26th EPS Conf. on Contr. Fusion and Plasma Physics, Maastricht, 14 - 18 June 1999 ECA Vol.23J (1999) 1813 - 1816

Near-Antenna Density Channels in a Magnetoplasma: Ionization Formation and Ducting Properties in the Lower-Hybrid Band

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Abstract.–An investigation is made of the structure and whistler ducting properties of steady-state magnetic-field-aligned plasma channels created due to a rf discharge in the vicinity of a magnetic antenna in the lower-hybrid band. Self-consistent distributions of the electron temperature and plasma density, obtained by numerically solving the heat conduction equation and the continuity equation, are discussed for the ionospheric conditions.

Introduction

Over the past decade there has been a substantial degree of interest in near-antenna rf discharge structures produced by ionizing a neutral gas in an external dc magnetic field. However, at present, there exists very little theory dealing with the ionization formation of density nonuniformities by the antenna near-zone fields when the exciting antenna is immersed in an already created (background) magnetoplasma. It is the purpose of the present paper to discuss the structure and ducting properties of a field-aligned density enhancement caused by a rf discharge in the vicinity of a magnetic antenna (current loop) immersed in a homogeneous magnetoplasma and operating in the lower-hybrid (whistler) band

$$\omega_{\rm LH} \ll \omega < \omega_{\rm H} \ll \omega_{\rm p}, \tag{1}$$

where $\omega_{\rm LH}$ is the lower-hybrid frequency, and $\omega_{\rm H}$ and $\omega_{\rm p}$ are respectively the gyrofrequency and the plasma frequency of the electrons.

The discharge is assumed to be excited by the field of a loop the electric current in which is given, with $\exp(i\omega t)$ time dependence dropped, by

$$\mathbf{J} = \hat{\phi}_0 2 I_0 \frac{\rho}{a^2} \exp(-\rho^2/a^2)\delta(z), \tag{2}$$

where I_0 is the total current magnitude, a is the radial extent of the source distribution, δ is the Dirac delta function, and ρ, ϕ, z are the cylindrical coordinates. We fuzzify the current in the radial direction in order to eliminate the field singularities arising in the case of a infinitesimally thin current ring. Note that the problem of discharge, in which the density profiles and thermal distributions are found self-consistently, is treated herein for the representative case of a background plasma modeled upon the ionosphere.

Field structure and temperature distribution

Let the plasma diffusion lengths along and across the external magnetic field $\mathbf{B}_0 = B_0\hat{\mathbf{z}}_0$ be much greater than the source dimension a, and let the effective electron-collision frequency $\nu_{\rm e}$ be such that $\nu_{\rm e} \ll \omega$. Then the field in the antenna near zone can be calculated with the use of the quasistatic theory neglecting plasma nonuniformity. Within this theory the electric field is expressible via a scalar potential ψ and a vector potential $\mathbf{A}$, $\mathbf{E} = -\nabla\psi - ik_0\mathbf{A}$, and the equations for ψ and $\mathbf{A}$ are

$$\nabla^2 \mathbf{A} = -\frac{4\pi}{c}\mathbf{J},$$
$$\varepsilon \nabla_\perp^2 \psi + \eta \frac{\partial^2 \psi}{\partial z^2} = -\mathrm{i} k_0 \nabla \cdot \hat{\varepsilon} \cdot \mathbf{A}, \qquad (3)$$

where $\nabla_\perp$ is the transverse (to $\mathbf{B}_0$) part of the *del* operator ∇, $k_0 = \omega/c$ is the wavenumber in free space, and $\hat{\varepsilon}$ is the plasma dielectric tensor, whose non-zero elements are $\varepsilon_{\rho\rho} = \varepsilon_{\phi\phi} = \varepsilon$, $\varepsilon_{\rho\phi} = -\varepsilon_{\phi\rho} = -\mathrm{i}g$ and $\varepsilon_{zz} = \eta$. In the frequency band (1) these elements are written as

$$\varepsilon = \omega_{\mathrm{p}}^2/(\omega_{\mathrm{H}}^2 - \omega^2),\ g = -\omega_{\mathrm{p}}^2\omega_{\mathrm{H}}/\omega(\omega_{\mathrm{H}}^2 - \omega^2),\ \eta = -\omega_{\mathrm{p}}^2/\omega^2. \qquad (4)$$

The solution for the electric field due to the source (2) can be found to be

$$E_\rho = \mathrm{i}E_0 g(\varepsilon - \eta)^{-1}\left[\mathcal{F}_1^{(1)} + \mathrm{i}\,(-\varepsilon/\eta)^{-1/2}\,\mathcal{F}_1^{(2)}\right],$$
$$E_\phi = E_0\,\mathcal{F}_1^{(1)},\ E_z = -\mathrm{sgn}(z)g(\varepsilon - \eta)^{-1}\left[\mathcal{F}_0^{(1)} - \mathcal{F}_0^{(2)}\right], \qquad (5)$$

where

$$\mathcal{F}_m^{(1,2)} = \int_0^{2\pi} \xi_{1,2}\,\mathrm{erfc}(\xi_{1,2})\,\exp(\xi_{1,2}^2 + \mathrm{i}m\zeta)\mathrm{d}\zeta,$$
$$\xi_1 = (\mid z \mid -\mathrm{i}\rho\,\cos\zeta)/a,\ \xi_2 = (\mathrm{i} \mid z \mid (-\varepsilon/\eta)^{1/2} - \mathrm{i}\rho\,\cos\zeta)/a,\ E_0 = I_0\pi^{1/2}k_0/c.$$

It is important that the field (5), with the tensor elements given by (4), turns out to be independent of the plasma density.

With the above field, we solve the equation governing the heating of electrons. Assuming that the axial and transverse diffusion lengths considerably exceed the corresponding heat-conduction lengths and that the transverse length of heat conduction is much smaller than a, which conditions are typical of the lower ionosphere, the equation for the steady-state electron temperature T_e reads

$$\frac{\partial}{\partial z}\left(\frac{T_{\mathrm{e}}}{m\nu_{\mathrm{e}}}\frac{\partial T_{\mathrm{e}}}{\partial z}\right) + Q - \delta_{\mathrm{e}}\nu_{\mathrm{e}}\,(T_{\mathrm{e}} - T_{\mathrm{e}0}) = 0, \qquad (6)$$

where the term

$$Q = \frac{\nu_{\mathrm{e}}e^2}{3m\omega_{\mathrm{H}}^2}\left[\mid E_\rho \mid^2 + \mid E_\phi \mid^2 + 4\frac{\omega_{\mathrm{H}}}{\omega}\mathrm{Im}\left(E_\rho E_\phi^*\right) + \frac{\omega_{\mathrm{H}}^2}{\omega^2} \mid E_z \mid^2\right] \qquad (7)$$

represents the Joule heating rate, e and m are the electron charge and mass, δ_{e} is the average relative fraction of energy lost by the electron in a collision with a neutral molecule, and $T_{\mathrm{e}0}$ is the background (unperturbed) electron temperature. In writing (7) use was made of the simplifying condition $\omega^2 \ll \omega_{\mathrm{H}}^2$ which is henceforth assumed. The dependence of δ_{e} and ν_{e} on the temperature T_{e} was taken from [1,2]. Equation (6) was solved numerically under the conditions $\partial_z T_{\mathrm{e}} = 0$ when $z = 0$, and $T_{\mathrm{e}} \to 0$ when $z \to \infty$.

Figure 1 shows the spatial distribution of Q for $\omega = 3 \times 10^6$ s^{-1}, $a = 2$ m and the following background plasma parameters typical of the ionosphere at altitude 150 km: $\omega_{\mathrm{p}0} = 3.1 \times 10^7\mathrm{s}^{-1}$, $\omega_{\mathrm{H}} = 8.8 \times 10^6\mathrm{s}^{-1}$ and $\nu_{\mathrm{e}0} = 1.4 \times 10^3\mathrm{s}^{-1}$. Figure 2 displays T_{e} at $\rho = 0, z = 0$ versus I_0 for the same parameters. The curve $T_{\mathrm{e}}(0, z)$ for $I_0 = 78$ A is presented in Fig.3.

Density distribution

The distribution of the plasma density N is governed by the continuity equation. Let $L_{\|}$ and $L_{\perp}$ be the characteristic spatial scales of the density variation along and across $\mathbf{B}_0$, respectively. Under the condition $\left(L_{\|}/L_{\perp}\right)^2 \gg D_{e\|}/D_{i\perp}$, the validity of which is confirmed by the forthcoming analysis, the steady-state plasma continuity equation reads

$$2D_{e\perp}\nabla_{\perp}^2 N + 2D_{e\|}\frac{\partial^2 N}{\partial z^2} + \nu N - \alpha N^2 + q_{ext} = 0, \qquad (8)$$

where $\nu = \nu_{ion} - \nu_T - \nu_a$; $\nu_T = -D_{e\perp}^{(T)}\nabla_{\perp}^2 T_e/T_e$; $D_{e\perp}$, $D_{e\|}$ and $D_{i\perp}$ are the coefficients of diffusion across ($\perp$) and along ($\|$) $\mathbf{B}_0$ for the electrons and ions in the background plasma; $D_{e\perp}^{(T)}$ is the transverse thermal-diffusion coefficient for the electrons; ν_{ion} represents the frequency of electron impact ionization of neutral molecules; ν_a is the electron attachment frequency; α is the electron-ion recombination coefficient; and q_{ext} is the intensity of the external ionization sustaining the background density N_0. Since the diffusion lengths are assumed to be fairly large, the quantities ν and α, which depend [1] on the known function $T_e(\rho, z)$, behave like the delta function in the (ρ, z) space, and hence can be replaced by functions $\nu(\tilde{r})$ and $\alpha(\tilde{r})$ with similar properties, where $\tilde{r} = (\rho^2 + \tilde{z}^2)^{1/2}$ and $\tilde{z} = z\left(D_{e\perp}/D_{e\|}\right)^{1/2}$, provided that the integrals of ν and α over the entire space are conserved. Then, introducing $D = 2D_{e\perp}$, we reduce eq.(8) to a more simple form:

$$D\left(\frac{\mathrm{d}^2 N}{\mathrm{d}\tilde{r}^2} + \frac{2}{\tilde{r}}\frac{\mathrm{d}N}{\mathrm{d}\tilde{r}}\right) + \nu(\tilde{r})N - \alpha(\tilde{r})N^2 + q_{ext} = 0. \qquad (9)$$

Equation (9) is readily solved in the $\tilde{r}$ space, together with the conditions $\partial_{\tilde{r}}N = 0$ when $\tilde{r} \to 0$, and $N \to 0$ when $\tilde{r} \to \infty$. The typical solution of eq.(9) is that corresponding to the formation of a field-aligned density channel with enhanced plasma density. The value of N/N_0 at $\rho = z = 0$ varies from $N/N_0 = 1.01$ to $N/N_0 = 10^2$ when I_0 varies from $I_0 = 75$ A to $I_0 = 80$ A. A representative example of the radial distribution of N at $z = 0$, obtained by numerically solving eq.(9) for $I_0 = 78$ A, is shown in Fig.4.

Ducting properties of the density channel

The created channel can be recognized as a duct capable of guiding the antenna-launched whistler-mode waves. Since in the case considered the creation of the channel is a result of the nonlinear interaction of the strong quasistatic antenna field with the plasma, and is not associated with the wave fields, which have rather small amplitudes, the excitation and propagation of whistler-mode waves away from the source can be treated within a linear approach. Using a full wave formulation [3], we calculate the total field in the channel. For computed density distribution (e.g., Fig.4), it is shown that there is strong coupling of helicon waves to quasielectrostatic (Trivelpiece-Gould) waves in the ducted whistler modes which slightly leak from the channel into quasielectrostatic waves of the background plasma. We found that the presence of such a duct can lead to a substantial increase in the radiation power of the source.

Conclusion

The analysis performed shows that the additional ionization of a plasma in the lower

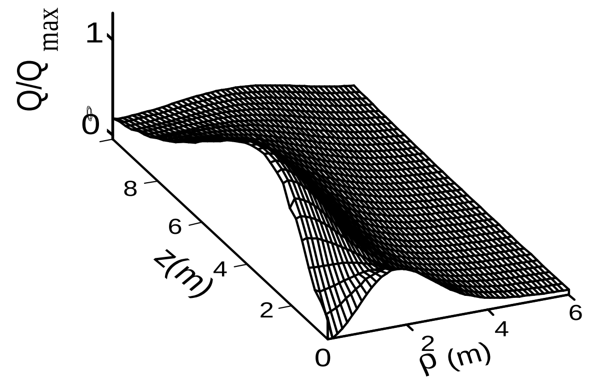

Figure 1: *The distribution of the Joule heating rate (see text for discussion)*

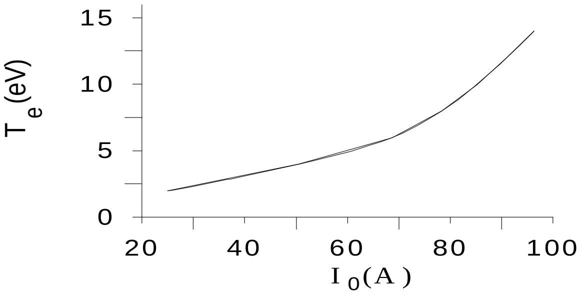

Figure 2: *Electron temperature* T_e *vs.* I_0 *for* $a = 2m$, $\omega = 3 \times 10^6 s^{-1}$, $\nu_{e0} = 1.4 \times 10^3 s^{-1}$, $B_0 = 0.5G$, *and background density* $N_0 = 3 \times 10^5 cm^{-3}$

ionosphere in the vicinity of a magnetic antenna can result in the formation of a field-aligned channel with enhanced plasma density. The channel can guide whistler mode waves. The guided waves turn out to be slightly leaky modes whose energy leaks from the channel into quasielectrostatic (lower hybrid) waves of the background plasma. The presence of the channel causes an increase in the power radiated from the source in the whistler band.

Acknowledgments.–This work was supported in part by the Russian Center for Fundamental Natural Sciences (KTsFE) under grant no. 97-0-8.2-15 and in part by the Russian programs "Universities of Russia" and "Leading Scientific Schools"(grant no. 96-15-96591).

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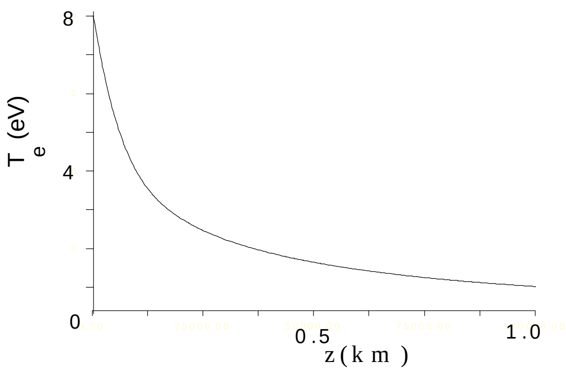

Figure 3: *Axial distribution of the electron temperature computed for* $I_0 = 78A$ *Source dimensions and plasma parameters are the same as for Fig.2*

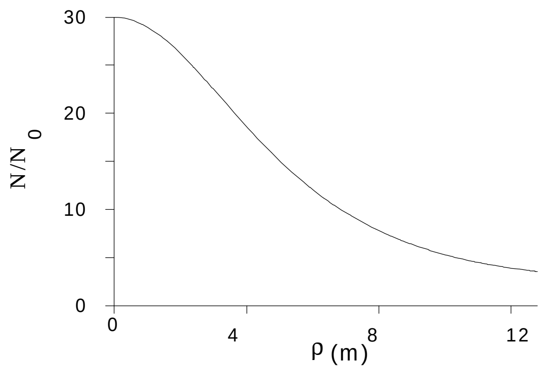

Figure 4: *Radial distribution of plasma density normalized to the background value* N_0, *for* $I_0 = 78A$. *Other conditions are the same as for Fig.3*