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TOKAMAKS

A Possible Way to Reverse the Impurity Influx by Alfvén Heating in a Tokamak Plasma

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Abstract—A study is made of the effect of local heating on impurity fluxes in the Pfirsch–Schlüter regime. When the effect of the thermoforce on impurity ions is taken into consideration, the impurity flux can be reversed by heating the impurities. This concept may be implemented in experiments on Alfvén heating of plasmas in tokamaks. The RF heating power required to reverse the impurity influx is estimated. © *2002 MAIK "Nauka/Interperiodica".*

1. INTRODUCTION

The development of various methods for reducing the impurity content in the plasma is motivated by the experimentally observed degradation of plasma parameters in the presence of heavy impurities. The impurity ions can be affected by both passive (divertors) and active methods. Among the active methods, which have begun to be developed in [1, 2], are the working-gas puffing (particle source); the heating of the main plasma component (energy source); and the transfer of momentum, e.g., from injected high-energy particles to the plasma. However, in [1, 2] (as well as in more recent papers), no account was taken of a weak thermal force acting on impurity ions, and, accordingly, the heat sources acting on impurities were excluded from consideration.

The distinguishing feature of active methods for affecting impurity fluxes is the need to create a particle (or momentum, or energy) source that is asymmetric with respect to the equatorial plane of the tokamak. While it is fairly easy to design particle and momentum asymmetric sources, a means of creating an asymmetric energy source is still lacking. On the other hand, as early as 1975, Messiaen *et al.* [3] pointed out that, in the ion cyclotron frequency range, the waves with $m = +1$ and $m = -1$ (where m is the azimuthal wavenumber) propagate in different ways. This effect stems from the plasma gyrotropy, and, in the model of a plasma cylinder in a longitudinal magnetic field, it is independent of the sign of the axial wavenumber. However, in a helical magnetic field, it breaks the axial symmetry of the plasma. Jaeger *et al.* [4] proposed to use this effect to control the current profile. Craddock and Diamond [5] showed that Alfvén heating can be used to generate a sheared electric field in a plasma in order to suppress edge plasma turbulence. In this paper, we investigate the possibility of reversing the impurity influx by means of Alfvén heating of a plasma in the Pfirsch–Schlüter regime. The analysis is carried out with allowance for the thermal force acting on impurities.

2. ASYMMETRIC HEATING OF IMPURITY IONS

With allowance for the thermal force acting on impurity ions, the expression for the impurity flux averaged over a magnetic surface in a stellarator was derived in [6]. Following [1, 2] and taking into account the results obtained in [6], we arrive at the following expression for the radial impurity flux Γ_I in a tokamak:

$$\Gamma_I = -\Gamma_i/Z_I = \frac{n_i 2q^2\rho_i^2}{Z_I\tau_{iI}T_i}\left[\left(C_1+\frac{C_2^2}{C_3}\right)\left(\frac{1}{n_i}\frac{\partial p_i}{\partial r}-\frac{1}{Z_I n_I}\frac{\partial p_I}{\partial r}\right)\right.$$

$$\left.-\frac{5C_2}{2C_3}\frac{\partial T_i}{\partial r}+\frac{5C_2'}{2C_3'}\frac{\partial T_I}{\partial r}\right]-\frac{n_i q^2\rho_i^2}{Z_I\tau_{iI}T_i}\frac{eB_tR}{c}$$

$$\times\left[\left(C_1+\frac{C_2^2}{C_3}\right)\frac{a_{\tau i}}{n_i}-\frac{C_2}{C_3}\frac{a_{Qi}}{n_iT_i}+\frac{C_2'}{C_3'}\frac{a_{Q1}}{n_IT_I}\right].$$

Here, p_α, $\mathbf{v}_\alpha$, n_α, $Z_\alpha e$, m_α, and T_α are, respectively, the pressure, velocity, density, charge, mass, and temperature of the ions of species α ($\alpha = i, I$); τ_{iI} is the scattering time of the bulk ions by the impurity ions; $q = rB_t/R_0B_p$; a and R_0 are the minor and major radii of the torus; B_t and B_p are the toroidal and poloidal components of the tokamak magnetic field; ρ_i is the Larmor radius of the bulk ions; $a_{Q\alpha}$ and $a_{\tau i}$ are the amplitudes of the components with $\sin\vartheta$ in the Fourier series expansion of the energy source $Q_\alpha(r, \vartheta)$ and the source $\tau_i(r, \vartheta)$ of the main plasma ions; and r and ϑ are the radius and the poloidal angle in the minor cross section (at the inner circumference of the torus, we set $\vartheta = 0$). Hence, under

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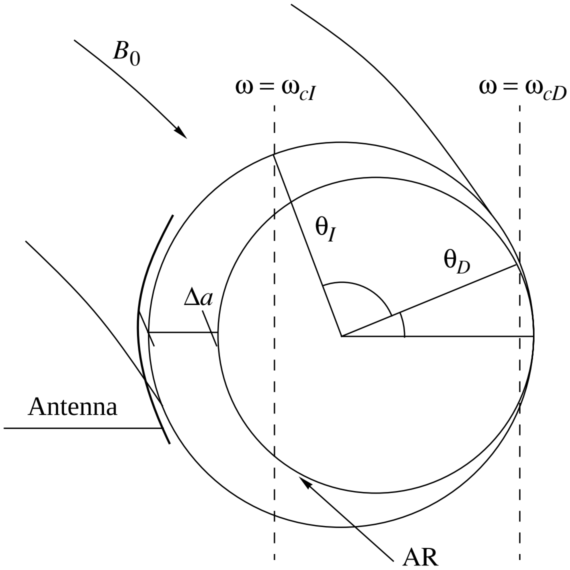

Fig. 1. Minor cross section of a tokamak. Shown are the position of AR and the zones of cyclotron resonances for the main plasma ions ($\omega = \omega_{cD}$) and for the impurity ions ($\omega = \omega_{cI}$).

the condition $B_0 a_{QI}(\bar{r}) < 0$, the impurity influx can be reversed by an asymmetric energy source acting on impurity ions.

Now, we consider whether an asymmetric energy source of this kind can be realized through Alfvén heating of a plasma. Let an antenna exciting the wave magnetic field $B_\parallel$ (which is parallel to the steady-state magnetic field) be positioned at the inner side of the torus. In what follows, we will be interested in the heating of a deuterium plasma. The radial dependence of the cyclotron frequency of deuterium ions is described by the expression $\omega_{cD}(R) = \omega_{cD}(R_0)(1 - \varepsilon_t \cos\vartheta)$. We choose the oscillator frequency ω to be $\omega = \omega_{cD}(R)$ at the outer circumference of the torus (at $\vartheta = \pm\vartheta_D$, where $\cos\vartheta_D = \dfrac{R_0}{a}\dfrac{\omega_{cD}(R_0) - \omega}{\omega_{cD}(R_0)}$). For impurity ions, we have $Z_I/M_I < 1/2$. Consequently, relative to the line $\omega = \omega_{cD}(R)$, the region of cyclotron resonance for impurity ions occurs at $\vartheta = \pm\vartheta_I$ in the inner part of the minor cross section of the torus (Fig. 1). For $\omega < \omega_{ci}$, the local Alfvén resonance (AR) zone (Fig. 1) arises in a tokamak on the inner side with respect to the line $\omega = \omega_{ci}$. In a plasma of sufficiently high density, the local AR zone is located in a narrow region of width Δa at the periphery of the plasma column. Since $n_I \ll k_\parallel \rho_{LI} n_e$, the impurities do not affect the propagation of fast Alfvén waves. From the condition $B_p/B_t = \varepsilon_t/q \ll 1$, we have $B \approx B_t$; consequently, the position of the cyclotron resonance is insensitive to the poloidal magnetic field. The AR position is determined by the expression $\varepsilon_1(r, \vartheta) = N_\parallel^2$, where $\varepsilon_1 = \dfrac{\omega_{pD}^2(r)}{2\omega_{cD}^2\varepsilon_t(\cos\vartheta_D - \cos\vartheta)}$, $N_\parallel = \dfrac{c}{R\omega}(l + m/q)$, and l and m are the toroidal and poloidal wavenumbers. In our analysis, we do not specify how the field excited by an antenna depends on the poloidal angle. Nevertheless, the bulk of the RF power fed into the plasma is, as a rule, carried by the waves with $|m| \le 2$ (see, e.g., [7]). Since $l \ge 10$, we have $l \gg m/q$, so that the poloidal magnetic field has essentially no impact on the AR position. For typical central plasma densities (10^{13}–10^{14} cm^{-3}), the AR occurs at the plasma periphery in a narrow region with the relative width $\Delta a/a \sim \varepsilon_t N_\parallel^2/N_A^2(0)$ (where $N_A^2 = \omega_{pD}^2/\omega_{cD}^2$). In the hydrodynamic approximation, we set $\varepsilon_1/\varepsilon_3 = 0$ in Maxwell's equations for a tokamak plasma with a zero rotational transform (at this point, the effect of the rotational transform can be neglected). As a result, we arrive at the equation

$$\frac{\omega^2}{c^2}B_\parallel - \left(\nabla \times \left[\frac{i\varepsilon_2}{\varepsilon^2 - \varepsilon_2^2}\frac{1}{R}\mathbf{\nabla}(RB_\parallel)\right]\right)_\parallel + R\nabla \cdot \left[\frac{\varepsilon_1 - N_\parallel^2}{\varepsilon^2 - \varepsilon_2^2}\frac{1}{R^2}\mathbf{\nabla}(RB_\parallel)\right] = 0, \quad (1)$$

where

$$\varepsilon = \varepsilon_1 - N_\parallel^2,\ \varepsilon_1 = 1 + \sum_i \frac{\omega_{pi}^2(r)}{\omega_{ci}^2(r,\vartheta) - \omega^2},$$

$$\varepsilon_2 = \sum_i \frac{\omega\omega_{pi}^2(r)}{\omega_{ci}(r,\vartheta)[\omega_{ci}^2(r,\vartheta) - \omega^2]},\ \text{and}\ N_\parallel = \frac{cl}{\omega R}.$$

Recall that $\Delta a/a$ is a small parameter in our problem. In order to solve Eq. (1) with the help of this parameter, we multiply Eq. (1) by $\exp(-im\vartheta)$ and integrate over the region S, whose width is $\Delta r \le \Delta a$ and which is bounded by contours L_1 and L_2, as is shown in Fig. 2. Using the Stokes and Gauss theorems, we convert Eq. (1) into the form

$$\frac{\partial}{\partial r}[R(r,\vartheta)B_\parallel(r,\vartheta)] - \frac{i\varepsilon_2(r,\vartheta)}{r\varepsilon(r,\vartheta)}\frac{\partial}{\partial\vartheta}[R(r,\vartheta)B_\parallel(r,\vartheta)]$$

$$= \frac{\varepsilon^2(r,\vartheta) - \varepsilon_2^2(r,\vartheta)}{\varepsilon(r,\vartheta)(N_\parallel^2 - 1)}\frac{aR(r,\vartheta)}{rR(a,\vartheta)}\frac{\partial}{\partial a}[R(a,\vartheta)B_\parallel(a,\vartheta)]$$

$$- \frac{R(r,\vartheta)}{r}\frac{\varepsilon^2(r,\vartheta) - \varepsilon_2^2(r,\vartheta)}{\varepsilon(r,\vartheta)}\left\{\frac{\omega^2}{c^2}\int_a^r r'B_\parallel(r',\vartheta)dr' \quad (2)\right.$$

$$+\frac{\partial}{\partial\vartheta}\int_a^r \frac{dr'}{R(r',\vartheta)} \frac{i\varepsilon_2(r',\vartheta)\frac{\partial}{\partial r'}[R(r',\vartheta)B_{\parallel}(r',\vartheta)]}{\varepsilon^2(r',\vartheta)-\varepsilon_2^2(r',\vartheta)}$$

$$\left.+\frac{\partial}{\partial\vartheta}\int_a^r \frac{dr'}{r'R(r',\vartheta)} \frac{\varepsilon(r',\vartheta)\frac{\partial}{\partial\vartheta}[R(r',\vartheta)B_{\parallel}(r',\vartheta)]}{\varepsilon^2(r',\vartheta)-\varepsilon_2^2(r',\vartheta)}\right\}.$$

Here, we have introduced the notation $\frac{\partial}{\partial r}R(r,\vartheta)B_{\parallel}(r,\vartheta)\Big|_{r=a} = \frac{\partial}{\partial a}R(a,\vartheta)B_{\parallel}(a,\vartheta)$. Now, Eq. (2) can be expanded in powers of the small parameter $\Delta a/a$, and the problem can be solved by the narrow layer method. This method was originally developed in [8] for the problem of surface waves. In [9], it was applied to the study of an AR in a one-dimensional plasma, and, in [10], it was generalized to a two-dimensional (r, ϑ) case. We assume that the quantities $B_{\parallel}(a, \vartheta)$ and $\frac{\partial}{\partial a}B_{\parallel}(a,\vartheta)$ are specified at the plasma boundary and use the small parameter $\Delta a \ll a, a/m, 1/k_{\parallel}$. Taking into account the relationships $\frac{\partial\varepsilon}{\partial\vartheta}/\frac{\partial\varepsilon}{\partial r} \sim \frac{\partial\varepsilon_2}{\partial\vartheta}/\frac{\partial\varepsilon_2}{\partial r} \sim \frac{\Delta a}{a} \ll 1$, we expand the wave magnetic field $B_{\parallel}$ in a power series in the small parameter $\Delta a/a$, $B_{\parallel} = B_{\parallel 0} + B_{\parallel 1} + \ldots$, to obtain

$$\frac{i\varepsilon_2}{r\varepsilon}\frac{\partial(RB_{\parallel 0})}{\partial\vartheta} - \frac{\partial(RB_{\parallel 0})}{\partial r} = \frac{\varepsilon^2-\varepsilon_2^2}{\varepsilon(N_{\parallel}^2-1)}\frac{\partial R(a,\vartheta)B_{\parallel}(a,\vartheta)}{\partial a}, \tag{3}$$

$$\frac{i\varepsilon_2}{r\varepsilon}\frac{\partial(RB_{\parallel 1})}{\partial\vartheta} - \frac{\partial(RB_{\parallel 1})}{\partial r} = -\frac{R}{r}\frac{\varepsilon^2-\varepsilon_2^2}{\varepsilon}\frac{\partial}{\partial\vartheta}$$
$$\times\left[\frac{1}{R}\left(\int_a^r \frac{i\varepsilon_2}{\varepsilon^2-\varepsilon_2^2}\frac{\partial RB_{\parallel 0}}{\partial r'}dr' + \frac{1}{a}\int_a^r \frac{\varepsilon}{\varepsilon^2-\varepsilon_2^2}\frac{\partial RB_{\parallel 0}}{\partial\vartheta}dr'\right)\right]. \tag{4}$$

Equations (3) and (4) constitute a boundary-value problem of the first kind, which can be solved by the method of characteristics. The equations of characteristics have the form

$$\frac{d\vartheta(r)}{dr} = -i\frac{\varepsilon_2[r,\vartheta(r)]}{r\varepsilon[r,\vartheta(r)]}, \tag{5}$$

$$\frac{d}{dr}RB_{\parallel 0}[r,\vartheta(r)] = \frac{\varepsilon^2[r,\vartheta(r)]-\varepsilon_2^2[r,\vartheta(r)]}{\varepsilon[r,\vartheta(r)](1-N_{\parallel}^2)}\frac{\partial}{\partial a}RB_{\parallel}[a,\vartheta(r)]. \tag{6}$$

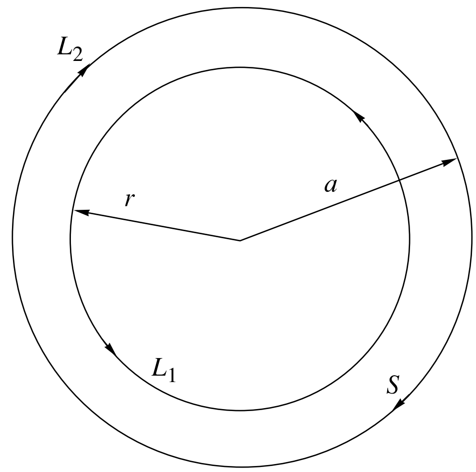

Fig. 2. Domain of integration S for Eq. (8) and contours L_1 and L_2.

An approximate solution to Eq. (5) can be written in the form

$$\Theta(\vartheta_0, a; r) \approx \vartheta_0 - i\int_a^r \frac{\varepsilon_2(r',\vartheta_0)}{r'\varepsilon(r',\vartheta_0)}dr'. \tag{7}$$

Hence, we have

$$\vartheta_0(r,\vartheta) = \Theta(\vartheta, r; a) \quad\text{and}\quad \vartheta'(r,\vartheta) = \Theta(\vartheta, r; r'). \tag{8}$$

Substituting solution (7) into Eq. (6) and using relationship (8), we obtain

$$R(r,\vartheta)B_{\parallel 0}(r,\vartheta) = R(a,\vartheta_0)B_{\parallel}(a,\vartheta_0)$$
$$+\int_a^r \frac{\varepsilon^2(r',\vartheta')-\varepsilon_2^2(r',\vartheta')}{\varepsilon(r',\vartheta')(1-N_{\parallel}^2)}\frac{\partial}{\partial a}R(a,\vartheta')B_{\parallel}(a,\vartheta')dr'.$$

Now, we expand the first term on the right-hand side of this equation in a series in $\vartheta - \vartheta_0$ and retain only the first term in the expansion. Setting $\vartheta' \approx \vartheta$ in the second term and introducing the notation $R(a,\vartheta)B_{\parallel}(a,\vartheta) = R_aB_{\parallel a}$, we arrive at the following equations for the wave magnetic field:

$$R(r,\vartheta)B_{\parallel 0}(r,\vartheta) = R_aB_{\parallel a} + \delta(RB),$$

$$\delta(RB) = i\frac{\partial}{\partial\vartheta}(R_aB_{\parallel a})\int_a^r \frac{\varepsilon_2(r',\vartheta)}{r'\varepsilon(r',\vartheta)}dr' \tag{9}$$

$$+\frac{\partial}{\partial a}(R_aB_{\parallel a})\int_a^r \frac{\varepsilon^2(r',\vartheta)-\varepsilon_2^2(r',\vartheta)}{(1-N_{\parallel}^2)\varepsilon(r',\vartheta)}dr'.$$

Equations (9) describe the behavior of the field in a narrow layer between the plasma boundary and the vicinity of the AR. Note that, formally, the integrals in Eqs. (9) have a pole, $\varepsilon(r',\vartheta) = 0$, at the AR point. However, when dissipative effects (e.g., Coulomb collisions) are taken into account, the pole actually lies in the region of complex values of r'.

Using Eqs. (9), we can calculate the poloidal RF power flux:

$$S_\vartheta = \frac{c^2}{8\pi\omega R}\mathrm{Re}\left\{\frac{B^*_{\|a}}{\varepsilon}\left[\frac{\varepsilon_2}{N_\|^2-1}\frac{\partial}{\partial a}(R_a B_{\|a}) + \frac{i}{a}\frac{\partial}{\partial\vartheta}(R_a B_{\|a})\right]\right\}. \tag{10}$$

Let us analyze this expression, assuming that $B_\|(a, \vartheta) \sim \exp(im\vartheta)$ at the plasma boundary. Since the quantity ε_2 is proportional to the plasma density and, accordingly, vanishes at the plasma periphery, the second term is dominant near the plasma boundary and we have $S_\vartheta \approx \frac{c^2}{8\pi\omega}\frac{m}{a}\frac{|B_{\|a}|^2}{N_\|^2-1}$. The waves with $m < 0$ carry the RF power in a clockwise direction, and the direction of the power carried by the waves with $m > 0$ is counterclockwise (Fig. 1). As the wave penetrates deeper into the plasma, the first term, which accounts for the plasma gyrotropy, plays an increasingly important role. Taking into account the inequality $\mathrm{Re}\left[B^*_{\|a}\frac{\partial}{\partial a}(R_a B_{\|a})\right] < 0$, we can see that, in this case, the power flux in a counterclockwise direction increases and the power flux in a clockwise direction decreases. This leads to an up/down asymmetric RF power flux in the minor cross section of the torus. The radial power flux can be represented as $S_r = S_{r0} + \delta S_r$, where

$$S_{r0} = \frac{c^2}{8\pi\omega R}\mathrm{Im}\left[\frac{B^*_{\|a}}{N_\|^2-1}\frac{\partial}{\partial a}(R_a B_{\|a})\right],$$

$$\delta S_r = \frac{c^2\omega}{4R^2}\left(\left.\frac{d\omega_{pi}^2}{dr}\right|_{r=a}\right)^{-1}\varepsilon_t(\cos\vartheta_D - \cos\vartheta) \tag{11}$$

$$\times\left|\frac{\omega}{\omega_{ci}}\frac{\partial}{\partial a}(R_a B_{\|a}) - \frac{m}{a}R_a B_{\|a}\right|^2.$$

The term δS_r describes the power deposited in the AR zone. Expressions (11) show that, in the lower part of the torus ($m > 0$), the RF power is absorbed more intensely than in the upper part ($m < 0$). In the vicinity of the AR, a fast magnetosonic (FMS) wave converts into a small-scale kinetic wave (KW). When the AR occurs at the plasma periphery, the ion contribution (on the order of $\omega^2\rho_{Li}^2/c^2$) to the dispersion of the KW can be neglected in comparison with the electron contribution (which is on the order of ε_3^{-1}, where $\varepsilon_3 = \frac{\omega_{pe}^2}{k_\|^2 v_{Te}^2}[1 + i\sqrt{\pi}z_e W(z_e)]$ with $z_e = \omega/\sqrt{2}k_\| v_{Te}$). For the parameters of the present-day tokamaks, we have $\varepsilon_3\omega^2\rho_{Li}^2/c^2 \sim \frac{T_i}{T_e}\varepsilon_t \ll 1$ for $z_e \ll 1$ and $\varepsilon_3\omega^2\rho_{Li}^2/c^2 \sim \frac{v_{Te}^2}{c^2}\varepsilon_t N_\|^2 \ll 1$ for $z_e \gg 1$. The KW propagates along the magnetic field lines toward the inner side of the torus, deviating only slightly from the magnetic surface toward lower plasma densities. Kinetic waves that originate in the sector $-\vartheta_I < \vartheta < \vartheta_I$ will reach the zone of cyclotron resonance for impurity ions and will be completely absorbed there. Hence, we have shown how asymmetric heating of impurities can be realized in tokamak experiments.

3. DISCUSSION OF THE RESULTS OBTAINED

First, we analyze how the process in question is affected by heat exchange between the impurities and the main plasma ions. In [2], it was noted that, with allowance for heat exchange, the change in the impurity temperature (including the change associated with the energy source) decreases by a factor of f, where $f = 1 + \frac{3m_i n_i q^2 R^2}{39 n_I \tau_{iI}\tau_{II} T_I}$. Setting $n_i(a) = 10^{12}$ cm^{-3}, $n_I/n_i = 2\times10^{-3}$, $m_I \approx 60m_p$, $T_I(a) = T_i(a) = 50$ eV, $Z_I = 4$, $q = 3$, and $R_0 = 150$ cm, we obtain $\tau_{iI} = 2.6\times10^{-3}$ s, $\tau_{II} = 1.3\times10^{-3}$ s, and $f = 2$. We thus can conclude that this effect is of a quantitative (rather than qualitative) character.

Second, we discuss another mechanism for the absorption of a fast Alfvén wave by impurity ions. The impurities make a small contribution to the tensor elements ε_1 and ε_2. The cyclotron absorption of a fast Alfvén wave by impurity ions is described by the imaginary part of this contribution. However, the fraction of the RF power that is absorbed by impurities is small in comparison with that transferred to the KW: $P_{ab}/P_{\mathrm{KW}} \sim n_I Z_I/n_i \ll 1$.

For a KW, the optical thickness τ of the zone of cyclotron resonance for impurity ions is described by the expression

$$\tau \sim \mathrm{Im}(N_\|)\Delta l \sim 9\sqrt{\frac{\pi}{2}}\frac{\omega_{pI}^2}{\omega_{cI}}\frac{Rq}{\varepsilon_t N_\| c},$$

where Δl is the length of the resonance zone along the magnetic field lines. Setting $k_\| = 0.1$ cm^{-1}, we obtain $\tau \geq 1$. Note that, in the resonance zone, we have $z_e \sim 6$, in which case Landau damping of the KW is negligible, $\tau_L \sim k_\| Rq\,\mathrm{Im}\varepsilon_3/|\varepsilon_3| \ll 1$. Consequently, the KW is completely absorbed by impurity ions. Denoting the total RF power fed into the plasma by P_{tot}, we can see that the power transferred to the KW is equal to $\frac{\varepsilon_t N_\|^2}{N_A^2(0)}P_{\mathrm{tot}}$, where $N_A^2(0) = \omega_{pi}^2(0)/\omega_{ci}^2(0)$. In order to obtain the total power absorbed by impurity ions, we must integrate expressions (11) over $\vartheta_D - \vartheta_I$. Note that the optimum position of cyclotron resonance for impurities is between $\vartheta_I = 0.88\pi$ for $Z_I/M_I = 1/4$ and $\vartheta_I = 2\pi/3$ for $Z_I/M_I = 1/3$. Consequently, the quantity Q_I' is equal to

$$Q_I' \approx \frac{\Delta\vartheta_I}{\pi}\frac{\sin\vartheta_I[(\vartheta_I - \vartheta_D)\cos\vartheta_D - \sin\vartheta_I + \sin\vartheta_D]}{(\pi - \vartheta_D)\cos\vartheta_D + \sin\vartheta_D} \times K_{as}\frac{\varepsilon_t N_\parallel^2}{N_A^2(0)}\frac{P_{\text{tot}}}{V}, \tag{12}$$

where the quantity $\Delta\vartheta_I = 3k_\parallel\rho_{LI}/\varepsilon_t$, $K_{as} \sim \Delta a/a$ describes the up/down asymmetry of the RF power distribution, $V = 2\pi Ra\Delta\vartheta_I\Delta r$ is the plasma volume in which the impurity ions absorb RF power, and Δr is the radial thickness of the absorption region. For estimates, we adopt $\Delta r \sim \Delta a$. Then, a necessary condition for reversing the direction of the impurity influx has the form

$$Q_I' \approx 2\frac{n_i(0)T_i(0)}{aR}\frac{cT_I(a)}{eB_t}\frac{n_I}{n_i}. \tag{13}$$

Substituting condition (13) into expression (12) yields

$$P_{\text{tot}} \approx 7\frac{n_i(0)T_i(0)}{K_{as}}\frac{cT_I(a)}{eB_t}\frac{n_I}{n_i}a.$$

In deriving this estimate, we took into account the fact that the doubled reciprocal of the second fraction in expression (12) is approximately equal to seven. For $n_i(0) = 10^{14}$ cm^{-3}, $n_I/n_i = 2 \times 10^{-3}$, $T_i(0) = 2$ keV, $B_t = 1$ T, $K_{as} = 0.1$, and $a = 50$ cm, we obtain $P_{\text{tot}} \approx 200$ kW.

4. CONCLUSION

Our analysis is based on the assumption that, at the periphery of the plasma column, impurity ions are transported in the Pfirsch–Schlüter regime. It has been shown that, when the small thermoforce affecting the impurities is taken into account, the impurity influx can be reversed by asymmetric (up-down) heating of the impurity ions. To achieve this, we propose to excite RF fields with frequencies lower than the cyclotron frequency of the main (deuterium) plasma ions on the inner side of the torus. We have shown that, after an FMS wave converts into a KW in the vicinity of the AR, the RF power is absorbed by impurity ions. For a medium-size tokamak, the RF power required to reverse the impurity influx is approximately equal to 200 kW. The efficiency of the method proposed here can be increased by displacing the antenna in the poloidal direction. The question of whether the method will prove feasible when heavy impurities are transported in the "plateau" regime or when their transport is anomalous requires a separate study.

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REFERENCES

1. P. H. Rutherford, Phys. Fluids **17**, 1782 (1974).
2. K. T. Burrell, Phys. Fluids **19**, 401 (1976).
3. A. M. Messiaen, P. E. Vandenplas, R. R. Weynants, and R. Koch, Nucl. Fusion **15**, 75 (1975).
4. E. F. Jaeger, M. D. Carter, L. A. Berry, *et al.*, in *Proceedings of the 24th European Conference on Controlled Fusion and Plasma Physics, 1997*, Rep. P2.122.
5. G. G. Craddock and P. H. Diamond, Phys. Rev. Lett. **67**, 1535 (1991).
6. D. L. Grekov, Fiz. Plazmy **27**, 1070 (2001) [Plasma Phys. Rep. **27**, 1011 (2001)].
7. A. V. Longinov and V. A. Tsurikov, Fiz. Plazmy **16**, 39 (1990) [Sov. J. Plasma Phys. **16**, 22 (1990)].
8. K. N. Stepanov, Zh. Tekh. Fiz. **35**, 1002 (1965) [Sov. Phys. Tech. Phys. **10**, 773 (1965)].
9. D. L. Grekov, K. N. Stepanov, and J. A. Tataronis, Fiz. Plazmy **7**, 752 (1981) [Sov. J. Plasma Phys. **7**, 411 (1981)].
10. S. V. Kasilov, Author's Abstract of Candidate's Dissertation in Physics and Mathematics (Kharkov, 1989).

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