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Control of Alpha Particle Transport by Spatially Inhomogeneous Ion Cyclotron Resonance Heating

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Control of the radial alpha particle transport by using Ion Cyclotron Range of Frequency waves is investigated in a large-aspect-ratio tokamak geometry. It is shown that spatially inhomogeneous ICRF-wave energy with properly selected frequencies and wave numbers can induce fast convective transport of alpha particles at the speed of order $v_{\text{alpha}} \sim (P_{RF}/n_\alpha \epsilon_0)\, \rho_p$, where P_{RF} is the ICRF-wave power density, n_α is the alpha density, ϵ_0 is the alpha birth energy, and ρ_p is the poloidal gyroradius of alpha particles at the birth energy. Application to ITER plasmas is studied and possible antenna designs to control alpha particle flux are discussed.

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I. Introduction

Control of alpha ash buildup is an important issue in obtaining sustained fusion reactions in a magnetically confined fusion reactor. The usual particle transport rate deduced from experimental observations implies that alpha ash may build up to an unacceptable level in a fusion device such as International Thermonuclear Engineering Reactor(ITER). We must be able to enhance selectively the particle transport rate for alphas without causing deterioration of the background plasma confinement in order to increase the probability of sustained fusion reaction.

Reference 1 introduced a scheme to control selectively the energetic ion transport by using a spatially inhomogeneous ICRF wave. In the present work, we will apply this scheme to control alpha particle transport in the central reaction region. Efficient coupling of ICRF power to the alpha particles requires us to use higher harmonic heating for alpha control, since the background deuterium at lower energy will be strongly coupled to the wave at the fundamental harmonic number. In order to yield radial transport in a tokamak geometry, the spatial inhomogeneity must be transformed into velocity space distortion in a preferred parallel direction (net friction) of the influenced particles.[2] ICRF waves will be very effective in delivering such a net friction to the high energy alpha paticles rather than to the low energy deuterons because the large Doppler broadening associated with the high energy particles makes different signs of parallel velocity $v_{||}$ be resonant with the waves at different locations in major radius. Ions with small kinetic energy will exhibit small spatial Doppler spread of the resonance locations and, thus, will all experience essentially the same ICRF-wave power.

New born alpha particles at high energy are valuable in that they provide energy to the background plasma through Coulomb collisional processes. It is required that the present scheme will not greatly enhance this higher energy alpha transport. This can be achieved by considering two techniques. One technique requires adjusting the Bessel function behavior in the ICRF-heating operator in such a way that the alpha birth energy corresponds to a minimum heating point. The other way is to use a small parallel resonance velocity, so that at higher energies the resonance particles exist only in the deeply trapped regime. The latter technique relies upon the fact that the fast bounce motion of trapped particles does not allow for uneven distortion of their distribution function, which is necessary for the present transport process. We find that the latter technique is very reliable in not enhancing the higher energy alpha transport, but some loss of efficiency is inevitable due to the heating of the non-participating high energy alphas in the trapped particle regime. The former technique may allow some enhancement of the high energy alpha transport, but there will not be loss of ICRF power to the non-participating alpha particles.

The other requirement that needs to be satisfied for ITER is that the total ICRF power should be much smaller than the total alpha power. In Section II, energetic ion transport calculation of Ref. 1, induced by a spatially inhomogeneous ICRF-wave power density, has been further developed to treat the alpha particle removal problem in the central burning region of ITER-like plasma. After careful assessment of the requirements, it is found that a proper selection of wave numbers and resonance harmonic numbers can successfully control alpha particle transport between thermal and birth energies. The present scheme can suppress the supply of alpha ash by making the lower-energy tail alphas move out of the central reaction region when the plasma is free of any MHD activity. However, when a mild level of sawtooth activity is present in the central burning region, the radial alpha flux can be sufficiently large inside the $q = 1$ surface. In this case, the present scheme can be applied to the region outside the sawtooth activity region in order to push the alpha particles further towards the edge region. This case is not included in the present analyses. Once the alphas are pushed away from the central reaction region, other loss mechanisms, like ripple loss,[3] can also take place.

In the present work, we emphasize the radially outward loss of the lower energy part of tail alpha particles. It is also possible, however, to induce radially inward transport of high energy ions by changing the direction of $k_{\parallel}$ or the spatial gradient of the wave power density. This scheme may be useful when there is difficulty concentrating the high energy ions within the central reaction region. Some discussions on the present scheme are presented in Section III.

II. Alpha particle flux induced by an inhomogeneous ICRF wave

The quasilinear form of the steady-state drift-kinetic equation used in Ref. 1 is

$$\frac{\partial f}{\partial t} + (v_{\parallel}\hat{n} + \vec{v}_d) \cdot \nabla f = C(f) + Q(f) + S \ , \qquad (1)$$

where f is the quasi-static part of the alpha particle distribution function averaged over the gyrophase, t represents time, $v_{\parallel}$ is the parallel component of the velocity to the magnetic field $\vec{B}$, C is the Coulomb collision operator, Q is the quasilinear RF-heating operator, S represents the alpha-birth function, and the applied Ohmic electric field term has been neglected because its effects are usually small for the ion species.[2]

By assuming that the gyroradius ρ is small compared to the radial scale length L for alpha pressure,

$$\delta \equiv \frac{\rho}{L} \ll 1 \ ,$$

that the alpha collision frequency ν_c is much smaller than the alpha transit frequency ω_T,

$$\hat{\nu}_c = \frac{\nu_c}{\omega_T} \ll 1 \ ,$$

and that the ICRF-heating rate is of first order in both δ and $\hat{\nu}_c$,

$$Q \sim \delta \hat{\nu}_c \, \omega_T \ .$$

Ref. 1 obtained

$$\Gamma_\psi = -\frac{\bar{D}_b}{v_0^2}\frac{Iv_0}{\overline{\Omega}}\mathcal{I} ,$$

$$\mathcal{I} = \sum_\ell \left\langle \int_{pass} d^3s \frac{s_\parallel}{s^2}\frac{\partial}{\partial s}\left[s^2\lambda\,\delta\left(\xi_R - \frac{\omega - \ell\Omega_R}{k_\parallel v}\right)\left|J_{\ell-1}\left(\frac{k_\perp v_{\perp R}}{\Omega_R}\right)\right|^2 \frac{\partial \hat{f}_0}{\partial s}\right]\right\rangle ,$$

where $\overline{D}_b$ and $\overline{\Omega}$ are the flux-surface-averaged values of the bounce-averaged quasilinear ICRF diffusion coefficient D_b and the gyrofrequency Ω, $\vec{s} = \vec{v}/v_0$ is the normalized velocity variable, λ is the poloidally varying part of D_b, and $\hat{f}_0$ is the lowest order alpha particle distribution function expressed in the normalized velocity $\vec{s}$. The velocity integral $\int_{pass} d^3s$ in $\mathcal{I}$ is evaluated over the passing regime only. This is because of the fact that asymmetric variation of the distribution function in $v_\parallel$ is necessary to produce a net radial transport, but the fast bounce motion of the trapped particles does not allow for any variation of the collisionless distribution function along the particle trajectory. We refer the readers to Ref. 1 for detailed explanation.

The normalized lowest order distribution function $\hat{f}_0$ satisfies

$$\int d^3s\,\hat{f}_0 = n_\alpha .$$

Using

$$\mathcal{S} = \frac{\beta_\alpha}{4\pi v^2}\delta(v - v_0)$$

for the alpha source function, where β_α is the alpha birth rate and v_0 is the alpha birth speed, and using the slowing-down approximation for the Coulomb operator, we can easily show[4] that $\hat{f}_0$ may be approximated by

$$f_0^{(0)} = \frac{\beta_\alpha \tau_s}{4\pi(s_c^3 + s^3)} \quad \text{for } v \le v_0,$$

and $f_0 = 0$ for $v > v_0$, where τ_s is the alpha-electron slowing-down time and s_c is the alpha-electron critical slowing-down speed normalized to v_o.

The bounce-averaged RF-heating diffusion coefficient D_b for $k_\parallel v_\parallel \ll \ell\Omega$ can be found from the following form of bounce-averaged RF-heating operator, as found in

Refs. 5 and 6.

$$\{Q\}(f_0) \simeq -\sum_{\ell} \frac{1}{v^2} \frac{\partial}{\partial v} \left[v^2 D_b \delta \left(\xi_R - \frac{\omega - \ell\Omega_R}{k_{\|} v} \right) \left| J_{\ell-1} \left(\frac{k_\perp v_\perp}{\Omega_R} \right) \right|^2 \frac{\partial f_0}{\partial v} \right] , \tag{2}$$

where $J_{\ell-1}$ is the Bessel function of order $\ell - 1$, $\xi_R = v_{\|R}/v$ is the cosine of the pitch angle at the resonance location, and $v_{\|R}(\Omega_R)$ is the value of $v_{\|}(\Omega)$ at the resonance location. D_b can be written as

$$D_b = \frac{\pi Z e^2}{8M^2} \frac{|E_+|^2}{|\dot{\Omega}|\tau_b} \bigg|_{\Omega = \ell\omega - k_{\|} v_{\|}} ,$$

where Z is the alpha charge number, E_+ is the left-handed circularly polarized component of the RF-electric field, M is the alpha mass, $\dot{\Omega} = d\Omega/dt$ along the particle parallel motion, and τ_b is the bounce frequency. Here, $\dot{\Omega}$ can be a function of pitch angle for the trapped and the trapped-passing boundary particles. Since only the passing particles contribute to the lowest order radial flux in the present problem, D_b will be approximated as a velocity-independent quantity. In the present work, we assume that the spatial inhomogeneity of D_b is given in the form of $\cos\theta$.

$$\lambda = \lambda_o \cos\theta ,$$

$$D_b = \overline{D}_b (1 + \lambda_o \cos\theta) .$$

This form of D_b can be obtained with $E_+|^2$ given in the form

$$|E_+|^2 = \overline{|E_+|^2} (1 + \lambda_o \cos\theta) .$$

The integral $\mathcal{I}$ can be evaluated numerically as follows: For each set of $(s_{\|}, s_\perp)$, the resonance angles θ_R can be found, using the resonance condition $\omega - \ell\Omega_R - k_{\|} v_{\|R} = 0$ and the conservation of energy and magnetic moment. Thus, the quantities Ω_R,

$v_{\perp R}$, ξ_R, and θ_R are found as functions of $(s_{\|}, s_{\perp})$ and the velocity integral $\mathcal{I}$ can be numerically evaluated accordingly using a toroidal field model. The δ-function in the integrand will ensure that only the resonant particles contribute, and at the same time it includes any multiple resonance locations, which can easily exist for higher ℓ-numbers in a large-size tokamak.

For the alpha particles in the central reaction region, we can use the large aspect ratio approximation and simplify the integration further. Now, the differences between $(v_{\perp R}, \Omega_R, \xi_R)$ and $(v_{\perp}, \Omega, \xi)$ are small for the passing particles, and we can set $v_{\perp R} = v_{\perp}$, $\Omega_R = \bar{\Omega}$ and $\xi_R = \xi$ to simplify the evaluation of the integral $\mathcal{I}$, valid to the lowest order in the inverse aspect ratio. Since the trapped particles do not participate in the present transport mechanism, by having small $v_{\|R}$ we can be sure that the highest energy newborn alpha particles at $v \sim v_0$ are not transported radially outward (by having the resonant alphas in the trapped regime at high energy). We thus require $v_{\|R}/v_0 \ll \sqrt{r/R}$. With the simplifications given above and after integration by parts in s, in which we ignore the end point term as a consequence of our additional requirement, we obtain

$$\mathcal{I} = -\frac{3\beta_\alpha \tau_s}{4\pi} \sum_{\ell} \left\langle \int_{pass} d^3 s \frac{s s_{\|} \Theta(1-s)}{(s_c^3 + s^3)^2} \lambda_o \cos\theta_R \, \delta\left(\xi - \frac{\omega - \ell\Omega_R}{k_{\|} v}\right) \left| J_{\ell-1}\left(\frac{k_\perp v_\perp}{\bar{\Omega}}\right)\right|^2 \right\rangle ,$$

where $\Theta(x)$ is the heavy-side step function, and the expression

$$\frac{\partial \hat{f}_0}{\partial s} = -\frac{\beta_\alpha \tau_s}{4\pi} \frac{3s^3}{(s_c^3 + s^3)^2}$$

is used.

The value of $\cos\theta_R$ is a function of s and $s_{\|}$, and their mutual relationship can be found from the resonance condition

$$\ell(\bar{\Omega} - \Omega_R) = k_\parallel \bar{v}_\parallel \tag{3}$$

for RF tuned at $\omega = \ell\bar{\Omega}$. Since $\bar{\Omega} - \Omega_R = \bar{\Omega}(r/R)\cos\theta_R$, Eq. (3) reduces to

$$\cos\theta_R = \frac{k_\parallel v_\parallel}{\ell\bar{\Omega}}\,\frac{R}{r} \tag{4}$$

for the particles with

$$\frac{|k_\parallel v_\parallel|}{\ell\bar{\Omega}} \le r/R. \tag{5}$$

If $|k_\parallel v_\parallel| > \ell\bar{\Omega}r/R$, the particle is not in resonance on the flux surface and does not contribute to transport. It is Eq. (5) which allowed us to neglect the $v_\parallel$-heating effect in the operator Q earlier in this section. Using Eq. (5) and the passing condition $v_\parallel/v > \sqrt{r/R}$, we can easily see that all the participating particles satisfy

$$\frac{v}{v_0} < \sqrt{\frac{r}{R}}\frac{\ell\bar{\Omega}}{k_\parallel v_0} < 1\,, \tag{6}$$

where the second equality is required for the neglect of the end points in the previous integration-by-part process. The participating region in velocity space can be seen from Fig. 1. By requiring $\ell\Omega/k_\parallel v_0 \sim 1$, it is clear that transport is enhanced for the lower energy part of tail alpha particles and not for the higher energy part.

Using Eq. (4), it is straightforward to obtain

$$\mathcal{I} \cong -\frac{3\beta_\alpha\tau_s}{4\pi}\lambda_o\frac{\bar{R}}{r}\sum_\ell \frac{k_\parallel v_0}{\ell\bar{\Omega}}\int_{pass} d^3s\,\frac{s s_\parallel^2 \Theta(1-s)\Theta(s_{\parallel max} - s_\parallel)}{(s_c^3 + s^3)^2}\left| J_{\ell-1}\left(\frac{k_\perp v_0}{\bar{\Omega}}s_\perp\right)\right|^2 \,,$$

where $s_{\parallel max} = r/(RA)$, $A = |k_\parallel v_0|/\ell\Omega$, and the flux-surface-average operator is not shown because the quantity to be averaged is a constant to the lowest order in the inverse aspect ratio. With this expression for $\mathcal{I}$, the particle flux Γ_ψ may be changed into

$$\Gamma_\psi = \frac{\bar{D}_b}{v_0^2}\,\frac{I v_0}{\bar{\Omega}}\beta_\alpha\tau_s\,\lambda_o\frac{\bar{R}}{r}\sum_\ell \frac{k_\parallel v_0}{\ell\bar{\Omega}}\hat{\mathcal{I}}_\ell \,,$$

where

$$\hat{\mathcal{I}}_\ell = \frac{3}{4\pi}\int_{pass} d^3s \frac{s s_\parallel^2 \Theta(1-s)\Theta(s_{\parallel max} - s_\parallel)}{(s_c^3+s^3)^2}\left|J_{\ell-1}\left(\frac{k_\perp v_0}{\bar{\Omega}} s_\perp\right)\right|^2 .$$

The quantity $\bar{D}_b$ can be expressed in terms of the RF-absorption power density $P_{\rm RF}$ from the relation

$$P_{\rm RF} = \left\langle \int d^3v \frac{1}{2} m v^2 Q(f_0) \right\rangle = \int d^3v \frac{1}{2} m v^2 \{Q\}(f_0) .$$

From Eq. (2), it is straightforward to obtain

$$P_{\rm RF} = M_\alpha \bar{D}_b \beta_\alpha \tau_s \sum_\ell \hat{\mathcal{J}}_\ell ,$$

where

$$\hat{\mathcal{J}}_\ell = \frac{3}{4\pi}\int d^3s \frac{s^3 \Theta(1-s)\Theta(s_{\parallel max} - s_\parallel)}{(s_c^3+s^3)^2}\left|J_{\ell-1}\left(\frac{k_\perp v_0}{\bar{\Omega}} s_\perp\right)\right|^2 .$$

Using this relation between $P_{\rm RF}$ and $\bar{D}_b$ to eliminate $\bar{D}_b$ from the expression for Γ_ψ, we obtain

$$\Gamma_\psi = \lambda_o \frac{\bar{R}}{r}\frac{P_{\rm RF}}{\epsilon_0} I\rho_0 \frac{\sum_\ell (k_\parallel v_0/\ell\bar{\Omega})\hat{\mathcal{I}}_\ell}{\sum_\ell \hat{\mathcal{J}}_\ell} ,$$

where $\rho_0 = v_0/\bar{\Omega}$ and $\epsilon_0 = 3.5$ MeV is the alpha birth energy.

If we define an average minor radial distance r for a noncircular flux surface using the toroidal flux ϕ,

$$\phi = \bar{B}_T \pi r^2 ,$$

we have

$$\frac{\partial\psi}{\partial r} = \bar{R}\bar{B}_P ,$$

and thus we obtain

$$\Gamma_r = \frac{1}{\bar{R}\bar{B}_P}\Gamma_\psi = \frac{P_{\rm RF}}{\epsilon_0}\rho_p \lambda_o \frac{\bar{R}}{r}\frac{\sum_\ell (k_\parallel v_0/\ell\bar{\Omega})\hat{\mathcal{I}}_\ell}{\sum_\ell \hat{\mathcal{J}}_\ell}$$

where $\rho_p = v_0/\bar{\Omega}_p$ is the poloidal gyroradius of alphas at the birth velocity. We find that the value of $R\lambda_o(k_\parallel v_0/\ell\Omega)\hat{\mathcal{I}}_\ell/r\hat{\mathcal{J}}_\ell$ can be adjusted to be of order unity for a proper choice of RF waves, as will be shown numerically. Thus, we can have the radial alpha particle flux of order

$$\Gamma_r \sim \frac{P_{RF}}{\epsilon_0} \rho_p$$

in magnitude.

Plots of Γ_r are shown in Fig. 2 as a function of $A = k_{\parallel} v_0 / \ell \bar{\Omega}$ and $C = k_{\perp} v_0 / \bar{\Omega}$ for $\ell = 3$. At $C = 30$ and $A = 0.5$, we get $\Gamma_r = 7 \times 10^{19} cm^{-1} s^{-1}$ from Fig. 2. $A = 0.5$ corresponds to $k_{\parallel} = 17 m^{-1}$ and $C = 30$ corresponds to $k_{\perp} = 340 m^{-1}$. The radial alpha flux is proportional to the ICRH power, hence larger radial flux can be obtained by increasing the wave power. For an ITER plasma with a surface area of approximately 750 m^2 and volume of approximately 1,000 m^3, the required alpha removal rate is approximately 3.6 $\times 10^{20}$ particles per second. If we were to put in the $1 + \cos\theta$ type of ICRF power over the whole plasma volume, we needed about 160 MW of ICRF wave power absorbed by the alpha particles to remove the required amount of alpha particles from ITER by using the present method with $C = 30$ and $A = 0.5$. If we use this scheme on the central burn area only, occupying one tenth of the total plasma volume, we may need about 16 MW of ICRF wave power. Further reduction of RF power is possible by better focusing of the wave power than the $1 + \cos\theta$ type of fucusing.

Since A is proportional to $k_{\parallel}/\ell$, higher harmonic heating will yield smaller radial transport for fixed $k_{\parallel}$. From Eq. (3) or (5), we also see that higher ℓ value will raise the energy of participating alpha particles unless the $k_{\parallel}$ value is raised at the same time. Positive values of the product $\lambda_o k_{\parallel}$ yield positive radial flux, and vice versa. For completeness, we write the expression for the total number of tail alpha particles n_α as follows:

$$n_\alpha = \int d^3 s \ \hat{f}_0 \simeq \beta_\alpha \tau_s \ln\left(\frac{v_0}{v_c}\right) \ .$$

We now discuss the design of an appropriate ICRF antenna. The possibility of asymmetric damping of the fast wave can be enhanced by tailoring the poloidal design of the antenna. The deposition pattern is determined by the Doppler width of the damping zone and the illumination of the damping layer by the antennas. An estimate of the deposition can be made using reduced order damping operators in conjunction with a fully three-dimensional wave solver, as a function of the poloidal mode spectrum, as shown in Ref. 7. For the case of interest here, the deposition may be asymmetrized by locating the radiating elements above and below the midplane, as shown in Fig. 3. A circular cross-sectional torus is used to simulate ITER geometry with $B_T = 4.55\ T$ and $n_e = 1 \times 10^{14} cm^{-3}$ at the magnetic axis, ICRF-wave frequency of 105 MHz corresponding to the third harmonic resonance surface through the magnetic axis, and $k_{\|} = 0.1 cm^{-1}$. The damping occurs at the third harmonic of the alphas and is maximum near the focus of the wave fields as determined by refraction. The resulting deposition is primarily on the fast alphas, by virtue of the large associated Doppler broadening of the resonance layer. Moreover, the antenna radiation pattern is aided by refraction in producing an asymmetric damping zone. A large degree of asymmetry in radiation power is easily obtained (the degree of asymmetry up to as large as 5), as demonstrated in the figures.

The toroidal wave spectrum may similarly be generated using a series of phased radiating elements to produce a nearly unidirectional antenna spectrum as demonstrated in Ref. 8. Estimates of the directivity indicate that 85 percent of the power can be channeled into a given toroidal direction using eight phased antennas.

III. Discussions

We have shown that a horizontally asymmetric power distribution of an ICRF wave can generate large radial transport of alpha particles whose energy is much less than their birth energy. The direction of the radial flux corresponds to the sign of the product $\lambda_o k_{\|}$. The RF-power needed in this scheme is much less than the alpha power. For ITER, about 16 MW of absorbed ICRF power is needed, compared to the estimated alpha power of 200 MW, to remove the required amount of alpha particles (3.6×10^{20} particles per second) from a quiescent central burning core (within one third of minor radius). This estimate is based upon a $1 + \cos\theta$ type variation of the ICRF power. With more effective localization of the RF power, as the one shown in Fig. 3, the required RF power can be further reduced. If we rely upon a moderate level of sawtooth activity to remove alpha particles from the central burning core, we can apply the present technique to the radius outside the $q = 1$ surface. A more detailed study of the present technique, which would further optimize the required wave characteristics and reduce the power requirement should be conducted to enable a more accurate prediction.

The ratio between the radial transport and birth rates of the alpha particles in the plasma core can be estimated as

$$\frac{\nabla \cdot \vec{\Gamma}_r}{\beta_\alpha} \sim \frac{\Gamma_r / L}{\beta_\alpha} \sim \frac{R}{r} \frac{P_{RF}}{\beta_\alpha \epsilon_0} \frac{\rho_0}{L},$$

where L is the alpha density scale length in the central reaction regime. Strictly speaking, the present analysis is valid for the case when the numerator ($\nabla \cdot \vec{\Gamma}_r$) is smaller than the denominator(β_α), due to the neglect of the particle sink term in Eq. (1). When RF power is high enough to violate this validity condition, a complete alpha density control can be achieved without relying upon any other loss mechanisms. As a matter of fact, ρ_0 at the birth velocity is not so small compared to the scale length L of the highly peaked central alpha density profile, and R/r is large at the central reaction radius; and thus we can easily see that the radial loss rate of the slowed-down tail-alpha density can be easily comparable to the alpha birth rate in the central reactive plasma, even for the RF-wave power density P_{RF} less than the alpha power density $\beta_\alpha \epsilon_0$. This simple comparison shows how powerful the present technique can be in controlling alpha transport. Since MHD activities can also remove alpha particles from the plasma (higher energy alphas, unfortunately), application of the present scheme in the presence of a moderate level of sawtooth activity can reduce the power requirement significantly.

Since the alphas participating in the present transport mechanism are at the lower energy range $v^2 \ll v_0^2$, the ratio of thermal energy loss $\nabla \cdot \vec{Q}_r$ carried out by this transport mechanism to the alpha energy birth rate will satisfy

$$\frac{\nabla \cdot \vec{Q}_r}{\beta_\alpha \epsilon_0} \sim \frac{(v^2/v_0^2)\nabla \cdot \vec{\Gamma}_r}{\beta_\alpha} \ll \frac{\nabla \cdot \vec{\Gamma}_r}{\beta_\alpha}.$$

Thus, the present mechanism is expected to be effective in removing alpha particles, but have little effect on alpha energy transport.

The wave numbers of the ICRF waves needed here are rather of standard values ($k_\parallel \sim 0.1\ cm^{-1}$, $k_\perp \sim 1\ cm^{-1}$). We have shown by a numerical simulation technique that a proper design of antenna arrays can yield in-out asymmetric power deposition of ICRF waves with mostly one-directional toroidal wave vectors at higher harmonic numbers.

We note here that the present scheme is not for removing alpha ash directly. It is to control the alpha ash population indirectly by removing the lower energy tail alphas before they are thermalized. Thus, the present scheme is intended to suppress the supply of alpha ash. In addition, the present scheme can also be applied to confine the alphas within the central reaction region, simply by changing the sign of $\lambda_o k_\parallel$ to the negative side. The change of the $k_\parallel$ sign will be easier to obtain and can be used to control dynamically the alpha density at the plasma center.

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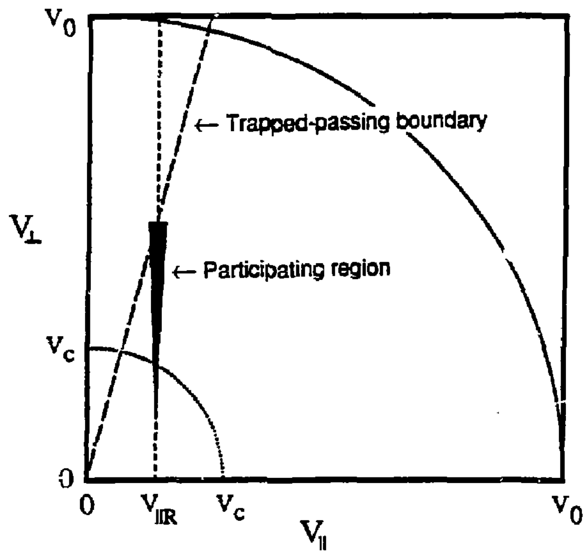
Figures

FIG. 1. Schematic diagram of the participating alpha particles in the velocity space. The emphasized part on the resonance line $v_\parallel = v_{\parallel R}$ represents the participating region.

FIG. 2. Radial alpha particle flux as function of $A = k_\parallel v_0/\ell\Omega$ for different values of electron temperature and $C = k_\perp v_0/\Omega$. The validity of the plots are limited to $A > \sqrt{r/R}$ due to the condition given by Eq. (6). $P_R = 2.3$ Watt$/cm^{-3}$, $\overline{B} = 4.85T$, $r = R/10$, $\ell = 3$, and safety factor=1.3 are chosen.

FIG. 3. Location of the antenna elements, and the radiation and power deposition profiles with horizontally asymmetric power deposition. (a) Radiation profile of the toroidal component of the RF magnetic field H_z. (b) Power deposition profile. About 3 to 1 asymmetry between out and in is achieved for this sample case. A circular cross-sectional torus is used with $B_T = 4.55\ T$ and $n_e = 1 \times 10^{14} cm^{-3}$ at the magnetic axis, ICRF-wave frequency of 105 MHz corresponding to the third harmonic resonance surface through the magnetic axis, and $k_\parallel = 0.1 cm^{-1}$.



V_0
← Trapped-passing boundary
$V_\perp$
← Participating region
V_c
0
0
$V_{\parallel R}$
V_c
V_0
$V_\parallel$

Fig. 1

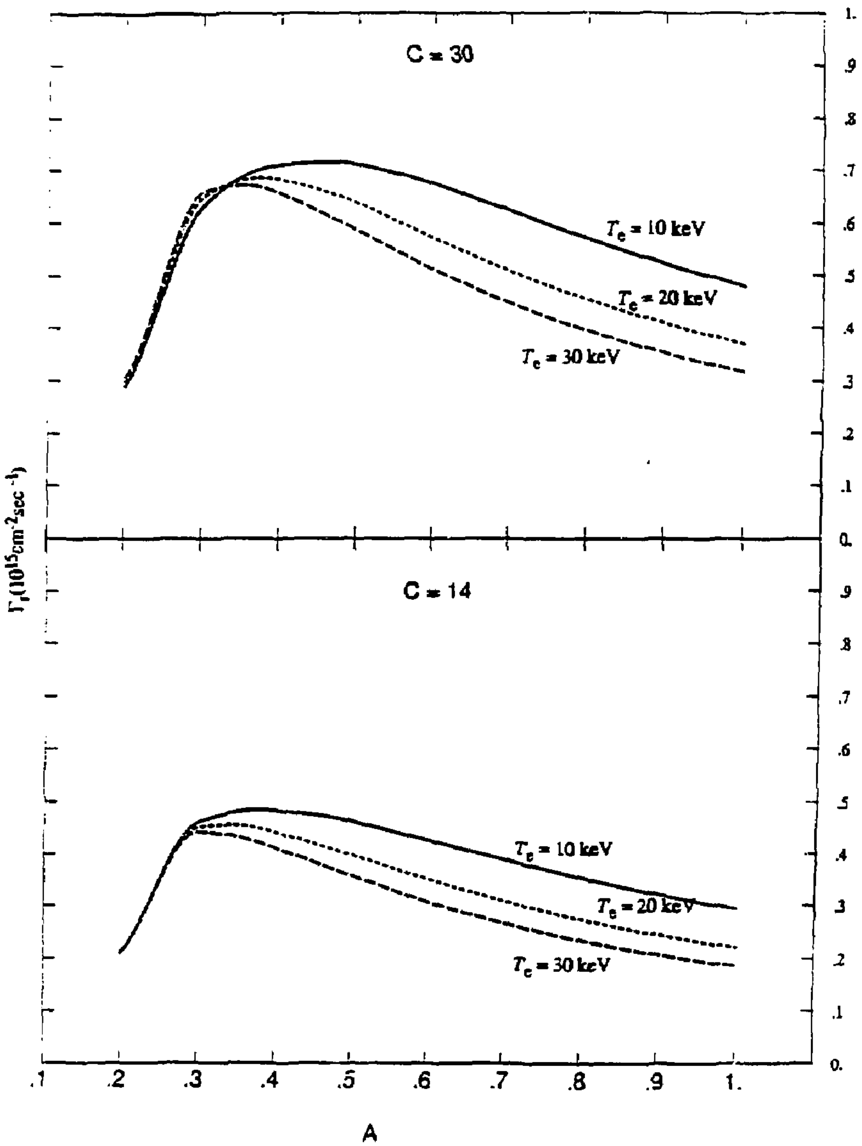

C = 30
C = 14
T_e = 10 keV
T_e = 20 keV
T_e = 30 keV
$\Gamma_i(10^{15}cm^{-2}sec^{-1})$
A
.1
.2
.3
.4
.5
.6
.7
.8
.9
1.
0.

Fig. 2

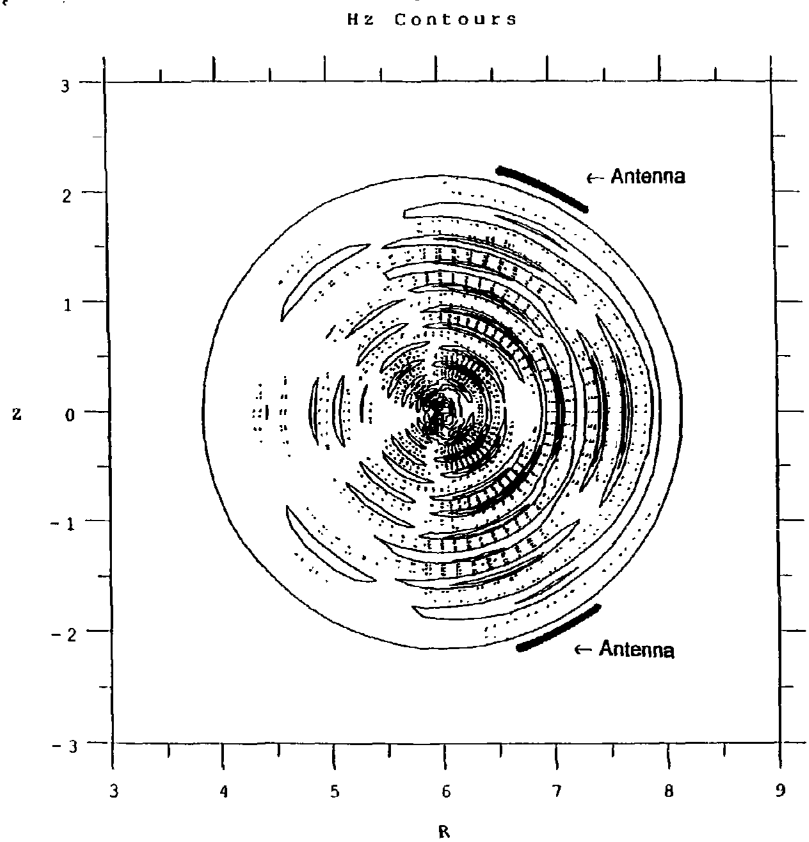


Fig. 3(a)

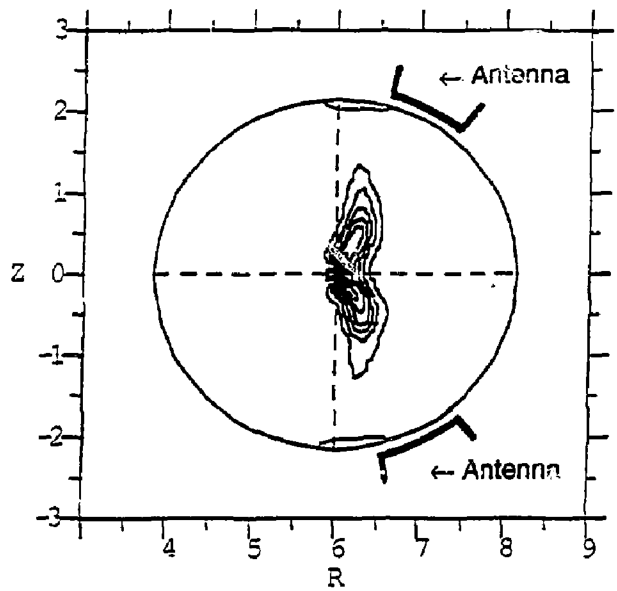


Fig. 3(b)