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Caigen LIU*, Mitsuru YAMAGIWA and Shangjie QIAN**

日本原子力研究所
Japan Atomic Energy Research Institute

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Production of Sheared Flow by Means of ICRF Heating in Tokamak Plasmas

Caigen LIU*, Mitsuru YAMAGIWA and Shangjie QIAN**

Department of Fusion Plasma Research
Naka Fusion Research Establishment
Japan Atomic Energy Research Institute
Naka-machi, Naka-gun, Ibaraki-ken

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Production of poloidal sheared flow during ICRF (ion cyclotron range of frequencies) heating is studied in tokamak plasmas. A new approach of inducing plasma poloidal rotation with rf wave is presented. The mechanism employed here is inducing a poloidal density asymmetry by rf wave and then destabilizing the anomalous Stringer spin up. The production of the poloidal density asymmetry during ICRF heating is analytically and numerically studied on the basis of the Fokker-Planck equation including the conventional collision and quasilinear rf diffusion operators. A model of jointing the rf power with the plasma rotation is presented to study the destabilization of the plasma rotation triggered by rf wave. A criterion upon destabilization of the plasma rotation in the presence of rf wave is given, which depends on the ratio of the electron slowing-down time to the ion collision time. The criterion is specified in the high rf power limit in the case of ICRF minority fundamental heating.

Keywords: Sheared Flow, Poloidal Rotation, ICRF Heating, Fokker-Planck Code, Tokamak

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* Visiting scientist from Southwestern Institute of Physics, P.R.China under the Scientist Exchange Program of Science and Technology Agency of Japan
** Southwestern Institute of Physics, Chengdu, Sichuan 610041, P.R.China

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トカマクプラズマにおけるICRF加熱によるシアー流の生成

日本原子力研究所那珂研究所炉心プラズマ研究部
Caigen LIU*・山極　満・Shangjie QIAN**

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トカマクプラズマのICRF（イオンサイクロトロン周波数帯）加熱下におけるポロイダルシアー流生成について検討を行い，プラズマのポロイダル回転を誘起する新しい手法の呈示を行った。ここで考慮した機構は，高周波によりポロイダル方向のプラズマ密度非一様性をつくりだすことによりStringerの提唱したプラズマスピンアップを不安定化させるというものである。ICRF加熱によるポロイダル密度非一様性の生成について，通常の衝突項および準線形高周波拡散項を含むフォッカープランク方程式に基づき，解析的かつ数値的に検討を行った。高周波パワーとプラズマ回転の関係を結びつけるモデルを提示し，高周波加熱により誘起される回転の解析を行った。密度非一様性に対して，回転をもたらし得るための条件を見いだし，それが衝突時間に対する減速時間の比に依存することを示した。特に，少数イオンを対象としたイオンサイクロトロン基本波加熱に対しては，ポロイダル回転を励起するための高周波パワーのしきい値の表式を解析的に見いだすことができた。

那珂研究所：〒311-01　茨城県那珂郡那珂町向山801－01
＊科学技術庁原子力研究交流制度に基づく研究員（中国西南物理研究所）
＊＊中国西南物理研究所

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1. Introduction

Improved confinement operations in tokamak plasmas are of numerous economic benefits to fusion reactors. In recent years, some researches have focused on how to operate tokamak under improved confinement modes especially in the fusion reactor conditions[1]. It is experimentally discovered that the improved confinement (H-mode) is often triggered by significant plasma rotation[2, 3]. Theoretically, it is also shown that the toroidal rotation of plasma can stabilize the MHD instability[4], and the sheared rotation either toroidal[5] or poloidal[6] can suppress the turbulence. Consequently it would improve the plasma confinement.

Therefore, one of approaches to achieve improved confinement operations is to induce plasma rotation. Some successful examples to conduct H-mode operations by rotating plasma are the imposed biased electrodes experiments performed in many tokamaks[7, 8, 9], and NBI experiments[5]. However, the methods of using biased electrodes and NBI to achieve H-mode have a fatal disadvantage; it is difficult to extend these methods to the reactor conditions.

Experimentally, the poloidal rotation observed in tokamak plasmas is significant only in the edge plasma region. It is desirable to seek for approaches of inducing plasma poloidal rotation, which are suitable for reactors, controllable externally and effective for the core plasma region. To this aim, some approaches of inducing plasma rotation with external rf waves have been proposed by several authors; Alfven wave[10], Ion Bernstein wave[11] and Ion cyclotron wave[12, 13], from the viewpoint of utilization of rf momentum and pondermotive force. It is observed experimentally that significant toroidal rotation is produced during ion cyclotron range of frequencies (ICRF) minority heating in JET[14].

Exploration of mechanisms of plasma rotation is also an attractive subject. Some research results of rotation mechanisms have been reported: ion orbit loss in the edge region[15]; spontaneous poloidal spin-up[16, 17] and its revised version[18, 19]; change of anomalous thermal conductivity[20]; Reynolds stress[21]; etc. Hassam has proposed[17] that any asymmetric particle or power deposition in tokamak plasmas might induce an additional rotation. In his serial works[18, 22, 23], NBI and gas puffing as ways of producing plasma rotation have been employed. Hassam's theory is also given the indirect support from the TFTR experiments[24].

During rf wave cyclotron heating in tokamak plasmas, resonant particles diffuse preferentially along the velocity direction perpendicular to the magnetic field. With the increase of rf power, more and more particles tend to be trapped, and tip points of the trapped particle orbit move close to the resonant layer. We refer to this phenomenon as the resonance localization. The resonance localization is a special feature of the strong rf cyclotron heating[25-29]. The cyclotron heating also makes the passing particles contribute to the density asymmetry. Some sequences of the resonance localization have been studied in several papers, such as the production of the poloidal electric field[26], additional particle radial flux[27], etc. The resonance localization can also result in accumulation of the resonant particles in the outer plasma region. Consequently, a poloidally asymmetric particle distribution is formed. The magnitude of asymmetry can be manipulated by changing the position of the resonant layer and the magnitude of the rf power. It is obvious that the rf wave power, not the rf momentum, is the origin of the poloidally asymmetric particle distribution. Therefore, with the increase of the rf wave power, the poloidal plasma rotation would be destabilized.

Experimentally, there seems to exist power threshold for H-mode achievement during auxiliary heating, for example, ICRH in ASDEX[30] and JET[31]; ECRH in DIII-D[32] and JFT-2M[33], indicating that the H-mode operation may have something to do with the rf power.

In this work, we explore the possibility of production of the poloidal rotation by ICRF heating in tokamak plasmas with the concentric circular magnetic surfaces. In the next section, we give a physical description of rf wave inducing plasma poloidal rotation. In Section 3, the poloidal density asymmetry due to the resonance localization is analytically and numerically studied by using the Fokker-Planck equation including the quasilinear rf diffusion operator. In Section 4, a set of MHD equations with a model of rf inducing the density asymmetry is presented to obtain the criterion upon the density asymmetry for triggering the poloidal plasma rotation in the presence of rf waves. The rf power threshold for excitation of the poloidal rotation in the high rf power limit in the ICRF fundamental minority heating is also presented. Finally, in Section 5, a brief summary and discussion are given.

2. Physical mechanism of rf inducing of plasma rotation

In tokamaks, the strength of the magnetic field can be approximately given by, $B = B_0 / h$, with $h = 1 + \varepsilon\cos\theta$, where ε is the inverse aspect ratio ε=r/R. When particles move along the magnetic field line, they feel lower magnetic field on the outer side of the torus and higher magnetic field on the inner side. In the absence of collisions and rf interactions, the guiding-center orbit of particle motions can be described by the three invariants; the energy $E = m\mathrm{v}^2/2$, the magnetic moment $\mu = mv_\perp^2/2B$ and the toroidal angular momentum $P = mv_{//}R + \Psi$, with Ψ being the poloidal magnetic flux. Therefore, only the particles with higher parallel velocity can pass the magnetic hill, while the particles with lower parallel velocity fail to pass the magnetic hill. Those particles that fail to pass the hill are reflected at the places where their parallel velocities vanish. They are called trapped particles. The point at which $v_{//} = 0$ is called the turning point. The effect of particle trapping results in more particles staying in the lower magnetic field region, that is in the outer side of the cross section of the torus. In this way, the poloidally asymmetric particle distribution is formed on a magnetic surface of a tokamak plasma. Physically, the poloidal density asymmetry is an intimate phenomenon in a tokamak plasma, but the magnitude of the density poloidal asymmetry is trivial in the normal situations.

The poloidal density asymmetry in a tokamak plasma can be enhanced by the neoclassical transport, even in the purely ohmic heating[34]. In the plateau regime, where the neoclassical transport coefficients are almost independent of the collisionality, (that is $1 >> v_c / \omega_t >> \varepsilon^{3/2}$ with v_c being the collision frequency, $\omega_t \equiv \mathrm{v}_{th}/L$ the transit frequency, L the scale length manipulated by changing the position of the resonant layer and the magnitude of the rf power. It is obvious that the rf wave power, not the rf momentum, is the origin of the poloidally asymmetric particle distribution. Therefore, with the increase of the rf wave power, the poloidal plasma rotation would be destabilized.

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$$\frac{\tilde{n}}{n_0} \sim \frac{\pi^{1/2}}{2} \frac{\rho_{i\theta}}{L_T} \varepsilon \sin\theta,$$

where n_0 is the flux surface averaged density of ions. The tilder denotes the poloidal fluctuation component of the ion density, $\rho_{i\theta}$ is the ion Lamour radius measured by the poloidal magnetic field, L_T the characteristic length of the ion temperature, and θ is the poloidal angle. In the collisionless or banana regime($v_c / \omega_t << \varepsilon^{3/2}$), where before particles suffer one collision, it may poloidally move several times, the ion density poloidal distribution is given by

$$\frac{\tilde{n}}{n_0} \sim 1.81 \left[\frac{qRv_i}{v_{ti}\varepsilon^{3/2}} \right] \frac{\rho_{i\theta}}{L_T} \varepsilon \sin\theta,$$

where $q \equiv rB_T / R_0 B_p$ is the safety factor, and v_i is the ion collision frequency. The magnitude of this poloidal density asymmetry is, however, too small to induce the plasma rotation. The electron density asymmetry has the similar form. Therefore, the ion density asymmetry is more significant than the electron density asymmetry because of $\rho_{i\theta} >> \rho_{e\theta}$.

In a tokamak plasma with rf wave cyclotron heating, the magnitude of the poloidal density asymmetry would be significantly enhanced because of perpendicular diffusion of particles. The mechanism of rf inducing of poloidal density asymmetry[26, 35] is given as follows (see Fig. 1). When the cyclotron resonant condition $\omega - k_{//}v_{//} = n\omega_c$ is satisfied, the resonant particles diffuse along the tangential direction of the curve of $v_\perp^2 + (v_{//} - \frac{\omega}{k_{//}})^2 = const$ in velocity space.

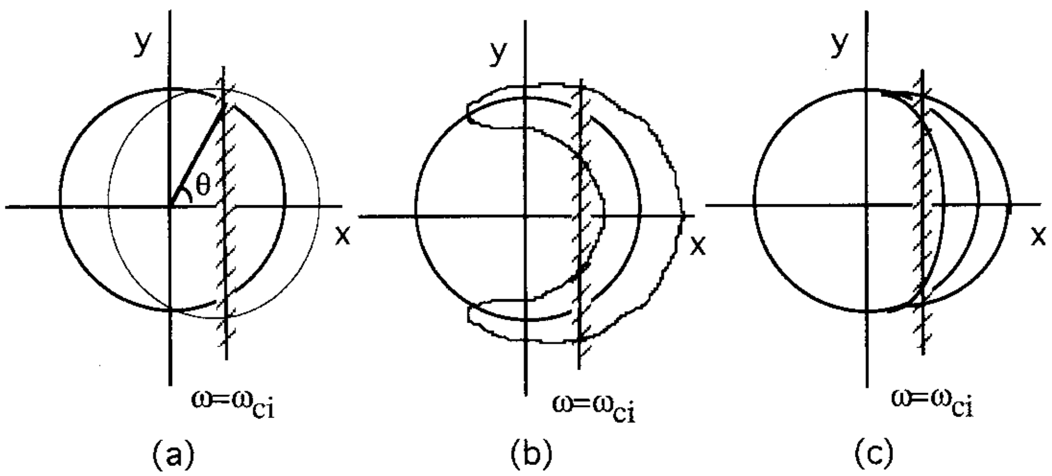

Fig. 1. Schematic draw of the evolution of the particle orbit while interacting with rf wave. a) A passing particle interacts with the ICRF wave and prefers to diffuse along the perpendicular velocity direction. b) After absorbing the rf power to certain degree, the particle is trapped. c) Further continuing to absorb the rf power, the turning point of the particle moves close to the resonance layer.

Figure 2 displays the numerical results of the relationship between the fraction of the trapped particles and the rf power density during ICRF heating with a two-dimensional bounce averaged Fokker-Planck code[36]. The population of the trapped particles increases with the rf power density. As the perpendicular energy of the particles becomes higher, the turning points of the trapped particles move close to the resonance layer. This is called resonance localization. If the resonance layer is located on the lower magnetic field side, the particles are likely to accumulate in the region of the outer torus with increase of the rf power. Thus, the poloidal density asymmetry is enhanced. The magnitude of the poloidal density asymmetry can be changed by shifting the resonance layer.

Hassam suggested an anomalous plasma spin-up mechanism[22, 23] on the basis of the Stringer spin-up theory[16]. He pointed out that under the condition of asymmetric particle deposition or extraction, a tokamak plasma has a tendency to spontaneous spin-up. Any approaches, which can produce a density in-out asymmetric distribution on a magnetic surface, would induce poloidal plasma rotation. The physical mechanism of this poloidal spin-up may be described as follows. For an asymmetric particle source $S(r,\theta)=\bar{S}(r)+\tilde{S}(\theta)$, $\tilde{S}$ would induce a parallel flow in order to avoid the poloidally localized accumulation of particles. For simplicity, this flow obeys the equation $\nabla_{//}\tilde{u}_{//}=\tilde{S}/n$. Recalling the force equilibrium equation in the parallel direction, $\tilde{u}_{//}$ would produce an inhomogeneous density $\tilde{n}$ along the magnetic surface,

$$n\frac{u_\theta}{r}\frac{\partial}{\partial\theta}\tilde{u}_{//}\approx-\frac{c_s^2}{qR}\frac{\partial}{\partial\theta}\tilde{n}, \tag{1}$$

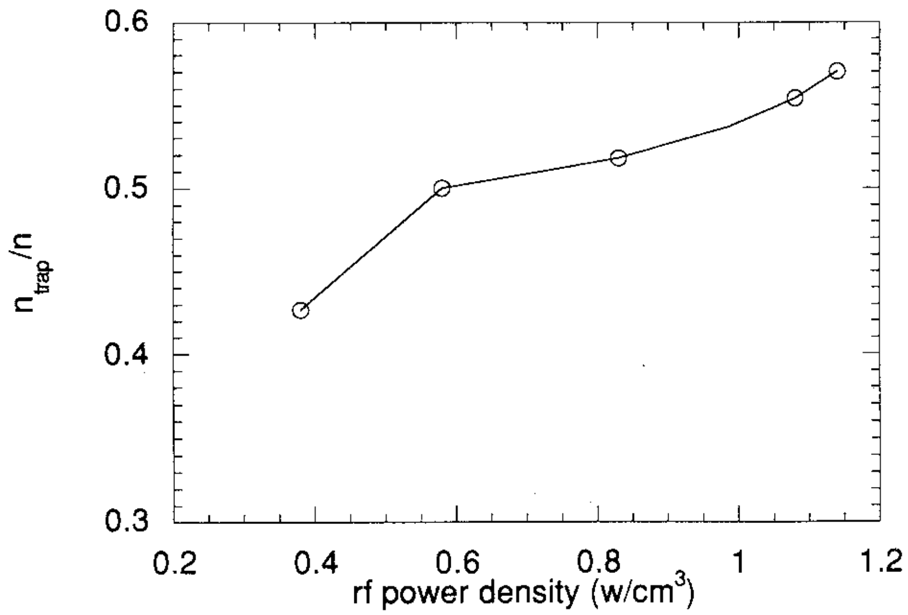

Fig. 2. Fraction of trapped particles versus rf power density during D(H) minority heating at ε=0.2.

where u_θ is the poloidal component of the flow. In a tokamak plasma, the effective gravity, $g_{eff} = {2c_s^2}/{R}$, is outward toroidally. This effective gravity would affect the plasma if the plasma density varies on the magnetic surface. The poloidal component of this force, $\vec{F} = \oint d\theta \tilde{n} m \vec{g}_{eff} \cdot \vec{\theta}$, would accelerate the plasma in the poloidal direction. The resultant equation of motion is

$$nm\frac{\partial u_\theta}{\partial t} = -\oint \frac{d\theta}{2\pi} \tilde{n} m \vec{g}_{eff} \cdot \vec{\theta}. \qquad (2)$$

It is easy to understand that for an in-out asymmetric source, $\tilde{S} \propto \cos\theta$, i.e., peaked in the low field side of the torus, a parallel flow $\tilde{u}_{//} \propto \sin\theta$ would be produced. This parallel flow will give rise to an up-down asymmetric density $\tilde{n} \propto -u_\theta \sin\theta$, and then $\tilde{n}$ will drive a poloidal flow due to the effective gravity. Similarly, we can conclude that $\tilde{S} \propto -\cos\theta$ will result in an oscillation of u_θ and that the rotation is stable for the case of the density up-down asymmetry. The nature of this kind of notation is of Stringer spin-up. In this way, a poloidal in-out density asymmetry produces a poloidal plasma rotation.

Combining the poloidal density asymmetry produced by the resonance localization in the rf wave cyclotron heating with the poloidal plasma spin-up destabilized by the in-out density asymmetry, the poloidal plasma rotation can be actively produced.

3. Production of poloidally inhomogeneous density by ICRF wave

Auxiliary heating by use of rf waves in tokamak plasmas has been developed for a few decades. Many papers have shown that the rf cyclotron wave can produce the particle density asymmetry. Hsu[26] et al. have studied the phenomenon of resonance localization during rf cyclotron heating. The research was carried out under the assumptions of banana regime, null parallel wave number, non-relativity and large aspect ratio tokamak. The induced density asymmetry can be evaluated as $\tilde{n} / n_0 \sim 0.6\delta\varepsilon$, where δ is the ratio of the effective rf heating rate to the collision frequency. Taguchi[27] has considered more general conditions. The result obtained by Taguchi is different from that of Hsu in that the sign of $\tilde{n}$ may be changed with varying plasma parameters. Contribution from the collisional diffusion to the density asymmetry is, however, neglected in their researches.

It is understandable that, in the collisionless regime, rf wave can induce a considerably large inhomogeneity of poloidal particle distribution because of the particle trapping. Even in higher collisionality regime, where banana orbits are deteriorated, rf cyclotron heating may produce an additional asymmetry of particles[37]. This asymmetry can be derived as $\frac{\tilde{n}}{n_0} \approx c\varepsilon \frac{\nu_{ee}}{\nu_t}\frac{\nu_{rf}}{\nu_t}$ on the basis of the drift kinetic equation, where ν_{ee}, ν_{rf}, and ν_t are the collision frequency, the effective rf heating rate and the transit frequency correspondingly, although the coefficient of the above where u_θ is the poloidal component of the flow. In a tokamak plasma, the effective gravity, $g_{eff} = 2c_s^2/R$, is outward toroidally. This effective gravity would affect the plasma if the plasma density varies on the magnetic surface. The poloidal component of this force, $\vec{F} = \oint d\theta \tilde{n} m \vec{g}_{eff} \cdot \vec{\theta}$, would accelerate the plasma in the poloidal direction. The resultant equation of motion is

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The temperature anisotropy is one of the special features in the rf cyclotron heated tokamak plasmas. From the Fokker-Planck equation including the collision diffusion and quasilinear rf diffusion operators[35]

$$\frac{\partial f}{\partial t} = C(f) + Q(f), \tag{3}$$

we can get the deformed particle distribution function. In the above equation, C(f) and Q(f) are the collision and quasilinear operators, respectively. Here,

$$C(f) = \frac{1}{v^2}\frac{\partial}{\partial v}(-\alpha v^2 f + \frac{1}{2}\beta v^2 \frac{\partial f}{\partial v}) + \frac{1}{4v^2}\frac{\partial}{\partial \mu}(1-\mu^2)\frac{\partial(\gamma f)}{\partial \mu},$$

$$Q(f) = \frac{1}{v_\perp}\frac{\partial}{\partial v_\perp} v_\perp D \frac{\partial f}{\partial v_\perp},$$

and $\mu = v_{//}/v$ is the cosine of the pitch angle, α, β, γ and D are the collision coefficients and quasilinear rf diffusion coefficient[35]:

$$\alpha = -\sum_j c_j l_j^2 (1 + \frac{A}{A_j}) G(l_j v) + \sum_j \frac{c_j}{2v^2}[\phi(l_j v) - G(l_j v)],$$

$$\beta = \sum_j \frac{c_j}{v} G(l_j v),$$

$$\gamma = \sum_j \frac{c_j}{v}[\phi(l_j v) + G(l_j v)],$$

$$\phi(x) = \frac{2}{\sqrt{\pi}}\int_0^x \exp(-y^2)dy,$$

$$G(x) = \frac{\phi(x) - x\phi'(x)}{2x^2},$$

$$D = \sum_l \frac{\pi z_i^2 e^2}{8 m_i^2}|E_\pm|^2 J_{l\pm1}^2(k_\perp \rho)\delta(\omega - k_{//} v_{//} - l\Omega),$$

$$c_f = \frac{8\pi n_f z_f^2 z^2 e^4 \ln\Lambda}{m_f},$$

with $l_j^2 = m_j / 2kT_j$, $J_l(x)$ being the Bessel function of the l-th order, and A the ion mass normalized by the proton mass. It is difficult to solve Eq. (3) analytically in two-dimensional velocity space except in some limited conditions[35, 38].

3.1 Analytical formulation

In the strong rf heating, the distribution function of the resonant particles is significantly anisotropic. With increasing the rf power, the perpendicular temperature of the resonant particles increases monotonously. The increase of the parallel temperature of the resonant particles is mainly dominated by the collisional pitch angle scattering. Therefore, the parallel distribution function may be considered as Maxwellian. In the case of lower rf power, the pitch angle scattering may sustain the isotropic distribution. For energetic ions in the ICRF fundamental heating, integrating Eq. (3) over $v_{//}$, we can obtain the distribution function in the perpendicular velocity direction as[35]

$$\ln f(v_\perp) = -\frac{2E}{(2+3\xi)T_e}\{1 + \frac{R_j[(2A+A_j)(2+3\xi)T_e - 4AT_j]}{2AT_j(2+2R_j+3\xi)} H(\frac{E}{E_j})\}, \qquad (4)$$

where $\xi = \langle p \rangle \tau_s / 3n_i T_e$, $\langle p \rangle$ is the absorbed rf power density, τ_s is the electron slowing down time,

$$H(x) = \frac{1}{x}\int_0^x \frac{du}{1+u^{3/2}} \approx 2.418x^{-1} - 2x^{-3/2} + 0.5x^{-3} - \ldots, \quad R_j = n_j z_j^2 l_j / n_e l_e$$

and $$E_j = \frac{AT_j}{A_j}\left[\frac{2+2R_j+3\xi}{2\varepsilon(2+3\xi)}\right]^{2/3}.$$

Defining an effective perpendicular temperature, $T_\perp = -[d(\ln f(v_\perp))/dE]^{-1}$, we can get

$$\frac{1}{T_\perp} = \frac{2}{(2+3\xi)T_e}\{1 + \frac{R_j[(2A+A_j)(2+3\xi)T_e - 4AT_j]}{2AT_j(2+2R_j+3\xi)} H(\frac{E}{E_j})\}. \qquad (5)$$

Note that the above equation is valid only for energetic ions.

The parallel temperature can be characterized as

$$T_{//} = 3.7kT_e[\frac{2A^{1/2}}{n_e}\sum_j n_j z_j^2]^{2/3}. \qquad (6)$$

The ratio of the perpendicular temperature to parallel temperature versus the rf power density is shown in Fig. 3.

The temperature anisotropy is correlated with the poloidal density asymmetry. The dependence of the magnitude of density asymmetry on the temperature anisotropy can be derived as follows. The force balance equation in the steady state may approximately expressed as,

$$\nabla \cdot \vec{P} - nq_i(\vec{E} + \vec{u} \times \vec{B}) = 0, \tag{7}$$

where $\vec{P} = nkT$ is the pressure tensor, $q_i = Z_i e$ the quantity of charge for ions. The parallel component of the equation is given by

$$\frac{\partial P_{//}}{\partial \theta} + \frac{P_\perp - P_{//}}{B}\frac{\partial B}{\partial \theta} = -nq_i\frac{\partial \phi}{\partial \theta}. \tag{8}$$

Dividing n into the magnetic surface average part and the poloidally fluctuating part; $n = n_0(\Psi) + \tilde{n}(\theta)$, expanding the above equation in orderings, and then integrating the first order equation over θ, we can get

$$\frac{\tilde{n}}{n_0} - \varepsilon(\frac{T_\perp}{T_{//}} - 1)\cos\theta = -\frac{q_i\tilde{\phi}}{kT_{//}}. \tag{9}$$

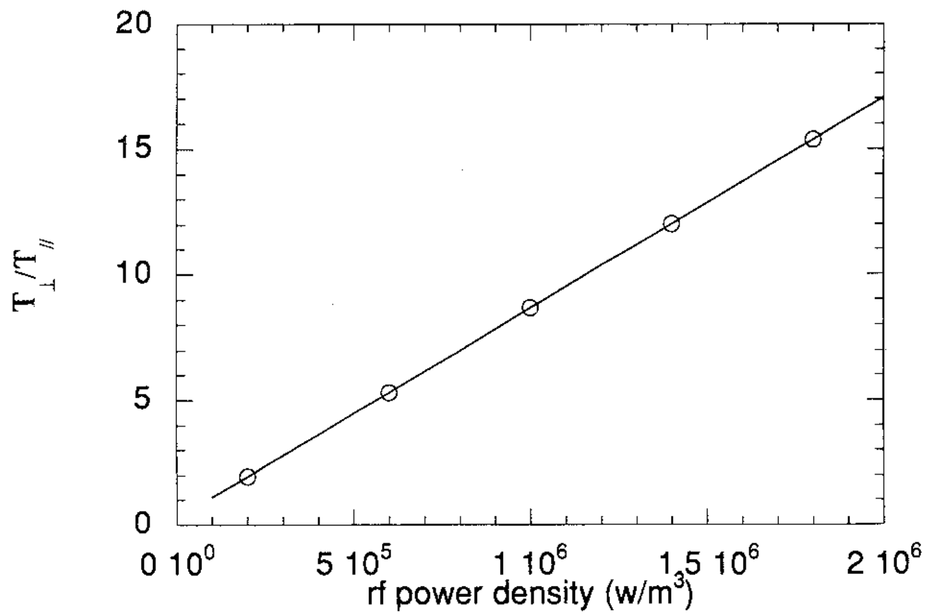

Fig. 3. The ratio of perpendicular temperature to parallel temperature vs. the rf power density. The plasma parameters are T_e=6 kev and n_e=5×10^{19} /m^3 in D(H) plasma.

By using $\frac{\tilde{n}_e}{n_{e0}}=\frac{e\tilde{\phi}}{T_e}$; $\frac{\tilde{n}_i}{n_{i0}}=-\frac{z_i e\tilde{\phi}}{T_i}$; $\frac{\tilde{n}_I}{n_{i0}}=-\frac{z_I e\tilde{\phi}}{T_I}$, where subscript I denotes impurity species and subscript i the background ion, and quasi-neutrality condition, we can get

$$\tilde{n}=(\frac{n_{e0}e}{kT_e}+\frac{z_i^2 e n_{i0}}{kT_i}+\frac{z_I^2 e n_{i0}}{kT_I})\frac{e\tilde{\phi}}{q_i}. \tag{10}$$

Combining above two equations, it gives

$$\tilde{n}=\varepsilon(\frac{T_\perp}{T_{//}}-1)N\cos\theta, \tag{11}$$

where $N^{-1}=\frac{1}{n_0}+\frac{q_i^2}{e^2 kT_{//}}(\frac{z_i^2 n_{i0}}{T_i}+\frac{z_I^2 n_{i0}}{T_I}+\frac{n_{e0}}{T_e})^{-1}$. Equation (11) indicates relationship between the temperature anisotropy and the poloidal density asymmetry.

From Eqs (5), (6) and (10), we can approximately get an asymptotic expression for relationship between the poloidal density perturbation and the absorbed rf power density in the ICRF fundamental minority heating,

$$\frac{\tilde{n}}{N}=(\frac{(2+3\xi)}{3.7[\frac{2A^{1/2}}{n_e}\sum_j n_j z_j^2]^{2/3}}-1)\varepsilon\cos\theta. \tag{12}$$

3.2 Numerical calculation

Equation (12) gives only an asymptotic value in the high rf power limit. For the moderate rf power, it is difficult to obtain an analytical expression. Further for a more strict treatment, analysis in two dimensional velocity space is necessary. This can be done only by numerical calculation. We study the evolution of the ion distribution function and the poloidal density inhomogeneity with a two-dimensional bounce-averaged Fokker-Planck code[36]. It is clearly shown that ICRF heating may create a significant poloidal density fluctuation even in moderate rf power heating.

In our numerical calculation, we employ the following parameters: the major radius R=3.0 m; the minor radius a=1.6 m; the toroidal magnetic field B_t=2.8 T; the ion temperature T_i=5 kev, the electron temperature T_e=6 kev; the electron density $n_e=5\times10^{19}$ /m^3. Figure 4 shows time evolution of the poloidal density profile when the rf electric field strength E_+=340 v/m at r=0.7 m for n_H/n_e=10 % in D(H) minority fundamental on-axis heating. Figure 5 shows poloidal density profiles for different rf power densities. It is shown that a 12 % density difference is produced at a moderate rf power density of p_{rf}=0.24 w/cm^3.

In the case of D(H) two-component fundamental heating, n_H/n_e=50 %, the rf power density

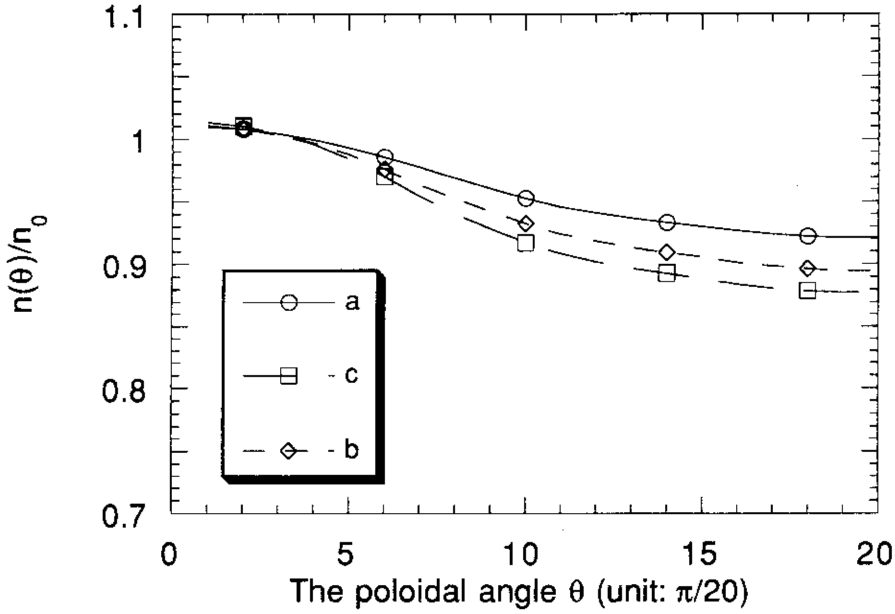

Fig. 4. Time evolution of the poloidal density profile during D(H) minority heating. The electric field E_{+}=340 v/m. r=0.7 m. a) t=20 ms; b) t=60 ms; c) t=100 ms.

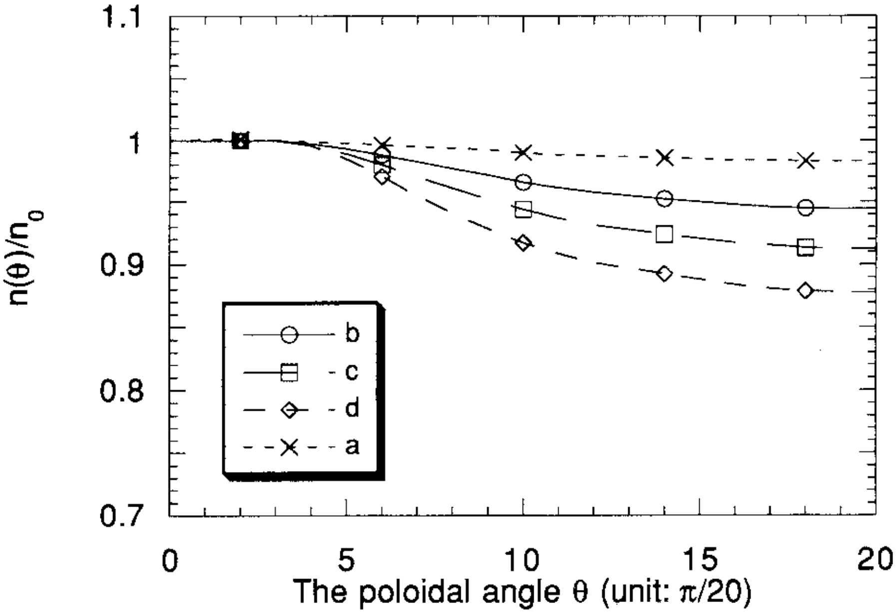

Fig. 5. Poloidal density profiles for different rf power density during D(H) minority heating. Parameters are t=0.1 s, r=0.7 m, p_{rf}=0.03 w/cm^3 for curve a; p_{rf}=0.1 w/cm^3 for curve b; p_{rf}=0.16 w/cm^3 for curve c; p_{rf}=0.24 w/cm^3 for curve d.

for producing the same level of the poloidal density perturbation is larger as shown in Fig. 6. Amplitude of the density difference increases with the rf power density as shown in Fig. 7.

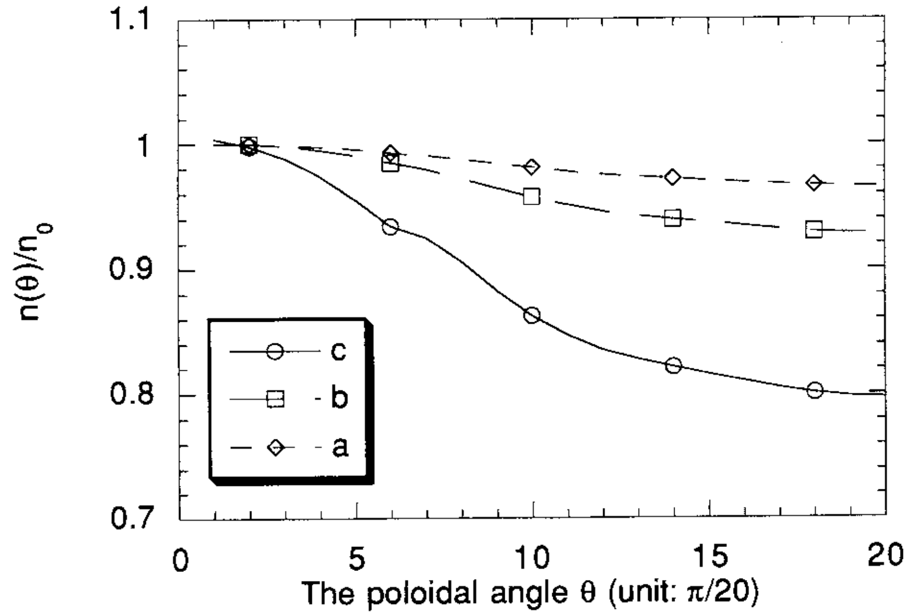

Fig. 6. Poloidal density profiles for different rf power density during D(H) two component heating. Parameters are t=0.1 s, r=0.5 m and p_{rf}=0.3 w/cm^3 for curve a; p_{rf}=0.67 w/cm^3 for curve b; p_{rf}=2.27 w/cm^3 for curve c.

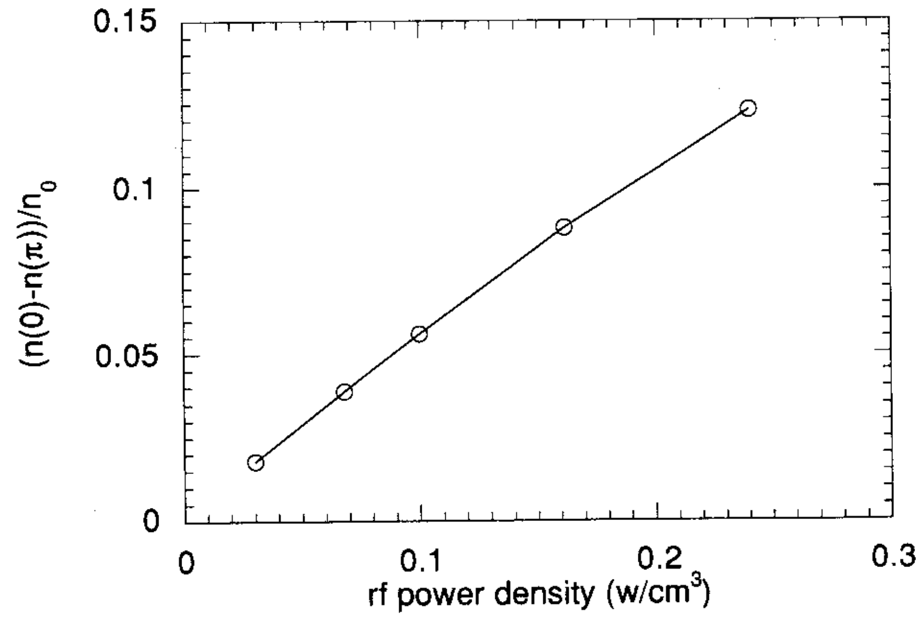

Fig. 7 Amplitude of the density perturbation versus the rf power input during D(H) minority heating. For t=0.1s and r=0.7m.

The resonant position may also affect the formation of poloidal density asymmetry. Figure 8 shows that the magnitude of the density perturbation increases as the resonant layer moves towards the lower field side.

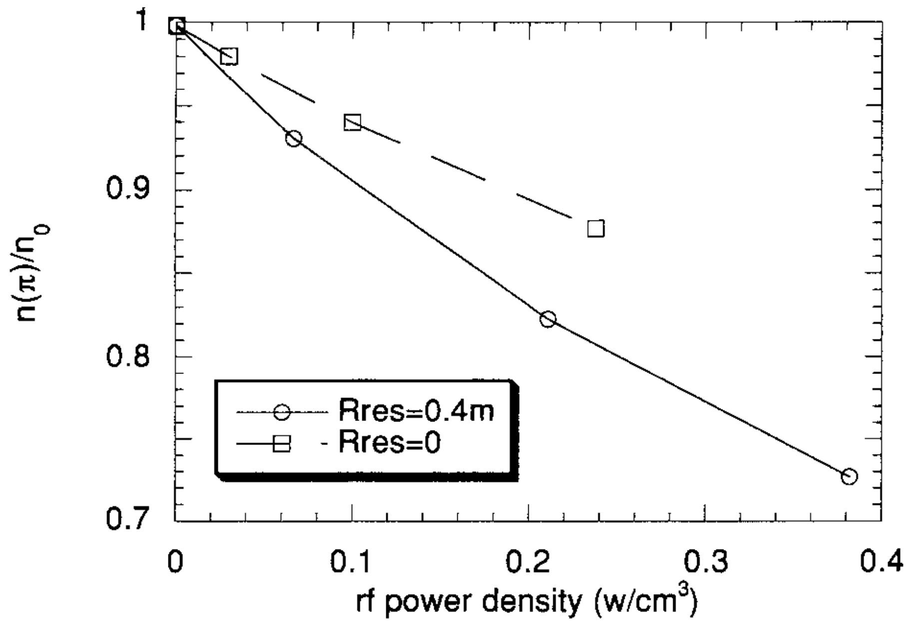

Fig. 8. The lowest density on a magnetic surface versus the rf power density during D(H) minority heating for different positions of the resonant layer of $r_{res}=0$ and $r_{res}=0.4$ m.

4. Poloidal rotation induced by density asymmetry

The equation describing the evolution of the poloidal rotation speed induced by the poloidal density inhomogeneity in the presence of rf waves is derived on the basis of MHD equations[23], of which, the continuity equation is

$$\frac{\partial n}{\partial t}+\nabla\cdot n\vec{u}_{\perp}+\vec{B}\cdot\nabla(nu_{//}/B)=S-\frac{1}{r}\frac{\partial}{\partial r}r\Gamma_r+\frac{\tilde{n}}{\tau_0}, \qquad (13)$$

where τ_0 is the characteristic time for the formation of the poloidal density profile induced by rf waves. The right hand side of the above equation represents the origin of the plasma rotation. The poloidally asymmetric source S and the poloidally asymmetric radial particle flux Γ_r can result in the poloidal density perturbation, and hence in the plasma rotation. These effects have been studied in Refs [19, 23]. We introduce the third term for modeling the contribution to the poloidal plasma rotation from the rf wave inducing the density perturbation. In the ICRF heating the slowing down of energetic ions is dominated by colliding with electrons. In the time scale of the

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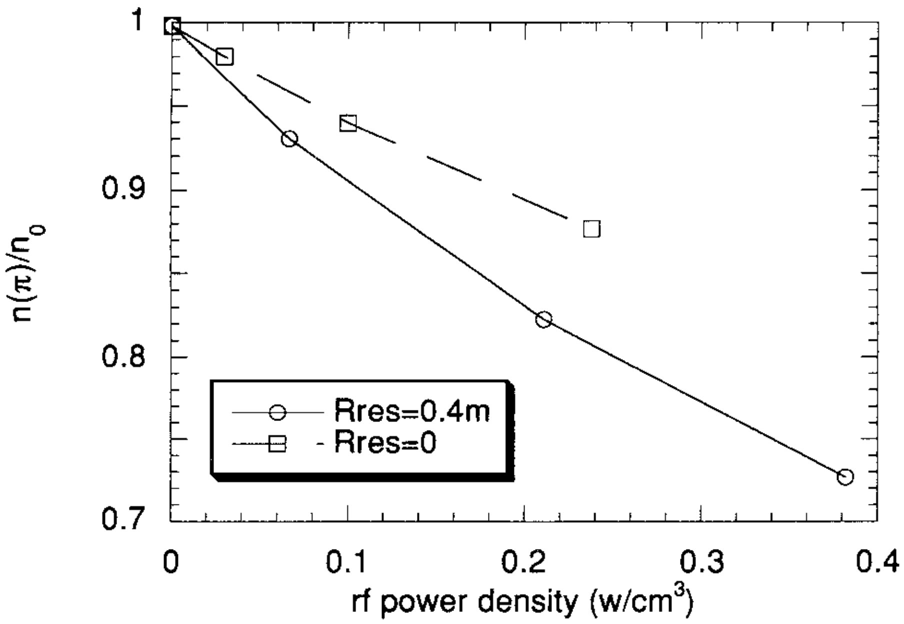

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$$\frac{\partial n}{\partial t} + \nabla \cdot n\vec{u}_{\perp} + \vec{B} \cdot \nabla(nu_{//} / B) = S - \frac{1}{r}\frac{\partial}{\partial r} r\Gamma_r + \frac{\tilde{n}}{\tau_0}, \tag{13}$$

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The other equations relevant to the analysis of the poloidal plasma rotation are as follows,

$$nm(\frac{\partial \vec{u}}{\partial t}+\vec{u}\cdot\nabla\vec{u}) = -T\nabla n+\vec{j}\times\vec{B}-mS\vec{u}, \tag{14}$$

$$\nabla\phi = \vec{u}\times\vec{B}, \tag{15}$$

$$\nabla\cdot\vec{j}=0, \tag{16}$$

where ϕ is the electric potential. In this paper, the effects of the particle source and the radial transport are neglected. The derivation procedure in the following is similar to that done by Hassam[23]. At the first, we assume that the plasma rotation is subsonic, and the ordering

$$\frac{\partial}{\partial t}\sim\frac{u}{a}\sim\frac{\tilde{n}}{n\tau_o}\sim\gamma_{MP}<<\frac{c_s}{qR},$$

is satisfied, where γ_{MP} is the magnetic pump damping rate, c_S is the sonic speed. Then, we expand Eqs (13) and (14) according to the above ordering. We can get an evolution equation of the poloidal rotation speed during strong rf cyclotron heating as,

$$\frac{\partial Y}{\partial t} = -GY+R, \tag{17}$$

where $Y = <\vec{u}\cdot\vec{B}_p/B_p^2>$ is proportional to the local poloidal velocity. G and R represent damping and driving contributions to the plasma poloidal spin-up:

$$G=\frac{<\tilde{n}\Delta_B>}{n\theta^2\tau_0}+\frac{\eta_0}{nm\theta^2}\frac{<(\vec{b}\cdot\nabla B)^2>}{<B^2>},$$

$$R=-\frac{<\tilde{n}\tilde{Y}B^2>}{n\theta^2<B^2>},$$

where $\Delta_b = {}^{(B^2-<B^2>)}/_{<B^2>}$. It is obvious that the poloidal rotation would be destabilized if G<0. If G>0, the poloidal rotation would be stabilized, but the plasma might be rotated poloidally in the steady state dependent on the value of R. The quantity R represents the contribution of the poloidal variation of the parallel flow. For the case of rf cyclotron heating, rf momentum input can be neglected and the density perturbation is in-out asymmetric, R vanishes because $\tilde{Y}$ is up-down asymmetric whereas B is asymmetric. Thus, if the following inequality holds

$$\frac{\eta_0 \tau_0}{m}\frac{<(\vec{b}\cdot\nabla B)^2>}{<B^2>} < -<\tilde{n}\Delta_B>, \tag{18}$$

Y will grow up exponentially, i.e., plasma will spin up poloidally. The left hand side of the above inequality represents stabilization due to the magnetic pump damping. The parallel ion viscous coefficient in the long mean free path region (i.e., banana region) is of the form $\eta_0 = \varepsilon^{-3/2} q^2 R^2 n_i m_i v_{ii}$ [19], where the ion collision frequency is given by $v_{ii} = n_i z^4 e^4 \ln\Lambda / 12\pi^{3/2}\varepsilon_0^2 m^{1/2} T_i^{3/2}$. Assuming $\tilde{n}(r,\theta) = \tilde{n}(r)\cos\theta = \tilde{n}\cos\theta$, we can get

$$\frac{\tilde{n}}{n_i} > \frac{1}{2\sqrt{\varepsilon}}\frac{\tau_s}{\tau_{ii}}. \tag{19}$$

The above inequality (19) indicates that the criterion of the poloidal density asymmetry for destabilizing the poloidal plasma spin-up is dependent on the ratio of the electron slowing down time to the ion-ion collision time. Because of $\frac{\tau_s}{\tau_{ii}} = 42.5\frac{z_i^2 n_i}{n_e}(\frac{m_i}{m_p})^{0.5}(\frac{T_e}{T_i})^{1.5}$, the ion species of plasma and the difference between the ion temperature and the electron temperature would affect the magnitude of the threshold of the density for destablizing the plasma spin-up. From the data calculated for D(H) minority fundamental heating, the absorbed rf power density being 0.6 w/cm^3 can make the resonant ion temperature $(T_i = \int d^3v f m v^2 / 2)$ over 40 keV. This means τ_s / τ_{ii} less than 0.19. From Eq. (7) we know that if $\tilde{n} / n$ is larger than 0.2, the poloidal plasma rotation will be destabilized. This amount of the poloidal density inhomogeneity can be produced only when the absorbed rf power density over 0.4 w/cm^3. It is achievable for present tokamaks. Therefore present rf power level can produce the density asymmetry required for destabilizing the poloidal plasma rotation.

Substituting Eq. (12) into Eq. (19), we can obtain the rf power threshold for excitation of the poloidal rotation, above which the anomalous Stringer spin-up would occur. In the high rf power limit in the ICRF fundamental minority heating, we get

$$\left(\frac{(2+3\xi)}{3.7[\frac{2A^{1/2}}{n_e}\sum_j n_j z_j^2]^{2/3}} - 1\right) = \frac{\chi}{2\varepsilon^{3/2}}\frac{\tau_s}{\tau_{ii}}. \tag{20}$$

Therefore, we have

$$< P >= \frac{n_i T_e}{\tau_s}[3.7(\frac{2A^{1/2}}{n_e}\sum_j n_j z_j^2)^{2/3}(\frac{\chi}{2\varepsilon^{3/2}}\frac{\tau_s}{\tau_{ii}}+1)-2], \tag{21}$$

where, $\chi = 1+\frac{q_i^2 n_0}{e^2 T_{//}}(\frac{z_i^2 n_{i0}}{T_i}+\frac{z_I^2 n_{i0}}{T_I}+\frac{n_e}{T_e})^{-1}$.

Because the resonance layer is characterized by $\omega = l\omega_{ci} = l\frac{eB_0}{m}(1-\varepsilon\cos\theta)$, where l is the harmonic number, it can be controlled by changing the toroidal magnetic field. Therefore, the location of sheared flow can be actively controlled. ICRF heating can also be used to induced plasma rotation in the core region, not only in the edge region. Thus, it is more practical and flexible.

5. Summary and discussion

Improved confinement operations are necessary in fusion reactors. Inducing the plasma rotation actively seems to be a reasonable choice for achievement of improved confinement. In this paper, we have proposed a new approach to produce the poloidal rotation by means of strong rf cyclotron heating in tokamak plasmas. This approach is actively controllable and available in core plasma. For simplicity, we have focused our attentions to the case of ICRF heating in a tokamak with concentric circular magnetic surfaces. With the increase in rf power, more and more particles are localized in the outer of torus. Hence, the poloidal density asymmetry is significantly enhanced by rf waves. When the magnitude of the density asymmetry reaches to a certain value, that is large enough to overcome the magnetic pump damping, it will trigger an anomalous spontaneous Stringer spin-up. Then, a poloidally sheared flow is produced. This is the physical mechanism of the present rotation production by means of rf cyclotron heating.

The poloidal density asymmetry due to the resonance localization has been studied analytically and numerically by using the Fokker-Planck code. A criterion upon the density asymmetry for triggering the poloidal plasma rotation has been obtained. The present rf power level has been found to be able to produce the density asymmetry required for destabilizing the poloidal plasma rotation. The expression for the rf power threshold for excitation of the poloidal rotation in the high rf power limit in the ICRF fundamental minority heating has also been presented.

In the present mechanism of production of the plasma rotation, the driving force originates from the rf power, not from the rf momentum input nor pondermative force[10-12]. So, the method can be extended to the case of other rf cyclotron heating such as ECRH which has the

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The H-mode operations have been performed with ICRF heating alone in many tokamaks, such as ASDEX[30], JET[31]. The rf power thresholds have been shown to be independent of the toroidal magnetic field and plasma current, and hardly dependent on the plasma density. Thus we could conclude that during strong rf cyclotron heating in tokamak plasmas, H-mode would be achieved in the following way; rf wave heating with high power would induce a significantly anisotropic distribution. This anisotropy enhances poloidal inhomogeneity of particle density on the magnetic surface. When this inhomogeneity reaches a certain value, it would make the plasma spin-up to trigger the H-mode. Further works to support this assumption are left for future study.

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