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Alfvén wave current drive

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Abstract

The resonant and nonresonant current drive produced by the Alfvén and fast waves are discussed both for collision and collisionless magnetized plasmas. The magneto-hydrodynamics approach is used to study collision plasma regimes, and the purely kinetic model, with the Landau and transit time magnetic pumping (TTMP) damping, is considered in the weak collision plasma case. The general form of the time-averaged ponderomotive force (and current drive as a result) generated by waves in magnetized plasma configurations is presented. It has been shown that standing Alfvén waves can drive non-inductive and nonresonant current in a magnetized plasma through Landau and TTMP damping. If all oscillations of the solar corona loops are related to standing Alfvén waves, which can introduce the density of dissipated wave power that is necessary to keep the plasma hot, the axial nonresonant current drive has been estimated to be enough for equilibrium and stability of corona loops. Experimental evidence of the Alfvén wave current drive in tokamaks and stellarators is demonstrated.

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1 Introduction

Existence of a current along magnetic field lines is necessary for an equilibrium and a stability of tokamak plasmas in laboratory experiments where this current is created by an electric field induced by a transformer. In space plasmas, for example, the Solar corona loop configuration (and stability) may be explained by this current too, and magnetic field aligned current can be observed in the auroral zone of the Earth magnetosphere. A tokamak with induction current is a pulsed system, which gives rise to excessive stress, that was one of the reasons for developing stationary tokamak-reactor studies. In the course of this study, the idea of current drive by radio-frequency (RF) waves was proposed by Wort [1].

Nowadays, RF plasma heating and current drive studies are very important parts of experimental programs on all large tokamaks (JET, D-III, TFTR, JT-60, Tor-Supra) (see, for example, [2, 3]). These experiments have shown that current drive can be produced by different kinds of waves (electron cyclotron, lower hybrid and fast waves). Alfvén wave heating and current drive are attractive for tokamaks and stellarators (see, for example, [4]) because of the availability of high power RF generators and the simplicity of wave launching structures. For the first time the Alfvén wave current drive was shown in the R-O stellarator experiments (see [5]). Recently the Alfvén current drive has also been demonstrated in the Phaedrus-T tokamak [6, 7]. In the nearest future current drive will be under study in TCA/Br (the Brazilian tokamak with Alfvén wave excitation [8]).

Current drive effects may exist in the space magnetized plasma. In solar corona loop conditions mentioned above, there is no obvious source of induction current to support the poloidal magnetic field during the long period of the corona loop life-time. The electron-ion collision diffusion time of the poloidal magnetic field is smaller by a factor of $a^2/R^2 \approx 10^{-6}$ than the diffusion time of the axial magnetic field, due to different scales in minor radius (a) and to the length of corona loops (πR). And there is a question about the source of the poloidal magnetic field in corona loops. One possible explanation of this phenomena may be a current drive [9] induced during Alfvén wave heating. The idea of the Alfvén wave corona heating has already been explored by many authors (see, for example, Refs. [10, 11, 12]). In the Earth magnetosphere physics, there are satellite instrumental measurements of the Alfvén wave excitation on the morning side of the Earth [13], which can be a source of the auroral electron precipitation and current drive [14, 15]. Alfvén wave current drive has been proposed in the extragalactic jets too [16]. So these were the substantial reasons to develop theoretical and computer methods to study the Alfvén wave current drive in magnetized plasmas.

The theoretical investigation of different aspects of the current drive has a long enough history, which is reviewed in Ref. [17]. In a part specifically related to the Alfvén wave current drive, this problem is observed in Ref. [4]. The simplest physical model of current drive is the transformation of the dissipated travelling wave momentum into plasma electron momentum, which is lost due to electron-ion collisions. Finally, in steady-state conditions, the momentum balance leads to the value of the current drive density j_{cd} , which is governed by the equation:

$$j_{cd} = eP/(m\nu_{ei}v_{ph}) \quad (1)$$

where e and m are respectively the charge and mass of electrons, ν_{ei} is the electron-ion

collision frequency, P is the density of dissipated power and v_{ph} is the phase velocity. This simple formula is valid for the phase velocity less than electron thermal velocity v_T . The result related to the whole range of phase velocities was computed in Ref. [18] where high efficiency current drive was demonstrated for the Alfvén waves with small phase velocity, $v_{ph} \ll v_T$, as well as large phase velocity, $v_{ph} \gg v_T$.

Now, it is well known that there are significant difficulties to use the Alfvén current drive in a tokamak reactor. There was an opinion, discussed in Refs.[17, 19, 20], that the efficiency of this current drive drops as a result of a wave pulse absorption by the trapped particles and because of the additional friction between untrapped electrons, which carry a current, and trapped electrons. This opinion is valid if current drive calculations takes into account only the travelling wave momentum imparted in to untrapped electrons (resonant model of the Alfvén wave current drive) discussed above. From another point of view (see [21]), in the weak collision limit (banana regime) the wave momentum, which is imparted by waves into trapped electrons, should be conserved as a canonical toroidal momentum of electrons:

$$m_e v_\zeta [1 + (r/R_0 \cos(\theta))] - \frac{e}{c} \int_0^r B_\theta dr = \text{const} \quad (2)$$

where ζ, θ are toroidal and poloidal angles, r and R_0 are minor and major radii of tokamak magnetic surfaces, B_θ is the poloidal component of a tokamak magnetic field, and c is the light velocity. As a result of this conservation law, the banana electrons suffer an inward Ware pinch as they gain momentum from the wave. In steady state conditions, due to collisional diffusion (inverse Ware pinch via collisions), this momentum can partially (approximately 50% of the imparted wave pulse) contribute to the untrapped electron current. So the efficiency of the Alfvén wave current drive should be high enough in a comparison with the current drive efficiency, which was calculated in Refs.[17, 19, 20]. Up to now, for tokamak plasmas, this question stays open due to difficult mathematical problems, and the complete physical model of the current drive has not yet been supported by the general solution of transport equations.

New hopes appeared, which was related to the possibility of increasing the current drive efficiency through helicity injection, as proposed by Ohkawa [22], and this idea was developed in Refs.[23]-[26]. In magnetized plasmas, helicity injection current can be maintained by means of ponderomotive forces, depending on the radio frequency (RF) field amplitude gradients. A general form of the ponderomotive forces in two-fluid MHD model of magnetized plasmas has been derived in Refs.[27] and [28] for a plane geometry. Some ponderomotive force related to dynamo effects was discussed in Ref. [29]. The general view of forces, which were produced by RF waves, was reviewed in Ref. [30] for cylindrical plasmas, and these driving forces in a general geometry were later calculated in Ref. [31].

Because collisionless wave dissipation stays out of MHD analysis, it is important to develop a kinetic approach of the current drive theory. The current drive produced by the Alfvén and fast waves was calculated in Refs.[4, 32, 33], where the kinetic approach was considered, to study time-averaged plasma currents driven by RF waves. The general kinetic study can be developed in the frame of Vlasov - Maxwell set of equations in three stages:

- On the first stage, the linearized Vlasov kinetic equation in an adequate geometry should be solved and the dielectric permeability (or conductivity) tensor can be

calculated. Thus, the connection between wave oscillations of current density and electric field are defined.

- On the next step, taking into account the obtained analytical expressions of the dielectric tensor, the Maxwell equations should be solved, in order to study the wave field structures, and to evaluate the dissipated wave power in complicated magnetic field geometries (for example, in toroidal or dipole configurations).
- On the last step, using the Vlasov equation, which is averaged over wave oscillations and over magnetic surfaces, and taking into account the kinetic collision operator (for example, in Landau form), the current drive can be calculated.

The plan of this paper is as follows. In section 2, we present and discuss the ponderomotive force in order to derive expressions for the plasma current in a cylindrical plasma column. In section 3, we show the direct kinetic evaluation of the axial current and present some results of the Alfvén wave current drive related to the plasma and RF field amplitude gradients. In section 4, the experimental evidence of the Alfvén wave current drive in the P-O stellarator and in the Phaedrus-T tokamak is discussed. Conclusions of this study are summarized in section 5.

2 MHD Ponderomotive forces and current drive

What will be relevant for the objectives of this section is the expression for the general ponderomotive force derived in Ref. [27]. In accordance with the theory of ponderomotive effects in a two-fluid plasma, each field of the plasma medium is expressed as a sum of a quasi-steady component and an RF component. The subscripts "0" and " ω " of currents, density and electromagnetic fields denote the time-averaged and the RF oscillating values, respectively and the subscript α is the index of particle species. Summarizing, the ponderomotive force, which can be proceed from the two fluid MHD equations, has the form,

$$\dot{F}_\alpha = \dot{F}_{n\alpha} + \dot{F}_{d\alpha} + \dot{F}_{p\alpha}.$$

where

$$\dot{F}_{n\alpha} = e_\alpha(n_{\alpha\omega}\dot{E}_\omega)_0, \quad \dot{F}_{d\alpha} = \frac{1}{c}([\dot{j}_{\alpha\omega} \times \dot{B}_\omega])_0, \quad \dot{F}_{p\alpha} = -4\pi\nabla_i \left(\frac{j_{\alpha\omega}^i j_{\alpha\omega}^i}{\omega_{p\alpha}^2} \right)_0. \quad (3)$$

where $n_{\alpha\omega}$ is the fluctuating plasma density, $\dot{j}_{\alpha\omega}$ is the fluctuating current density and e_α is the charge of species α . The vectors $\dot{F}_{D\alpha}$ and $\dot{F}_{n\alpha}$ are related to the dynamo effect and electrostrictive effect forces, that may contribute to current drive by wave helicity injection [22]-[26],[32]. The next part of the total ponderomotive force is $\dot{F}_{p\alpha}$, which is a dynamic pressure force [30] related to the macroscopic oscillating flow velocity of the fluid:

$$\dot{F}_{p\alpha} = -m_\alpha \nabla_i (n_{\alpha 0} V_{\alpha\omega}^i \dot{V}_{\alpha\omega}^i) \quad (4)$$

In Eq.(3) the force $\dot{F}_{p\alpha}$ was expressed in terms of the RF current density, $\dot{j}_{\alpha\omega}$, and the reasonable assumption that the wave phase velocity is much greater than the time averaged fluid velocity is taken into account.

In the presence of the ponderomotive force, $\dot{F}_\alpha$, the average fluid motion of particle species α is governed by the force equation,

$$m_\alpha \frac{\partial n_{\alpha 0} \dot{U}_\alpha}{\partial t} = -\nabla P_{\perp \alpha 0} - \nabla_i [\dot{h}_0^i \dot{h}_0 (P_{\parallel \alpha 0} - P_{\perp \alpha 0})] - m_\alpha \nabla_i (n_{\alpha 0} U_{\alpha 0}^i \dot{U}_{\alpha 0}) + e_\alpha n_{\alpha 0} (\dot{E}_0 + \frac{1}{c} [\dot{U}_\alpha \times \dot{B}_0]) + \dot{R}_0^\alpha + \dot{F}_\alpha. \quad (5)$$

The term $\dot{R}_0^\alpha$ accounts for collision friction and particle sources, $\dot{P}_\alpha$ is the thermal pressure tensor, $\dot{h}_0 = \dot{B}_0/B_0$ and $\dot{U}_\alpha$ is the time-averaged velocity, defined by the expression,

$$n_{\alpha 0} \dot{U}_\alpha = (n_\alpha \dot{V}_\alpha)_0. \quad (6)$$

Note, here, the current associated with a pressure tensor stays out of discussion because of low β approximation.

2.1 RF current drive

As an application of the ponderomotive force theory that we have outlined, we consider in this section an axial plasma current drive produced by RF waves in a cylindrical plasma configuration. We assume a plasma consisting of electrons and charged ions with equilibrium number densities, $n_{e0}(r)$ and $n_{i0}(r)$, respectively. As we wish to model a tokamak plasma, we further assume that the equilibrium magnetic field has two components, a poloidal component, $B_{0\theta}(r)$, and an axial component, $B_{0z}(r)$, both of which depend on r . The axial component of the magnetic field simulates the toroidal component of a tokamak. Superimposed on this configuration is an RF wave that produces a steady current via the effects of a ponderomotive force. Our analysis of this configuration will be done in cylindrical coordinates, (r, θ, z) .

Ponderomotive force effects of RF waves produce steady currents in addition to transport. The axial component of the RF induced electron current density, j_{0z}^{RF} , averaged over the magnetic surfaces of the cylindrical plasma, has been derived in Ref. [28]. It can be expressed in the following form,

$$j_{0z}^{RF} = -\frac{B_{0z}}{e B_0 n_{e0} \eta_{\parallel}} \left(\frac{B_{0\theta}}{B_0} \langle F_{e\theta} \rangle + \frac{B_{0z}}{B_0} \langle F_{ez} \rangle \right), \quad (7)$$

where $\eta_{\parallel}$ represents the plasma resistivity in the direction parallel to the equilibrium field. Our estimates of wave induced transport and current, which we will make in a later section of this paper, will be based on Eqs.(7).

To complete Eqs.(7), it is necessary to express $F_{e\theta}$ and F_{ez} in terms of cylindrical components of the RF wave fields from Eqs.(3). where only the wave electric field and current density appear. We assume a single harmonic wave propagating along the plasma column,

$$\{\vec{E}_\omega(r, \theta, z, t); \vec{j}_\omega(r, \theta, z, t)\} = \{\hat{E}(r); \hat{j}(r)\} e^{i(kz + m\theta - \omega t)}. \quad (8)$$

where m and k are, respectively, the poloidal mode number and axial wavenumber. The wave amplitudes, $\hat{E}(r)$ and $\hat{j}(r)$, are connected by the linearized equations that govern the harmonic wave dynamics of the cylindrical plasma configuration. Substituting Eqs.(8) in Eqs.(3), we find the following expressions for the poloidal and axial components of the average electron ponderomotive force.

$$\langle F_{e\theta} \rangle = \frac{m}{r\omega} P_e - \left\{ \frac{1}{4r^2} \frac{\partial}{\partial r} \left[r^2 \hat{j}_r^* \left(\frac{\hat{E}_\theta}{i\omega} + \frac{4\pi \hat{j}_\theta}{\omega_{pe}^2} \right) \right] + c.c. \right\}, \quad (9)$$

$$\langle F_{ez} \rangle = \frac{k}{\omega} P_e - \left\{ \frac{1}{4r} \frac{\partial}{\partial r} \left[r \hat{j}_r^* \left(\frac{\hat{E}_z}{i\omega} + \frac{4\pi \hat{j}_z}{\omega_{pe}^2} \right) \right] + c.c. \right\}, \quad (10)$$

where *c.c.* designates the complex conjugate of the preceding term and P_e is the time-averaged RF power density absorbed by electrons.

$$P_e \equiv \frac{1}{2} \text{Re} \left(\hat{E} \cdot \hat{j}^* \right), \quad (11)$$

Observe that both $\langle F_{e\theta} \rangle$ and $\langle F_{ez} \rangle$ can be expressed as the sum of two components.

$$\langle F_{e\theta} \rangle = F_{\theta P} + F_{\theta \partial}, \quad \langle F_{ez} \rangle = F_{z P} + F_{z \partial}, \quad (12)$$

where the subscripts P and ∂ label, respectively, force components associated with electron power absorption and radial derivatives, as they appear in Eqs.(10). Both the electrostrictive force, $\vec{F}_{ne}$, and the electron dynamo force, $\vec{F}_{de}$, defined, respectively, by Eqs.(3), contribute to the absorption and derivative forces of Eqs.(12). On the other hand, the dynamic pressure force, defined by Eq.(4) contributes only to the derivative forces in Eqs.(9) and (10).

The gradient forces have an interesting implication regarding the transfer of linear and angular momentum to the plasma. It may be easily shown that the two gradient forces are zero when dissipation is absent. Therefore, nonzero dissipation, such as collisions, Landau damping, wave transformation in a resonance region, etc., is essential for these ponderomotive force components. It may be also shown that the gradient forces may dominate over the momentum absorption forces, $F_{\theta P}$ and $F_{z P}$, in many cases relevant to RF plasma heating in tokamaks and also in low temperature plasma devices used for plasma processing.

Substituting Eqs.(9) and (10) in Eq.(7) we express the induced axial current density in terms of the RF electric field and oscillating current density,

$$j_{0z}^{RF} = -\frac{B_{0z}}{en_{e0}\eta_{||}B_0^2} \left\{ \frac{mB_{0\theta} + krB_{0z}}{r\omega} P_e - \left\{ \frac{B_{0\theta}}{4r^2} \frac{\partial}{\partial r} \left[r^2 \hat{j}_r^* \left(\frac{\hat{E}_\theta}{i\omega} + \frac{4\pi \hat{j}_\theta}{\omega_{pe}^2} \right) \right] + \frac{B_{0z}}{4r} \frac{\partial}{\partial r} \left[r \hat{j}_r^* \left(\frac{\hat{E}_z}{i\omega} + \frac{4\pi \hat{j}_z}{\omega_{pe}^2} \right) \right] + c.c. \right\} \right\}. \quad (13)$$

Equation (13) are composed of two fundamentally different groups of terms, which can be viewed as sources of the induced electron current. The first consists of the single

term that depends directly on the absorbed power, P_e , and therefore is associated with the transfer of momentum from the wave to the electrons. It is the source that is likely dominant in a variety of RF current drive schemes. In addition to wave momentum transfer, Eq.(13) has a current source that is connected with radial derivatives. These terms can contribute to a net axial current under certain conditions that are dependent on the radial profiles of the equilibrium plasma. An example of a configuration that does not receive a contribution to the net current from the second and third terms of Eq.(13) is one where the electron number density, n_{e0} , and the parallel resistivity, $\eta_{||}$, are constants, i.e. independent of the radius, and the poloidal component of the magnetic field, $B_{0\theta}$, varies linearly, implying that $B_{0\theta}/r$ is a constant. For this case, the integral of the second and third terms over the circular cross section of the plasma is zero if $\hat{j}_r$ vanishes at the plasma boundary. On the other hand, these two terms would likely contribute to the net current if the parallel resistivity were dependent on radius.

3 Kinetic model of Alfvén wave current drive

Here, the kinetic approach developed in Ref. [9, 32] was used to analyze the current drive produced by Alfvén waves in the cylindrical plasma column. Density, temperature, and current profile are proposed to be diffusive over the radius and homogeneous along the magnetic field lines as discussed in the previous section. This cylindrical plasma model is the reasonable for tokamaks and corona loops with low plasma pressure, ($\beta \ll 1$), and a small toroidicity parameter ($a/R \ll 1$, where a is the minor and R is the major radii of tokamaks or corona loops). We consider that the plasma is magnetized, so that the Larmor radii of electrons and ions are smaller than the radial inhomogeneity of the plasma parameters ($v_{Te,i}/\omega_{ce,i} \ll L_r = n/(dn/dr), T/(dT/dr)$). Weak collisions are taken into account ($\nu_{ei} \ll \omega$), where ν_{ei} is the electron-ion collision frequency, and ω is the frequency of the waves). The Landau form of the collision operator is considered in Vlasov's kinetic equation for electrons and ions.

To realize the program of the kinetic approach proposed in introduction we start from the last step related to averaging of the kinetic equations for electrons and ions over time or a spatial period of Alfvén waves over axial and poloidal directions. For this purpose we represent the distribution function of ions and electrons as a sum ($f_{0,e,i} + \tilde{f}_{e,i} \exp[i(kz + m\theta - \omega t)]$) of the quasistationary and oscillating parts as proposed for MHD variables in the section 2. Then, multiplying the Vlasov equation for averaged distribution functions electrons and ions, $F_{0,e,i}$ by the parallel momentum $m_{e,i}v_{||}$, integrating in velocity space, and combining these two equations, we obtain from the equation for the average density of the current:

$$j_{0z}^{RF} = -\frac{B_{0z}}{en_{e0}\eta_{||}B_0} \left\{ \frac{k_{||}}{\omega} P_e - \left[\frac{B_{0\theta}}{4r^2B_0} \frac{\partial}{\partial r} \left(r^2 \frac{\hat{j}_r^* \hat{E}_\theta}{i\omega} \right) + \frac{B_{0z}}{4rB_0} \frac{\partial}{\partial r} \left(r \frac{\hat{j}_r^* \hat{E}_z}{i\omega} \right) + c.c. \right] \right\}, \quad (14)$$

where $k_{||} = mB_{0\theta} + krB_{0z}/rB_0$, and the oscillations of plasma density and magnetic field are expressed from the equation of continuity and from the induction equation:

$$n_{e\omega} = -\frac{i}{e\omega} \nabla \cdot \vec{j}_{\alpha\omega}, \quad \dot{B}_\omega = -\frac{ic}{\omega} [\nabla \times \vec{E}_\omega]; \quad (15)$$

In Eq.(14) the current corrections related to oscillations of ions in the wave fields was neglected because of small electron-ion mass ratio m_e/m_i . Note, the same result as in Eq.(14) was obtained another method in Ref. [33] using the kinetic approach and this equation can be directly calculated from the general Ohm's law [34] for weak collision plasmas.

Comparing Eq.(13) obtained from two fluid model with Eq.(14) we observe the difference related to dynamic pressure force $F_{p\alpha}$ which is absent in the kinetic approach. This difference may be explained by two reasons:

- Two fluid MHD approach is not valid in weak collision limit because the parallel motion MHD equation is not closed (it need to take into account high moment equations, for example, for temperature and so on).
- In the kinetic approach the nonlinear corrections of the equilibrium distribution function, which is related to the equilibrium modified by waves, is not taken into account in Eq.(14).

This disagreement is not so important for the Alfvén waves with phase velocity less than thermal velocity because in this case the influence of the dynamic pressure force is very small in comparison with another forces ($F_{p,e\parallel} \approx (v_{ph}/v_T)^2 F_{d,e\parallel}$). In any case this question is under the study and the results will be published elsewhere.

Considering the smooth current profile in Eq.(14), we estimate the value of the magnetic field aligned (parallel) current drive:

$$j_{\parallel}^{RF} \approx -\frac{|e|k_{\parallel}}{m_e \nu_{ei} \omega} \left\{ P_e + \frac{1}{8\pi k_{\parallel} r} \frac{\partial}{\partial r} \left[\text{Im}(r \hat{j}_r^{(e)} \hat{E}_{\parallel}^*) \right] \right\}. \quad (16)$$

For homogeneous plasmas, the first term of the above equation is the current drive, calculated in Eq.(1), which is proportional to the density of dissipated power, P_e , as defined by an analogy to the homogeneous plasma case where $\bar{P}_e \approx (\omega/8\pi) \text{Im}(\epsilon_{33}) |E_{\parallel}|^2$. The second term is the gradient current related to the gradient of the density and the wave amplitude over the radius. This current contains the helicity injection current (see discussion, for example, in Ref. [32]):

$$j_h = -\frac{|e|}{8\pi \nu_{ei} m_e} \text{Im} \left[\frac{i\epsilon_{12}^{(e)}}{r} \frac{\partial}{\partial r} (r \hat{E}_b \hat{E}_{\parallel}^*) \right].$$

3.1 Wave propagation analysis

To clarify the current drive dependence in Eq.(14) from plasma parameters, we should fulfill the first and the second items of the program mentioned in introduction. On the first step, we had to defined the relation between the current oscillations and the electric field oscillations through a dielectric tensor:

$$j_{s,\omega}^{(e)} = -(i\omega/4\pi) \epsilon_{sq}^{(e)} E_{q\omega}$$

where the electron parts of the dielectric permeability tensor components for low frequencies $\omega \ll \omega_{ce}$ can be taken from [32]:

$$\begin{aligned}
\varepsilon_{11} &= \varepsilon_1 = \frac{\omega_{pi}^2}{\omega_{ci}^2 - \omega^2}, \quad \varepsilon_{12} = \varepsilon_{12}^{(i)} + \varepsilon_{12}^{(e)} = i\varepsilon_1 \frac{\omega_{ci}}{\omega} - \frac{\omega_{pe}^2}{\omega|\omega_{ce}|}; \\
\varepsilon_{22} E_b &= (\varepsilon_1 + \varepsilon_{22}^{(e)}) E_b = \varepsilon_1 \left\{ 1 - \frac{2}{n_e r} \frac{\partial}{\partial r} \left[\frac{\eta_e n_e v_{Te}^2}{\omega^2} \frac{\partial E_b}{\partial r} \right] \right\}; \\
\varepsilon_{23}^{(e)} E_{\parallel} &= -\frac{(1 - \eta_e) \omega_{ci}}{\omega \varepsilon_1 k_{\parallel} r} \frac{\partial}{\partial r} (r \varepsilon_1 E_{\parallel}) : \quad \varepsilon_{32}^{(e)} E_b = \varepsilon_1 \frac{(1 - \eta_e) \omega_{ci}}{\omega k_{\parallel} r} \frac{\partial}{\partial r} (r E_b); \\
\varepsilon_{33}^{(e)} &= \frac{\omega_{pe}^2}{k_{\parallel}^2 v_{Te}^2} (1 - \eta_e); \quad \eta_e = \frac{S_{ph}}{\sqrt{\pi}} \int \frac{\exp(-S^2)}{S_{ph} - S} dS, \quad S_{ph} = \frac{\omega}{\sqrt{2} |k_{\parallel}| v_{Te}}.
\end{aligned} \tag{17}$$

Here, the tensor indexes (1,2,3) mean radial (r), binormal (b), and parallel ($\parallel$) components respectively.

On the next step, we propose the geometric optic approximation for the waves traveling over the radial coordinate,

$$E_{r,b,\parallel} = E_{r,b,\parallel}^{(0)} \exp(i \int_0^r k_r dr), \tag{18}$$

to study the wave polarization. Substituting the above equation into the Maxwell's equations, it is easy to obtain the ratio between the parallel and binormal electric field components for low frequency waves ($\omega \ll \omega_{ci}$):

$$\frac{E_{\parallel}}{E_b} \approx \frac{(\varepsilon_{11} - N_{\parallel}^2)(\varepsilon_{22} - N_r^2 - N_{\parallel}^2) - N_b^2(\varepsilon_{22} - N_{\parallel}^2)}{N_r^2 N_b N_{\parallel} + \varepsilon_{23}(N_b^2 + N_{\parallel}^2 - \varepsilon_{11}) + \varepsilon_{13} N_r N_{\parallel}} \tag{19}$$

where $N_{r,b,\parallel} = k_{r,b,\parallel} c / \omega$ are the components of refractive index. This formula will be useful for current drive analysis.

For kinetic Alfvén wave (KAW) with the dispersion relation

$$N_r^2 \approx (\varepsilon_{11} - N_{\parallel}^2) \varepsilon_{33} / \varepsilon_{11}; \tag{20}$$

we find

$$E_{\parallel} \approx -E_b \left[(\varepsilon_{11} - N_{\parallel}^2) N^2 + i N_b^2 \text{Im}(\varepsilon_{22}) \right] / N_r^2 N_b N_{\parallel} \tag{21}$$

So, KAW may appear near the threshold of the Alfvén continuum, which can be defined by the equation:

$$\omega_{sh}(r) = k_{\parallel}(r) c_A(r), \tag{22}$$

when the frequency is much smaller than an ion cyclotron frequency, ω_{ci} . In the above equation, $c_A(r)$ is the spatially dependent Alfvén speed, defined from the equation:

$$\frac{1}{c_A^2(r)} = \frac{1}{c^2} \sum_i \left(\frac{\omega_{p,i}}{\omega_{c,i}} \right)^2. \tag{23}$$

In the course of tokamak plasma heating by Alfvén waves, partially ionized heavy impurities, such as oxygen or iron, introduce zones of ion cyclotron resonances than can

modify the dispersion characteristics of these waves. An important issue concerns the effects of lighter impurities, such as deuterium, helium or carbon, on the spectral and dispersive properties of Alfvén waves [35]. Note, the influence of the lighter impurities is very important for the Alfvén wave study in tokamak reactor because of working gas is proposed to be a deuterium and tritium (D-T) mixture. When the Alfvén wave frequency is much smaller than the ion cyclotron frequency of both the bulk plasma and the fully ionized impurities, only small changes in the properties of the Alfvén eigenmodes will occur, see Eq.(23). However, when the ion cyclotron frequencies of the light impurities are comparable to the Alfvén wave frequency, significant modifications of the Alfvén wave propagation and heating properties are likely. With a small population of impurities added to the plasma, gaps will appear in the Alfvén continuum about the cyclotron resonances of the impurity ions, $\omega \approx \omega_{cZ}$. This can be understood by considering the local dispersion relation that governs the Alfvén continuum,

$$\omega^2 \varepsilon_{11}(r) = k_{\parallel}^2(r) c^2. \quad (24)$$

where $\varepsilon_{11}(r)$ is the transverse component of the permeability tensor, which can be taken from Eq.(17). The Alfvén continuum is defined through the solutions of Eq.(24) for the frequency as a function of the radius, $\omega_{sh}(r)$. When the frequency is close to the cyclotron frequency of an impurity ion, Eq.(24) has no real solutions, implying the presence of a gap in the continuum. The solution of Eq.(24) for the Alfvén continuum frequency has two branches, as $\omega_{sh}^+(r)$ and $\omega_{sh}^-(r)$, which can be written for a plasma with two ion species in the following manner,

$$\left(\frac{\omega_{sh}^{\pm}}{\omega_{cZ}} \right)^2 = \frac{1 + \rho_z + \kappa^2(1 + A^2) \pm \sqrt{[1 + \rho_z + \kappa^2(1 + A^2)]^2 - 4\kappa^2[\rho_z + A^2(1 + \kappa^2)]}}{2(\rho_z/A^2 + 1 + \kappa^2)}. \quad (25)$$

In the above equation, ρ_z designates, $n_z M_z / n_H M_H$, the ratio of impurity mass density to the hydrogen density, A represents, $M_z e / M_H e_z$, the impurity charge-mass ratio normalized to the hydrogen mass-charge ratio, κ replaces $k_{\parallel} c / \omega_{pH}$, where ω_{pH} is the plasma frequency of the hydrogen ions. The ion-ion (upper) branch of the Alfvén continuum, $\omega_{sh}^+(r)$, has a minimum at the ion-ion hybrid frequency.

Fig. 1. The frequency of the surface Alfvén wave (SAW) and the Global Alfvén (GA) wave versus $k_{||}c/\omega_{pH}$ for two sets of wave parameters: $m = 1$, $1 \leq n \leq 12$ (labeled with boxes) and $m = -1$, $-12 \leq n \leq -1$ (labeled with diamonds). The impurity can be either carbon with $n_C/n_H = 0.078$, or helium with $n_{He}/n_H = 0.235$. For comparison, the frequency of the GA wave in a pure hydrogen plasma is also plotted (labeled with crosses).

In Fig. 1, we illustrate the gap in the Alfvén continuum for a hydrogen plasma with carbon as the impurity ion species with the density ratio $n_C/n_H = 0.078$ in a comparison with the shown case of the pure hydrogen plasma, $n_C/n_H = 0.0$. We have plotted specifically the threshold frequency of the Alfvén continuum, normalized to ω_{cD} , the cyclotron frequency of deuterium, versus κ , the normalized parallel wavenumber. Instead of the Alfvén continuum changing mentioned above, the spectrum of Global Alfvén (GAW) waves is changed that is also demonstrated on Fig.1. Note, sometimes GA waves are named as fast waves or ion cyclotron waves. As the mass density of the impurity ions increases, as a result, for example, of a larger impurity number density or impurity mass, the upper branch of the shear Alfvén continuum moves together with the surface Alfvén wave (SAE) in the direction of higher frequencies. The frequency of the GA eigenmodes also increases as impurity density increases.

Usually, GAW wave are situated below the Alfvén continuum (see, for example, [4]) but in the presence of the impurities the upper branch of GA waves is immersed into the Alfvén continuum and converts into SAE with diminishing parallel wavenumber, $k_{||}$. The best RF heating regimes had been found in Phaedrus-T tokamak experiments [37] around parameters when GAW waves was immersed the Alfvén continuum, $\omega \approx 1.4\omega_{cD}$. The Phaedrus-T heating experiments were performed in the high frequency range of the Alfvén spectrum, i.e. 7 MHz, with a short wavelength antenna, corresponding to $k_{||}$ approximately equal to 0.1 cm^{-1} , and a broad band of toroidal wavenumbers, from

approximately $n = 6$ to $n = 20$.

3.2 Estimation of Alfvén wave current drive

Now it is possible to clarify the influence of different terms in Eq.(16) on the value of the Alfvén wave current drive. Using the Eq.(21) we estimate the KAW current drive (which is related to the electron Landau damping) as

$$\langle j_{\parallel} \rangle \approx j_{cd} \left[1 - \frac{2\omega_{ci} \text{Re}(k_r) k_b^3 v_{Te}^2 \text{Im}(\eta_e)}{\text{Re}(\varepsilon_{33}) \omega^3 k_r} \right] \quad (26)$$

where j_{cd} is momentum transfer current drive mentioned in introduction and the second term is the helicity injection current. If the kinetic Alfvén wave is a standing wave along toroidal and poloidal directions, we should consider the current drive from waves with $k_{\parallel} = \pm |k_{\parallel}|$ and $k_b = \pm |k_b|$, and finally we obtain only the helicity injection (or gradient) current. The local efficiency (which is the ratio between a current and dissipated power) of KAW current drive is higher by one order than the simple current drive efficiency of travelling KAW ($j_h/\bar{P}_e \approx 10j_{cd}/\bar{P}_e$).

For Global Alfvén wave (GAW) and fast magneto-sonic waves (FMW), the value of the parallel oscillatory current will be zero, due to a high parallel conductivity. Based on this condition, we evaluate:

$$\varepsilon_{33} E_{\parallel} \approx -\hat{c}_{32} E_b \quad (27)$$

which is the particular case of Eq.(19). After substituting the above equation into Eq.(16), we obtain:

$$\langle j_{\parallel} \rangle \approx j_{cd} \left[1 + \frac{\text{Re}\eta_e}{\text{Im}(\eta_e)} \frac{\chi_N - 2\text{Im}(k_r)}{\text{Re}(k_r)} \right] \quad (28)$$

where χ_N is the radial inhomogeneous density parameter. In above equation (as for the KAW) the first term is the simple current drive and the second and third terms are the gradient (and helicity) current drive. If the GAW are standing waves ($k_{\parallel} = \pm |k_{\parallel}|$ and $\text{Re}k_r = \pm |k_r|$) the first and second terms of Eq.(28) will disappear and only the helicity current drive will exist.

The value of the GAW helicity (nonresonant) current drive will be

$$j_{\parallel}^{RF} \approx \frac{|e| k_{\parallel} \bar{P}_e (1 - \text{Re}\Lambda_e)}{m_e \nu_{ei} \omega} \frac{\chi_N}{k_r} \quad (29)$$

So, the local efficiency of the Alfvén wave current drive will be increased by one order due to nonresonant gradient forces when the KAW phase velocity is small $c_A \ll v_{Te}$ and GAW phase velocity is large $c_A > v_{Te}$. This additional nonresonant current drive will not depend on trapped particle effects, which reduce strongly the Alfvén current drive in tokamaks as was mentioned in the introduction.

Now let us begin to discuss the value of the Alfvén wave current drive in corona loop conditions. For estimations we use the standard plasma parameters of corona loops (see, for example [10, 11]) such as: axial magnetic field $B = 5 \cdot 10$ G, density $n_i m_i = 3 \cdot 10^{-15} \text{g/cm}^3$, temperature $T = 2.5 \cdot 10^6$ K, corona loop length $\pi R = 10^{10}$ cm, corona loop radius $a = 0.5 \cdot 10^9$ cm. Alfvén speed $v_A = 2.5 \cdot 10^8$ cm/s, and velocity oscillation amplitude $\delta v_r \approx \delta v_b = 3 \cdot 10^6$ cm/s.

Because in corona loop conditions the Alfvén waves should be standing waves as eigenvalues of the system with fixed ends, we had to consider the current drive from waves with $k_{\parallel} = \pm|k_{\parallel}|$ and $k_b = \pm|k_b|$. Finally, we obtain only helicity injection (or gradient) current because of standing wave momentum is equal zero. If all radiated power is brought into the corona loop plasma due to electron heating by the kinetic Alfvén waves with the density of dissipated power $P = 10^{-3} \text{ erg cm}^{-3}\text{s}^{-1}$, the local value of helicity current drive will be $\approx (10^8) \text{ statA/cm}^2$. This value exceeds any possible corona loop current density in (10^5) times. In reality, the average value of the KAW current drive should be smaller, because KAW weren't excited in corona loops directly. These waves appear due to a conversion of surface and combined Alfvén modes (see [12]). The KAW propagates from the conversion point to the plasma center in an interval that equals KAW dissipation length ($L_r \approx 1/\text{Im}k_r \approx 10^4 \text{ cm}$). For supporting of a diffuse current profile it needs to have a spectrum of surface or combined Alfvén modes to overlap the current profile over a corona loop minor radius and to compensate the back current in the regions outside the conversion point. Thus, the average current drive efficiency will be smaller by a factor ($\approx 1/(\text{Im}k_r) \approx 10^{-5}$) and, finally, in this case the value of the current drive density will be $\langle j_h \rangle \approx 10^3 \text{ statA/cm}^2$. This value of the current will be enough to explain the corona loop poloidal magnetic field order of ($\sim 10 \text{ G}$), which needs for plasma equilibrium and stability.

For GAW or fast magneto-sonic waves in the corona loop plasma conditions discussed above ($\text{Re} \eta_e / \text{Im} \eta_e \approx \omega / k_{\parallel} v_{Te} \approx 1$), we estimate the current drive to be:

$$\langle j_{gr} \rangle = 10^6 \frac{P_e}{\omega \langle W \rangle} \text{statA/cm}^2 \quad (30)$$

where $\langle W \rangle = n_i m_i \langle v_r^2 \rangle / 2$ is the observed density of Solar corona loops oscillation energy. If the GAW quality resonance parameter $Q \approx 1000$ (where $Q \approx \omega W / P_e$), the value of the gradient current drive in Eq.(10) will be enough to sustain the corona loops poloidal magnetic field ($\approx 10\text{G}$).

4 Experimental evidence of the Alfvén wave current drive

4.1 R-0 stellarator experiments

The first time the Alfvén wave current drive had been demonstrated in RF heating experiments with travelling Alfvén waves in the R-0 stellarator (Sukhumi/former USSR)[5]. The parameters of the R-0 stellarator were: major and minor radii were 50 cm and 3.5 cm, antenna radius 5cm, magnetic field B_t was $\leq 8 \text{ kG}$, rotational transform angle i_0 was $\leq .8$, electron temperature was $\leq 80 \text{ eV}$, and the averaged plasma density $\bar{n}_e$ was varied in the interval $2\text{-}8 \cdot 10^{13} \text{ cm}^{-3}$. The direct measurements of the current drive was done with Rogovsky coil. This current have changed the direction with the changing direction of the travelling Alfvén waves, which was launched by multi-pole helical antenna with poloidal and toroidal wavenumbers $m/n = \mp 2 / \pm 2$. Showing a good agreement with Eq.(1) the magnetic probe measurements indicated the hollow profile of driving current because of the hollow profile of the dissipated power (see for details Ref. [4]). The maximum of

dissipated power was situated near the conversion point position, which can be defined from the equation, $\varepsilon_{11}(r) = k_{\parallel}^2 c^2 / \omega$. The conversion of RF fields can be related to the kinetic Alfvén waves, which have the dispersion relation derived by Eq.(20).

In the next series of the RF heating experiments in the R-0 stellarator (see Ref. [36]) the additional current to the momentum transfer current drive had been shown. This current can change the direction of the current drive during the same RF pulse.

Fig.2. The temporal evolution of the plasma current, absorbed RF power, energy content and averaged plasma density for the case travelling Alfvén waves with the frequency 1.2 MHz and the poloidal and toroidal wave numbers $m/n = -2/2$, magnetic field $B_t = 3.1 \text{ kG}$, and rotational transform .31.

The results of this experiments are presented in Fig. 2. The variations of the plasma current, dissipated RF power from the travelling wave (with the frequency 1.2 MHz and wavenumber combinations $m/n = -2/2$), energy content and averaged plasma density are shown. On the first stage of the RF pulse, when the plasma density was stabilized on the level of $2 \cdot 10^{13} \text{ cm}^{-3}$, the momentum transfer current drive (approximately 70 A), appeared. With growing of the density up to $3 \cdot 10^{13} \text{ cm}^{-3}$ the opposite direction current appeared, which overcame the momentum transfer current drive. This effect may be related to decreasing of the momentum transfer current drive efficiency due to density

(and electron-ion collision frequency) growing. Another reason can be appeared due to the maximal efficiency of gradient (helicity) current drive when the Alfvén phase velocity is the order of the thermal velocity. In this case the calculated helicity current drive had the opposite direction in a comparison with momentum transfer current drive. In the same conditions the value of the current drive efficiency induced by standing waves ($m/n = \mp 2/\pm 2$) was in 1.5 times more than caused by travelling waves and the current flow was in the direction of the helicity injection current drive.

4.2 Current drive in the Phaedruss-T experiments

During the last years the main goal of the Phaedruss-T experimental program was to study the Alfvén wave heating and current drive and to understand factors, $k_{\parallel}$ spectrum, toroidal geometry influence and trapped electron effect as well as determining dissipated power and current drive profiles and efficiency. On the first step of the Phaedruss-T experiments[37] the efficient Alfvén wave heating regime had been found. Here, RF fields were launched by π -phasing two strap antenna located on the high field side. The electron heating experiments was performed with parameters: $R=93$ cm, $a=26$ cm, $B_0=6.5-9$ kG, $I_p=60-85$ kA, $T_{e0}=400-600$ eV, $\bar{n}_e=.2-1.5 \cdot 10^{13}\text{cm}^{-3}$, working gas hydrogen. The frequency of RF field was 7 MHz and the maximum of the corresponding k_{tor} spectrum was approximately equal to 0.1 cm^{-1} . The efficient electron heating $\Delta T_{e0} \approx 100$ eV and loop voltage drop $\approx 15\%$ were achieved with dissipated power 300 kW when the density was $\bar{n}_e=1 \cdot 10^{13}\text{ cm}^{-3}$. These RF heating conditions in Phaedruss-T can be explained by the mode conversion theory of GAW, which immersed into ion-ion the Alfvén continuum (see discussion in subsection 3.1). Recently, very similar experiment on RF heating of D-⁴He-³He plasma mixture in TFTR tokamak had been reported in Ref.[38].

On the next step of the Phaedruss-T experiments [6] (see Ref. [7] for details) the first Alfvén wave current drive in tokamaks had been successfully demonstrated with $\pi/2$ -phasing of two strap antenna. The Phaedruss-T parameters were the same as for heating regimes but the plasma density was low, $\bar{n}_e=.2 \cdot 10^{13}\text{ cm}^{-3}$. The on-axis magnetic field was .68 T, which corresponds to $\omega/\omega_{cH} = .7$. Under these conditions the deuterium cyclotron resonance was on the high field side at the edge of the plasma and the hydrogen cyclotron resonance was outside of the plasma on the low field side of the toroidal magnetic field.

Fig.3. The temporal behavior of loop voltage and the radial profiles of electron temperature in the Phaedrus-T tokamak with RF heating.

The results of the experiments taken from [6] are shown in Fig.3. A loop voltage drop $\Delta V/V = 36\%$ can be observed. Electron temperature measurements indicate that the loop voltage drop can not be related to an increase of T_e . The constancy of the electron temperature in time may be caused by increasing of electron transport and obvious rising of impurities, $\Delta Z_{eff}/Z_{eff} \approx 50\%$, which was detected by the visible bremsstrahlung emission. Increasing loop voltage (after switch of RF power) in comparison with the loop voltage before the RF pulse ($\Delta V/V \approx 15\%$ because of the increasing resistance) confirms the rising of Z_{eff} . During RF pulse the plasma position and current was kept constant, $I_{pl}=64$ kA. Although the internal inductance of the plasma column, l_{int} , was not measured, the contribution of dl_{int}/dt can be neglected. Finally, the value of the current drive, $I_{cd} = 24 - 36$ kA, can be estimated. This current drive value was in a good agreement with the theoretical predictions of the current drive efficiency in Ref. [19] and Eq.(1) when the phase velocity was of the order of the thermal velocity and dissipated power in the plasma core, W_{RF} , was ≈ 130 kW, which can be found from cylindrical kinetic code calculations [35].

5 Conclusion

In conclusion, we can say that:

- A general form of the time-averaged ponderomotive force produced by RF waves in a magnetized plasma has been evaluated. The ponderomotive force in two fluid MHD approximations includes contributions from oscillating plasma flow, dynamo effect, and electrostrictive force.
- RF driven current in a cylindrical plasma configuration has been found in two fluid MHD and weak collision kinetic approximations. Assuming parameters that are characteristic of the tokamaks and solar corona loops, our calculations indicate that spatial gradient effects (or helicity) may dominate momentum current drive. In addition, axial current may be driven by poloidal ponderomotive forces, which are composed of effects associated with poloidal momentum dissipation and radial gradients of the wave fields

and the plasma profiles.

- It has been shown that standing waves may produce helicity injection current. In the specific case of corona loops, if the value of the corona loop's oscillation energy is related only to standing GAW, the surface, or the combined Alfvén eigenmodes, then the value of the GAW gradient (or SAE and CAE helicity) current drive will be enough to sustain the proposed corona loop's poloidal magnetic field ($\approx 10G$).

- The maximal local efficiency of the gradient current drive is caused by the kinetic Alfvén wave but the total efficiency of this current drive, which is averaged over the radius, is less than that caused by the Global Alfvén wave. The reason is that the KA wave is excited by conversion of SAE or CAE and produce back current at the interval out the conversion point. This back current should be compensated with the broad spectrum of SAE or CAE.

- Experimental evidence of the Alfvén wave current drive was obviously demonstrated in the R-0 stellarator and Phaedrus-T tokamak during RF heating experiments with standing and travelling waves.

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