damped by the electrons in the outer due to the energy channeling. © *2011 American Institute of Physics*. [doi:10.1063/1.3561831]

I. INTRODUCTION

Energetic particles can be created in magnetically confined plasma through auxiliary heating such as neutral beam injection and radio frequency heating and nuclear reactions. The interaction between the energetic ions and the background could be twofold. First they can drive instabilities such as energetic-particle mode,[1] Alfvén eigenmodes,[2–4] etc. Second, microturbulence could affect the confinement of the energetic ions. The confinement of the energetic particles is a critical issue in the International Thermonuclear Experimental Reactor (ITER),[5] since the ignition relies on self-heating by the fusion products. Alpha particle distributions are generally isotropic and contain both kinds of particles, trapped and untrapped. The main difference between trapped and passing ions is the pitch angle, the angle between the particle velocity vector, and the local magnetic field. This quantity is strictly related to the orbital motion and the Larmor radius of a charged particle in a tokamak. It also plays an important role in determining the interaction between energetic ions and plasma waves. Whenever this interaction is significant, the transport of energetic ions can become important. Some recent theoretical[6] and computational[7–10] studies also suggested a significant transport level of the energetic particles driven by the microturbulence. In recent publications,[11,12] the energetic-particle transport of slowing down distribution is found negligible in high energy regime.

Kolesnichenko *et al.*,[13,14] have proposed a new feature of instabilities driven by energetic ions (neutral beam) and found that these instabilities can affect the plasma heating and rotation by channeling the energy and momentum of the energetic ions to the region where the destabilized waves are damped. Due to this energy channeling, the plasma core cannot be heated, as observed experimentally on the stellarator Wendelstein 7-AS (W7-AS) (Ref. 15) and the national spherical torus experiment (NSTX).[16] Experiments on the stellarator W7-AS (Ref. 15) have shown that Alfvén instabilities can considerably reduce the plasma energy and lead to strong thermal crashes (the temperature dropped by up to 30% during instability bursts in the discharge). Experiments on NSTX (Ref. 16) have shown that broadening of the electron temperature correlates with Alfvénic activity and the heat transfer can differ at different locations: the calculated heat conductivity coefficient increases with injected power in the core region and decreases at the periphery. Recently, we[17] explore the effect of magnetic shear on neutral beam driven lower hybrid fluctuations on tokamak, in which the destabilized mode is evanescent in the inner and outer region while propagating waves in the intermediate region.

Valeo and Fisch[18] explored the possibility of excitation of the large $k_\parallel$ ion Bernstein wave (IBW) in tokamak, so as to divert alpha particle power to ions, which can enhance the plasma reactivity. Shalashov *et al.*[19] observed kinetic instabilities of IBW, driven by fast ion distributions, resulting from both radial and tangential neutral beam injections on the W7-AS stellarator. Kumar and Tripathi[20] carry out the local and nonlocal approach to study the excitation of IBW and ion cyclotron wave by considering a gyrating ion beam in plasma column. Brambilla[21] investigated the electron Landau damping of the IBW in plasma, in which the waves are excited by linear mode conversion during fast wave heating in tokamak by using toroidal axisymmetric full wave code TORIC.

In this paper, we study the destabilization of IBW by strongly localized alpha particles in a tokamak when the Bernstein wave leaks into the outer region where it acquires a large parallel wave number due to magnetic shear and is damped on electrons.

We model the tokamak by a plasma slab, with x, y, and z directions corresponding to radial, poloidal, and toroidal directions in the tokamak configuration. We divide the plasma slab radially into two regimes: central and outside region. The magnetic filed profile in the two regimes can be written as follows:

$$\mathbf{B} = \begin{cases} B\hat{z} & \text{for} -a < x < a \\ B\hat{z} \pm B_y\hat{y} & \text{for } \ a < |x| < a_1, \end{cases}$$

as shown in Fig. 1, where B is the static magnetic field (toroidal magnetic field) and, in the outer region, there is an additional small component of magnetic field in the y direction, which appears due to the magnetic shear, and is considered to be constant in magnitude. The equilibrium distribution function of the plasma, which is the mixture of deuterium ions, tritium ions, and electrons, is considered to be Maxwellian. We carry out the instability analysis for the following two regimes.

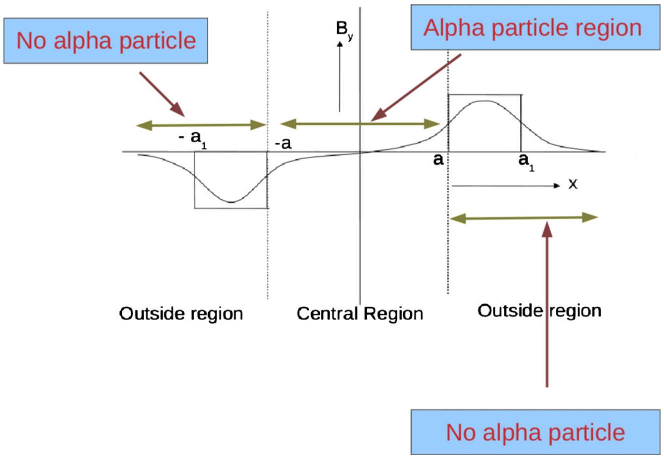

FIG. 1. (Color online) Schematic diagram of magnetic field profile in slab geometry.

A. Central region

Alpha particles are born in the core region of the tokamak. We take their distribution to be isotropic in velocity space but very peaked radial distribution

$$f_{0\alpha} = n_{0\alpha} L_{n\alpha} \delta(x + v_y/\omega_{c\alpha})\, \delta(v - v_\alpha)/v_\alpha^2, \tag{1}$$

where $\omega_{c\alpha}=2eB/m_\alpha$, $2e$, and m_α are the charges and mass of the alpha particle, respectively. We perturb the equilibrium by an electrostatic perturbation $\phi_c=\phi(x)e^{-i(\omega t-\mathbf{k}\cdot\mathbf{r})}$, where $\mathbf{k}=k_\perp\hat{y}+k_z\hat{z}$. The response of alpha particles is governed by the Vlasov equation

$$\frac{\partial}{\partial t} f_\alpha + \mathbf{v}\cdot\nabla f_\alpha - \frac{Z_\alpha e}{m_\alpha}[\nabla\phi_c - \mathbf{v}\times\mathbf{B}]\cdot\nabla_v f_\alpha = 0. \tag{2}$$

Expanding f_α about the equilibrium $f_\alpha=f_{0\alpha}+f_{1\alpha}$ and linearizing the Vlasov equation, one obtains

$$f_{1\alpha} = \frac{2e}{m_\alpha}\int_{-\infty}^{t}\left[\nabla\phi_c\cdot\left(\frac{\partial}{\partial\mathbf{v}} f_{0\alpha}\right)\right]_{t'} dt', \tag{3}$$

where the integration is over the unperturbed trajectory

$$x' = x + \frac{v_\perp}{\omega_{c\alpha}}[\sin\{\omega_{c\alpha}(t'-t)+\theta\} - \sin\theta],$$
$$y' = y - \frac{v_\perp}{\omega_{c\alpha}}[\cos\{\omega_{c\alpha}(t'-t)+\theta\} - \cos\theta]. \tag{4}$$

In the central region of the tokamak with static magnetic field $\mathbf{B}=B\,\hat{z}$, Eq. (3) simplifies to give

$$f_{1\alpha} = -\frac{n_{0\alpha}L_{n\alpha}Z_\alpha e\phi_c}{4\pi m_\alpha v_\alpha^2}\sum_l J_l\left(\frac{k_\perp v_\perp}{\omega_{c\alpha}}\right)\sum_l J_n\left(\frac{k_\perp v_\perp}{\omega_{c\alpha}}\right) \times \frac{e^{i(l-n)\theta}}{\omega - l\omega_{c\alpha} - k_z v_z}\Bigg[k_y\delta(v-v_\alpha)\frac{\partial}{\partial v_y}\delta(x+v_y/\omega_{c\alpha}) + \frac{l\omega_{c\alpha}+k_z v_z}{v}\delta(x+v_y/\omega_{c\alpha})\frac{\partial}{\partial v}\delta(v-v_\alpha)\Bigg]. \tag{5}$$

The perturbed density of alpha particle turns out to be

$$n_\alpha = \int_0^\infty\int_0^{2\pi}\int_{-\infty}^{\infty} f_{1\alpha} v_\perp dv_\perp d\theta dv_z = -\frac{k^2\epsilon_0}{2e}\chi_\alpha\phi_c, \tag{6}$$

$$\chi_\alpha = (i)4\pi\frac{\omega_{p\alpha}^2\omega_{c\alpha}L_{n\alpha}}{k_z v_\alpha^2 k^2}\sum_l\sum_n e^{i(l-n)\sin^{-1}\{-x\omega_{c\alpha}/(v_\alpha^2-\varsigma^2)^{1/2}\}}\Bigg[-v_\alpha k_y J_l(\xi)J_n(\xi)\left\{\frac{(-i)(l-n)}{(v_\alpha^2-x^2\omega_{c\alpha}^2-\varsigma^2)} + \frac{x\omega_{c\alpha}}{(v_\alpha^2-x^2\omega_{c\alpha}^2-\varsigma^2)^{3/2}}\right\} + \frac{k_\perp v_\alpha\omega}{\omega_{c\alpha}(v_\alpha^2-\varsigma^2)^{1/2}}\frac{J_l'(\xi)J_n(\xi)+J_n'(\xi)J_l(\xi)}{(v_\alpha^2-x^2\omega_{c\alpha}^2-\varsigma^2)^{1/2}} + \omega v_\alpha J_l(\xi)J_n(\xi)\left\{\frac{(i)(l-n)x\omega_{c\alpha}}{(v_\alpha^2-\varsigma^2)(v_\alpha^2-x^2\omega_{c\alpha}^2-\varsigma^2)} - \frac{1}{(v_\alpha^2-x^2\omega_{c\alpha}^2-\varsigma^2)^{3/2}}\right\}\Bigg], \tag{7}$$

where $\omega_{p\alpha}^2=4n_{0\alpha}e^2/m_\alpha\epsilon_0$, $\varsigma=(\omega-l\omega_{c\alpha})/k_z$, and $\xi=k_\perp(v_\alpha^2-\varsigma^2)^{1/2}/\omega_{c\alpha}$. One can write the susceptibilities of the plasma electrons and ions due to IBW as

$$\chi_e^c = \frac{2\omega_p^2}{k^2v_{th}^2}\left[1+\frac{\omega}{k_z v_{th}}\sum_l Z\left(\frac{\omega-l\omega_c}{k_z v_{th}}\right)I_l(b)e^{-b}\right],$$
$$\chi_i = \frac{2\omega_{pD}^2}{k^2v_{thD}^2}\left[1-\sum_l\frac{\omega I_l(b_D)e^{-b_D}}{\omega-l\omega_{cD}}\right] + \frac{2\omega_{pT}^2}{k^2v_{thT}^2}\times\left[1-\sum_l\frac{\omega I_l(b_T)e^{-b_T}}{\omega-l\omega_{cT}}\right], \tag{8}$$

in terms of which the electron and ion density perturbations are $n_e=k^2\epsilon_0\chi_e^c\phi_c/e$ and $n_i=-k^2\epsilon_0\chi_i\phi_c/e$, where $k^2=k_\perp^2+k_z^2-\partial^2/\partial x^2$, where the x variation is small; ω_p, ω_{pD}, and ω_{pT} are the electron, deuterium, and tritium plasma frequency, respectively; ω_{cD} and ω_{cT} are the cyclotron frequency of the deuterium and tritium ion, respectively; and $b=(k_\perp\rho)^2/2$, $b_D=(k_\perp\rho_D)^2/2$, $b_T=(k_\perp\rho_T)^2/2$, ρ, ρ_D, and ρ_T are the Larmor radius of the electron, deuterium, and tritium ion, respectively. Using these in Poisson's equation $[\nabla^2\phi_c=(e/\epsilon_0)\times(n_e-n_i-2n_\alpha)\phi_c]$ in the absence of the alpha particle term, one can write

$$\frac{d^2\phi_c}{dx^2} + \frac{S}{E}\phi_c = 0, \tag{9}$$

where

$$S = k_y^2 + k_z^2 + \frac{2\omega_{pT}^2}{v_{thT}^2} \times \left\{ 1 - \alpha_{0T} - \frac{\omega\alpha_{1T}}{\omega - \omega_{cT}} \right\} + \frac{2\omega_{pD}^2}{v_{thD}^2} \times \left\{ 1 - \alpha_{0D} - \frac{\omega\alpha_{1D}}{\omega - \omega_{cD}} \right\} + \frac{2\omega_p^2}{v_{th}^2}\left\{ \frac{k_y^2\rho^2}{2} - \frac{k_z^2 v_{th}^2}{\omega^2} \right\},$$

$$E = \frac{2\omega_{pT}^2}{v_{thT}^2}\left(\beta_{0T} + \frac{\beta_{1T}\omega}{\omega - \omega_{cT}} \right) + \frac{2\omega_{pD}^2}{v_{thD}^2}\left(\beta_{0D} + \frac{\beta_{1D}\omega}{\omega - \omega_{cD}} \right) - \left(\frac{\omega_p^2}{\omega_c^2} + 1 \right), \tag{10}$$

$$\alpha_{j,D,T} = I_j e^{-b_{D,T}}|_{k_x=0} + \frac{\partial}{\partial b_{D,T}}(I_j e^{-b_{D,T}})|_{k_x=0} \frac{k_y^2 \rho_{D,T}^2}{2},$$

$$\beta_{j,D,T} = \frac{\rho_{D,T}^2}{2}\frac{\partial}{\partial b_{D,T}}(I_j e^{-b_{D,T}})|_{k_x=0}, \quad \text{with } j = 0, 1.$$

If the alpha particle term is not zero, we may take ϕ_c to remain largely unmodified; however, the eigenvalue changes

$$\frac{d^2\phi_c}{dx^2} + \lambda^2\phi_c = -\frac{k^2\chi_\alpha}{E}\phi_c. \tag{11}$$

By employing Eq. (9) in Eq. (11) and multiplying the resulting equation by ϕ_c^* and integrating over the entire central region, we obtain

$$\lambda^2 - \frac{S}{E} = -\frac{k^2\int_{-a}^{a}\phi_c^*\chi_\alpha\phi_c dx}{E\int_{-a}^{a}\phi_c^*\phi_c dx} \tag{12}$$

and the solution of Eq. (11) can be written as

$$\phi_c = A_1 \cos(\lambda x). \tag{13}$$

B. Outer region

In this region, there are no alpha particles and the plasma contains electrons and ions only. We consider $\phi_0 = \phi(x)e^{-i(\omega t - \mathbf{k}\cdot\mathbf{r})}$ to be the perturbing potential in this region and one can write down the electron susceptibility as

$$\chi_e^o = \frac{2\omega_p^2}{k^2 v_{th}^2}\left[1 + \frac{\omega}{k_\parallel v_{th}}\sum_l Z\left(\frac{\omega - l\omega_c}{k_\parallel v_{th}} \right) I_l(b) e^{-b} \right], \tag{14}$$

with

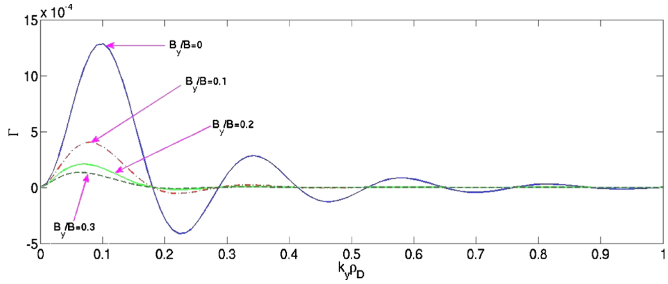

FIG. 2. (Color online) Growth rate of the IBW as a function of the normalized wave number for different values of the magnetic shear parameter $B_y/B = 0, 0.1, 0.2, 0.3$.

$$k_\parallel = \frac{k_z B + k_y B_y}{(B^2 + B_y^2)^{1/2}}. \tag{15}$$

Poisson's equation can be written as

$$\frac{d^2\phi_o}{dx^2} - \frac{F}{R}\phi_o = 0, \tag{16}$$

where

$$F = -(k_y^2 + k_z^2) - \frac{2\omega_{pD}^2}{v_{thD}^2}\left\{ 1 - \alpha_{0D} - \frac{\omega\alpha_{1D}}{\omega - \omega_{cD}} \right\} - \frac{2\omega_{pT}^2}{v_{thT}^2} \times \left\{ 1 - \alpha_{0T} - \frac{\omega\alpha_{1T}}{\omega - \omega_{cT}} \right\} - \frac{2\omega_p^2}{v_{th}^2}\left[\frac{k_y^2\rho^2}{2} - \frac{k_\parallel^2 v_{th}^2}{\omega^2} + i\sqrt{\pi}\left(1 - \frac{k_y^2\rho^2}{2} \right)\frac{\omega}{k_\parallel v_{th}}\exp\left(-\frac{\omega^2}{k_\parallel^2 v_{th}^2} \right) \right], \tag{17}$$

$$R = \frac{\omega_{pD}^2}{v_{thD}^2}\left(\beta_{0D} + \frac{\beta_{1D}\omega}{\omega - \omega_{cD}} \right) + \frac{\omega_{pT}^2}{v_{thT}^2}\left(\beta_{0T} + \frac{\beta_{1T}\omega}{\omega - \omega_{cT}} \right) - \left(1 + \frac{\omega_p^2}{\omega_c^2} \right) + i\sqrt{\pi}\frac{\omega_p^2}{\omega_c^2}\frac{\omega}{k_\parallel v_{th}}\exp\left(-\frac{\omega^2}{k_\parallel^2 v_{th}^2} \right).$$

The solution of Eq. (16) can be written as

$$\phi_0 = A_2 e^{\sqrt{F/R}x} + A_3 e^{-\sqrt{F/R}x}. \tag{18}$$

Employing the continuity of ϕ_c and ϕ_0 and its first derivative at $|x| = a$ and by considering that $\phi_0 = 0$ at $|x| = a_1$, we get the dispersion relation

$$\tan(\lambda a) = \frac{\sqrt{F/R}}{\lambda}\frac{[e^{\sqrt{F/R}(a-2a_1)+e^{-\sqrt{F/R}a}}]}{[-e^{\sqrt{F/R}(a-2a_1)+e^{-\sqrt{F/R}a}}]}. \tag{19}$$

The above transcendental equation can be solved graphically and let m be the solution of the above transcendental equation. The potential of the two regions can be written as

$$\phi = \begin{cases} \phi_c = A_3\dfrac{\sqrt{F/R}}{\lambda\sin(\lambda a)}\cos(\lambda x) \times [e^{\sqrt{F/R}(a-2a_1)} + e^{-\sqrt{F/R}a}] & \text{for } -a < x < a \\ \phi_0 = A_3[-e^{\sqrt{F/R}(x-2a_1)} + e^{-\sqrt{F/R}x}] & \text{for } a < |x| < a_1 \end{cases}$$

and the growth rate of the destabilized wave can be written as

$$\Gamma = -\frac{Im}{\frac{d}{d\omega}[Re]}, \tag{20}$$

where

$$\mathrm{Im} = \{\omega(\omega_{CD}+\omega_{CT}) - \omega_{CD}\omega_{CT}\}2\omega_p^2 k_z^2 \frac{\omega_p^2}{\omega_c^2}\left(\frac{\omega\sqrt{\pi}}{k_\parallel v_{th}}\right)\exp\left(-\frac{\omega^2}{k_\parallel^2 v_{th}^2}\right) - 7.16 G\omega^2\omega_{cD}\omega_{cT}\omega_{p\alpha}^2\omega_{c\alpha}\frac{k_y L_{n\alpha}}{k_z v_\alpha^3}J_0 J_1,$$

$$G = \frac{2\omega_{pD}^2}{v_{thD}^2}\beta_{0D} + \frac{2\omega_{pT}^2}{v_{thT}^2}\beta_{0T} - \left(\frac{\omega_p^2}{\omega_c^2}+1\right),$$

$$Re = \left\{A + 2.38 G\omega_{cD}\omega_{cT}\omega_{p\alpha}^2\omega_{c\alpha}\frac{k_y L_{n\alpha}}{k_z v_\alpha^3}J_0 J_1 - m^2 H\right\}\omega^2 + \{B - m^2 N\}\omega + C - m^2 Q,$$

$$\begin{aligned} A = &\left\{k_y^2\left(1+\frac{\omega_p^2}{\omega_c^2}\right) + k_z^2 + 2\frac{\omega_{pD}^2}{v_{thD}^2}(1-\alpha_{0D}) + 2\frac{\omega_{pT}^2}{v_{thT}^2}(1-\alpha_{0T})\right\}\left\{2\frac{\omega_{pD}^2}{v_{thD}^2}\beta_{0D} + 2\frac{\omega_{pT}^2}{v_{thT}^2}\beta_{0T} - \left(\frac{\omega_p^2}{\omega_c^2}+1\right)\right\}\omega_{cD}\omega_{cT} - 2\omega_p^2 k_z^2 \\ &\times\left\{\frac{2\omega_{pD}^2}{v_{thD}^2}\beta_{0D} + \frac{2\omega_{pT}^2}{v_{thT}^2}\beta_{0T} - \left(\frac{\omega_p^2}{\omega_c^2}+1\right)\right\} - 2\omega_p^2 k_z^2\left\{\frac{2\omega_{pD}^2}{v_{thD}^2}\beta_{1D} + \frac{2\omega_{pT}^2}{v_{thT}^2}\beta_{1T}\right\}, \end{aligned} \tag{21}$$

$$\begin{aligned} H = &-\left\{k_y^2\left(1+\frac{\omega_p^2}{\omega_c^2}\right) + k_z^2 + 2\frac{\omega_{pD}^2}{v_{thD}^2}(1-\alpha_{0D}) + 2\frac{\omega_{pT}^2}{v_{thT}^2}(1-\alpha_{0T})\right\}\left\{2\frac{\omega_{pD}^2}{v_{thD}^2}\beta_{0D} + 2\frac{\omega_{pT}^2}{v_{thT}^2}\beta_{0T} - \left(\frac{\omega_p^2}{\omega_c^2}+1\right)\right\}\omega_{cD}\omega_{cT} \\ &+ 2\omega_p^2 k_\parallel^2\left\{\frac{2\omega_{pD}^2}{v_{thD}^2}\beta_{0D} + \frac{2\omega_{pT}^2}{v_{thT}^2}\beta_{0T} - \left(\frac{\omega_p^2}{\omega_c^2}+1\right)\right\} + 2\omega_p^2 k_\parallel^2\left\{\frac{2\omega_{pD}^2}{v_{thD}^2}\beta_{1D} + \frac{2\omega_{pT}^2}{v_{thT}^2}\beta_{1T}\right\}, \end{aligned}$$

$$B = 2\omega_p^2 k_z^2(\omega_{cD}+\omega_{cT})\left\{\frac{2\omega_{pD}^2}{v_{thD}^2}\beta_{0D} + \frac{2\omega_{pT}^2}{v_{thT}^2}\beta_{0T} - \left(\frac{\omega_p^2}{\omega_c^2}+1\right)\right\} + 2\omega_p^2 k_z^2\left\{\frac{2\omega_{pD}^2}{v_{thD}^2}\beta_{1D}\omega_{cT} + \frac{2\omega_{pT}^2}{v_{thT}^2}\beta_{1T}\omega_{cD}\right\},$$

$$N = -2\omega_p^2 k_\parallel^2(\omega_{cD}+\omega_{cT})\left\{\frac{2\omega_{pD}^2}{v_{thD}^2}\beta_{0D} + \frac{2\omega_{pT}^2}{v_{thT}^2}\beta_{0T} - \left(\frac{\omega_p^2}{\omega_c^2}+1\right)\right\} - 2\omega_p^2 k_\parallel^2\left\{\frac{2\omega_{pD}^2}{v_{thD}^2}\beta_{1D}\omega_{cT} + \frac{2\omega_{pT}^2}{v_{thT}^2}\beta_{1T}\omega_{cD}\right\},$$

$$C = -2\omega_p^2 k_z^2\omega_{cD}\omega_{cT}\left\{2\frac{\omega_{pD}^2}{v_{thD}^2}\beta_{0D} + 2\frac{\omega_{pT}^2}{v_{thT}^2}\beta_{0T} - \left(\frac{\omega_p^2}{\omega_c^2}+1\right)\right\}, Q = 2\omega_p^2 k_\parallel^2\omega_{cD}\omega_{cT}\left\{2\frac{\omega_{pD}^2}{v_{thD}^2}\beta_{0D} + 2\frac{\omega_{pT}^2}{v_{thT}^2}\beta_{0T} - \left(\frac{\omega_p^2}{\omega_c^2}+1\right)\right\}.$$

In order to assess the results, we consider that the width of the inner region is $a=30$ cm, the background electron density $\sim 2\times 10^{20}$ m^{-3}, $T_e=30$ KeV, the magnetic field at the center of the tokamak is $B=5$ T, the temperature of deuterium and tritium is ~ 10 keV, and the ratio of the mixture to be 50:50 is $n_{0\alpha}\approx 6\times 10^{19}$ m^{-3} and $L_{n\alpha}=0.3a$.

In Fig. 2 we have plotted the growth rate of the destabilized IBW wave as a function of the normalized wave number for different magnetic shear factor (B_y/B). From Eq. (15), it is clear that with the increasing of shear factor (B_y/B), $k_\parallel$ also becomes large. Figure 2 shows that the destabilized IBW wave has large amplitude in the core region; in the outer region, due to the large $k_\parallel$, the wave gets damped on electrons and the amplitude of the destabilized IBW become small. So, energetic-ion-driven instabilities can channel the energy of the energetic ions from the birth region of the alpha particle to the region where the destabilized waves are damped. This phenomenon leads to cooling of the plasma core and heating the periphery, as most recently mentioned by Kolesnichenko *et al.*,[13,14] where Alfvén eigenmodes, destabilized by the neutral beam, play the same role as IBW. It follows from our analysis that the energy can be transported by the wave rather than by electron heat transport. Note that this energy channeling due to alpha particles is a new phenomenon since the alpha channeling predicted earlier[22] does not redistribute the energy and momentum along the plasma radius, which leads to heating of ions rather than electrons.

In our calculations, we have considered that the plasma minor radius is ~ 0.6 m (outer boundary of the plasma $\sim a_1$) and the wave (ϕ) vanishes at this point. The Larmor radius (ρ_α) for the alpha particle is ~ 0.1 m (for $B=5$ T). In choosing the inner boundary of the plasma, we consider that the inner boundary is about the half of the minor radius and the alpha particles, which are generated in the fusion process and highly localized in the core region of the tokamak, will move around the field lines and it may come out from the central region of the tokamak due to successive collisions or some other reasons. Fisch and Rax[23] and Kupfer *et al.*,[24] showed strong quasilinear diffusion of alpha particles which tends to displace alpha particles outward by radio frequency waves before they can slow down.

In summary, it is found that the regions of destabilized IBW emission by the alpha particle and the wave absorption by the plasma are different. Because of this, waves can channel energy from one region to another region, leading to cooling the plasma core, which is supplemented by a qualitative analysis.

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