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QUASISTATIC IMPEDANCE OF A THIN CONDUCTOR IN A MOVING ISOTROPIC COLLISIONAL PLASMA

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The equation for a thin cylindrical antenna, located in a plasma moving along the axis of the conductor, is derived in quasistatic approximation under the condition of excitation of forced oscillations in the system taking electron collisions into consideration; the impedance of the antenna is computed. The plasma is assumed homogeneous, isotropic, and nondispersive in the reference system bound to the plasma.

The excitation of characteristic quasistatic oscillations of the potential in a system consisting of a moving magnetoactive plasma and a thin cylindrical antenna (TCA) has been investigated in [1]. The problem of forced oscillations of the potential and the associated parameters of the radiator in such a system is also of considerable interest. The quasistatic input impedance of a thin ideal conductor located in a moving isotropic collisionless plasma was computed earlier in [2] by using the Green's function of Poisson's equation for an oscillating point charge.

It is known [2, 3] that in the case of forced oscillations the electromagnetic field of a elementary radiator in a moving plasma in quasistatic approximation is characterized by the presence of "trails" of Langmuir-type oscillations besides the "vacuum" term; these are carried away from the source by the flow of the surrounding medium. A "plasma-wave" field is generated as a result of interference of two waves with wave numbers

$$k_j = (\omega \pm \omega_0)/u, \tag{1}$$

where ω is the radiation frequency, ω_0 is the electron plasma frequency, and u is the flow velocity of the plasma ($v_{Te} \ll u \ll c$, v_{Te} is the mean thermal velocity of electrons, c is the speed of light); here and below $j = 1, 2$. Correspondingly, in the problem of determination of the impedance of a macroscopic source large "plasma-wave" parameters appear along with the usual large "statistical" parameters [4]

$$B = \ln(2L/a), \quad a \ll 2L \tag{2}$$

(2L is the length, a is the radius of the conductor); these are (for the axis of the cylindrical antenna oriented along the velocity vector of the plasma) [2]

$$A_j = K_0(|k_j|a) \simeq \ln|k_j|a, \quad |k_j|a \ll 1 \tag{3}$$

[$K_0(x)$ is Macdonald function]. At the plasma resonance ($\omega = \omega_0$) the amplitude of the wave with subscript j = 2 corresponding to the difference frequency in (1) increases indefinitely [2, 3, 5]. As a result the parameter of a radiator located in a moving plasma may have singularities at $\omega \to \omega_0$. In this case a more detailed analysis and, in particular, the consideration of the dissipative processes in the medium become necessary.

In the present article the equation of a TCA in a moving isotropic collisional plasma is derived by using convolution of Poisson's equation over plane waves under the same assumptions as in [1, 2] and its input

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impedance is computed; the derivation of the equation by itself is of independent interest.

By decomposing the solution of Poisson's equation into three-dimensional Fourier integral the potential of the electric field of a cylindrical source in a plasma can be expressed in the following way:

$$\varphi(r, z)=\frac{1}{\pi}\int_{-L}^{L} d\zeta\,\sigma(\zeta)\int_{0}^{\infty} dk_{\perp}k_{\perp}J_0(k_{\perp}a)\int_{-\infty}^{\infty}\frac{\exp[ik_z(z-\zeta)]}{k^2\varepsilon_c(\omega, k)}\,dk_z. \tag{4}$$

Here r, z are the cylindrical coordinates of the point of observation (0z axis is directed along the axis of the antenna which, in turn, is directed along vector **u**), $k_\perp$, k_z are cylindrical coordinates of wave vectors **k** in space, and $J_0(x)$ is Bessel function of first kind and zero index; $\sigma(z)$ is the linear charge density,

$$\varepsilon_c(\omega, k)=1-\omega_0^2/(\omega-ku)(\omega-ku+i\nu) \tag{5}$$

is the dielectric constant of the moving collisional plasma in hydrodynamic approximation [2], and ν is the effective electron collision frequency; the time dependence is taken in the form $\exp(-i\omega t)$.

Substituting (5) into (4) and carrying out integration over k_z and then over $k_\perp$ in (4) we get

$$\varphi(r, z)=\int_{-L}^{L}\sigma(\zeta)G_c(r, z-\zeta)\,d\zeta; \tag{6}$$

$$G_c(r, z)=\frac{1}{\pi\sqrt{ar}}Q_{-1/2}\left(\frac{r^2+z^2+a^2}{2ar}\right)-k_0^2[P(r, z)1(z)+$$

$$+\tilde{P}(r, z)1(-z)]+\frac{2ik_0^2}{k_{0c}}[I_0(k_{1c}a)K_0(k_{1c}r)e^{ik_{1c}z}-I_0(k_{2c}a)K_0(k_{2c}r)e^{ik_{2c}z}]\,1(z); \tag{7}$$

$$\begin{Bmatrix}P(r, z)\\ \tilde{P}(r, z)\end{Bmatrix}=\pm\frac{i}{k_{0c}}\sum_{j=1}^{2}(-1)^{j+1}\begin{Bmatrix}S_{jc}(r, z)\\ \tilde{S}_{jc}(r, z)\end{Bmatrix}; \tag{8}$$

$$\begin{Bmatrix}S_{jc}(r, z)\\ \tilde{S}_{jc}(r, z)\end{Bmatrix}=\int_0^{\infty}\frac{e^{\mp k_\perp z}J_0(k_\perp a)J_0(k_\perp r)}{k_\perp\pm ik_{jc}}\,dk_\perp; \tag{9}$$

$$k_{jc}=(2k\pm k_{0c}+ik_c)/2,\quad k_{0c}=(4k_0^2-k_c^2)^{1/2},\qquad 2k_0>k_c,$$
$$k=\omega/u,\quad k_0=\omega_0/u,\quad k_c=\nu/u. \tag{10}$$

In (7) $Q_{-1/2}(x)$ is spherical function of the second kind and zero order of degree $-1/2$; $1(z)$ is unity function; $I_0(z)$ is Bessel function of second kind and zero order; the following formulas have been used in the derivation of (7):

$$\int_0^{\infty}e^{-k_\perp|z|}J_0(k_\perp a)J_0(k_\perp r)\,dk_\perp=\frac{1}{\pi\sqrt{ar}}Q_{-1/2}\left(\frac{r^2+z^2+a^2}{2ar}\right),$$

$$\int_0^{\infty}k_\perp J_0(k_\perp a)J_0(k_\perp r)(k_\perp^2+k_{jc}^2)^{-1}dk_\perp=I_0(k_{jc}a)K_0(k_{jc}r), \tag{11}$$

$$r>a,\quad \mathrm{Re}\,k_{jc}>0.$$

Here and below in expressions of type (8) and (9) the upper (correspondingly lower) variants of the formula in the left-hand side corresponds to the upper (lower) variant of the equation and (or) to upper (lower) sign in the right-hand side. Furthermore, all the quantities pertaining to the case of collisional plasma and having their counterparts for medium without loss (which will be denoted by the same letters) are furnished with an additional subscript "c."

We note that using the identity $1/p=\int_0^{\infty}\exp(-p\xi)\,d\xi$ $(\mathrm{Re}\,p>0)$ and the first of formulas (11), function (9) can be written in the following form:

$$\begin{Bmatrix}S_{jc}(r, z)\\ \tilde{S}_{jc}(r, z)\end{Bmatrix}=\frac{1}{\pi\sqrt{ar}}\int_0^{\infty}e^{\mp ik_{jc}\xi}Q_{-1/2}\left[\frac{(\xi\pm z)^2+r^2+a^2}{2ar}\right]d\xi. \tag{12}$$

We now write the expression for the field potential at the antenna itself in the following way assuming that the antenna is thin:

$$\varphi(z) \equiv \varphi(r = a + 0,\ z) = \Phi_e[\sigma, z] - P[\sigma, z] + \Phi_i[\sigma, z]; \tag{13}$$

$$\Phi_e[\sigma, z] = \frac{1}{\pi a}\int_{-L}^{L} \sigma(\zeta)\, Q_{-1/2}\left[\frac{(z-\zeta)^2 + 2a^2}{2a^2}\right] d\zeta;$$

$$P[\sigma, z] = \int_{-L}^{L} \sigma(\zeta) P_e(z - \zeta)\, d\zeta, \tag{14}$$

$$P_e(z) = k_0^2 \begin{cases} P(z), & z > 0 \\ \tilde{P}(z), & z < 0 \end{cases}, \qquad \begin{Bmatrix} P(z) \\ \tilde{P}(z) \end{Bmatrix} \equiv \begin{Bmatrix} P(a, z) \\ \tilde{P}(a, z) \end{Bmatrix}; \tag{15}$$

$$\Phi_i[\sigma, z] = \int_{-L}^{z} \sigma(\zeta) K_c(z - \zeta)\, d\zeta; \tag{16}$$

$$K_c(z) = 2ik_0^2 k_{0c}^{-1}\left(A_{1c} e^{ik_{1c}z} - A_{2c} e^{ik_{2c}z}\right); \tag{17}$$

$$A_{jc} = K_0(k_{jc} a). \tag{18}$$

Using the integral representation of the spherical function [6]

$$Q_{-1/2}(x) = \frac{1}{\sqrt{2}}\int_0^{\pi} \frac{d\vartheta}{\sqrt{x - \cos\vartheta}} \tag{19}$$

integral $\Phi_e[\sigma, z]$ and the functions

$$\begin{Bmatrix} S_{jc}(z) \\ \tilde{S}_{jc}(z) \end{Bmatrix} \equiv \begin{Bmatrix} S_{jc}(a, z) \\ \tilde{S}_{jc}(a, z) \end{Bmatrix} \tag{20}$$

reduce to the form

$$\Phi_e[\sigma, z] = \int_{-L}^{L} \frac{\sigma(\zeta)\, d\zeta}{R(z - \zeta)}; \tag{21}$$

$$\begin{Bmatrix} S_{jc}(z) \\ \tilde{S}_{jc}(z) \end{Bmatrix} = \int_0^{\infty} \frac{e^{\mp ik_{jc}\xi}\, d\xi}{R(\xi \pm z)}; \tag{22}$$

$$\frac{1}{R(z)} = \frac{1}{\pi}\int_0^{\pi} \frac{d\vartheta}{[z^2 + 4a^2 \sin^2(\vartheta/2)]^{1/2}}. \tag{23}$$

The functional $\Phi_e[\sigma, z]$ in (13) has the significance of potential distribution along a cylindrical conductor of radius a and length 2L with charge distribution described by $\sigma(z)$ in vacuum [4]. Integral $\Phi_i[\sigma, z]$ in the right-hand side of (13) is related to the motion of the plasma and is due to the excitation of the oscillations trail (see above). Finally, the functional $P[\sigma, z]$ in (13) appears due to the presence of near fields of longitudinal waves in the moving medium.

For a transition to the case of stationary plasma in (13), it is necessary to consider that functions $S_{jc}(z)$ and $\hat{S}_{jc}(z)$ can be written in the following way after integrating once by parts in (22):

$$\begin{Bmatrix} S_{jc}(z) \\ \tilde{S}_{jc}(z) \end{Bmatrix} = \mp \frac{i}{k_{jc}}\left[\frac{1}{R(z)} - \begin{Bmatrix} T_{jc}(z) \\ \tilde{T}_{jc}(z) \end{Bmatrix} e^{ik_{jc}z}\right]; \tag{24}$$

$$\begin{Bmatrix} T_{jc}(z) \\ \tilde{T}_{jc}(z) \end{Bmatrix} = \int_{\pm z}^{\infty} \frac{\xi e^{\mp ik_{jc}\xi}}{R^3(\xi)}\, d\xi; \tag{25}$$

$$\frac{1}{R^n(z)} = \frac{1}{\pi} \int_0^{\pi} \frac{d\vartheta}{[z^2 + 4a^2 \sin^2(\vartheta/2)]^{n/2}}. \tag{26}$$

Substituting (24) into (8), (15), and (13) we get

$$\varphi(z) = \Phi_e[\sigma, z]/\varepsilon_{0c}(\omega) - Q[\sigma, z] + \Phi_i[\sigma, z], \tag{27}$$

where

$$\varepsilon_{0c}(\omega) = 1 - \omega_0^2/\omega(\omega + i\nu) \tag{28}$$

is the dielectric constant of stationary collisional plasma;

$$Q[\sigma, z] = \int_{-L}^{l} \sigma(\zeta)\, Q_c(z - \zeta)\, d\zeta, \quad Q_c(z) = k_0^2 \begin{cases} Q(z), & z > 0 \\ \tilde{Q}(z), & z < 0 \end{cases}; \tag{29}$$

$$\begin{Bmatrix} Q(z) \\ \tilde{Q}(z) \end{Bmatrix} = \frac{1}{k_{0c}} \sum_{j=1}^{2} \frac{(-1)^j}{k_{jc}} \begin{Bmatrix} T_{jc}(z) \\ \tilde{T}_{jc}(z) \end{Bmatrix} e^{ik_{jc}z}. \tag{30}$$

We note that in the integral representation (23) of function $1/R(z-\zeta)$ there is a quantity $[(z-\zeta)^2 + 4a^2 \times \sin^2(\vartheta/2)]^{1/2}$, i.e., the distance between two points located on the surface of the cylinder, where ϑ is the angle between the planes passing through these points and the axis of the cylinder. Thus geometrically $R(z-\zeta)$ is the mean distance along the azimuth between the source point and the observation point. Formula (24) represents the expansion of functions $S_{jc}(z)$ and $\tilde{S}_{jc}(z)$ in power series of $1/R(z)$.

Taking into consideration the asymptotic form of Macdonald function $K_0(x)$ for small and large values of the argument in the formal limits $u \to \infty$ and $u \to 0$ ($v_{Te} \to 0$), respectively, from formulas (13) and (27) we get the usual expressions for the potential of a cylindrical conductor in vacuum and in a collisional plasma at rest. If $\omega_0 = 0$, from (13) we again get the "vacuum" expression for the potential of the wire. Finally, for $\nu = 0$ formulas (13) and (27) yield the well-known results for the case of a moving collisionless plasma [2].

Now using the boundary condition

$$\varphi(z) = \Phi_0, \quad \Phi_0 = \text{const} \tag{31}$$

and the corresponding estimate [3] for functional $\Phi_e[\sigma, z]$ under the usual assumptions of the theory of thin antennas we write the approximate equation for the charge distribution along the conductor:

$$\sigma(z) + \mu\Phi_i[\sigma, z] = \mu\Phi_0, \quad \mu = 1/2B. \tag{32}$$

An antenna equation of the same type is obtained if in solution (4) of Poisson's equation we immediately put $J_0(k_\perp a) = 1$ ($a \to 0$) for the thin conductor. In this case in expressions (21)-(26) a quantity $R_0(z-\zeta) = [(z-\zeta)^2 + a^2]^{1/2}$ will be present instead of $R(z-\zeta)$. At the same time, the transition from $R(z-\zeta)$ to $R_0(z-\zeta)$ is the basic approximation of the theory of thin antennas (for example, see [7] and the literature cited therein). The replacement of R by R_0 is usually done without the required justification similar to the one presented above and is mathematically incorrect, since it means that in the equation of the source the regular kernel R_0 is substituted for the singular kernel R. On the other hand, this approximation leads to a correct (i.e., the same as in the case of conductor in vacuum [4]) logarithmic singularity in the charge distribution over the antenna (see below).

Thus, kernel $G_c(r, z-\zeta)$ of integral operator (6) is "insensitive" to the replacement of Bessel function $J_0(k_\perp a)$ by unity or equivalently to the replacement of nonnegative function $[J_0(k_\perp a)]^2$ in (4) (for $r = a$) by the alternating function $J_0(k_\perp a)$ in the approximation of the theory of thin antennas. The fact that this replacement is possible, indicates the correctness of the approach and the method of Green's function [2] developed here. It is completely justified to assume that this replacement is admissible even in more complex situations when it is not possible to trace explicitly and in fine details the transition from the expressions, obtained by retaining function $J_0(k_\perp a)$, to analogous formulas with $J_0(k_\perp a)$ replaced by unity, for example, in the computation of the impedance of TCA in a moving plasma taking into consideration the transverse radiation field or the "intrinsic" spatial dispersion of the medium.

Equation (32) is easily solved using Laplace transform defined in [2]; for $\nu \ll 2\omega_0$ the desired function $\sigma(z)$ is of the form

$$\sigma(z) = \mu\Psi_0 \left[\gamma_{0c} + \gamma_{1c} e^{ip_{1c}(L+z)} + \gamma_{2c} e^{ip_{2c}(L+z)}\right], \quad \gamma_{0c} = \frac{k_{1c} k_{2c}}{p_{1c} p_{2c}}, \tag{33}$$

$$\gamma_{jc} = \frac{(p_{jc} - k_{1c})(p_{jc} - k_{2c})}{p_{jc}(p_{1c} - p_{2c})}, \quad p_{jc} = (1/2)\{2k + \mu k_0 (A_{2c} - A_{1c})$$

$$+ ik_0 \pm k_0 [\mu^2 (A_{2c} - A_{1c})^2 - 4\mu (A_{1c} + A_{2c}) + 4]^{1/2}\}.$$

The quasistatic input impedance $Z = U/I_0$ of a symmetric antenna, for which

$$\int_{-L}^{0} \sigma(z)\, dz = -q; \quad \int_{0}^{L} \sigma(z)\, dz = q,$$

where the amplitude of the charge at each conductor is given and equal to q, U is the potential difference of the antenna whiskers, and the current in the gap

$$I_0 = I(0) = i\omega \int_{-L}^{0} \sigma(z)\, dz = -i\omega q$$

can be written in the following form:

$$Z = (i/\omega\mu)(J_1^{-1} + J_2^{-1}); \tag{34}$$

$$J_1 = \gamma_{0c} L + i(\delta_{1c} - \delta_{2c}),$$
$$J_2 = \gamma_{0c} L + i(\delta_{1c} e^{ip_{1c}L} - \delta_{2c} e^{ip_{2c}L}), \quad \delta_{jc} = (\gamma_{cj}/p_{jc})(1 - e^{ip_{jc}L}). \tag{35}$$

Formulas (33)-(35) show that the characteristics of the source in a moving collisional plasma do not have singularities at the plasma frequency. In particular, in the case of infrequent collisions ($\nu \ll \omega$) the components of the input impedance $Z^0 = R^0 - iX^0$ are equal at $\omega = \omega_0$:

$$R^0 \simeq (A_c/u)[1 - \cos(\alpha\mu A_c)], \quad X^0 \simeq -(A_c/u)\cos(\alpha\mu A_c)\operatorname{ctg}(\alpha\mu A_c/2),$$
$$\alpha = kL, \quad A_c = K_0(k_c L/2), \quad \nu/\omega \ll 2\mu A_c, \quad k_c L \ll 1. \tag{36}$$

Thus at the plasma frequency the amplitudes of the input resistance and reactance of the source are determined by the radius of the conductor, the electron collision frequency, and the flow velocity of the plasma, while the length of the antenna gives the period of oscillations of these quantities. If $\nu \to 0$ $Z^0 \to \infty$ [8], which is due to the intense build-up of the characteristic oscillations of the medium.

For the case of collisionless plasma and frequencies not too close to the plasma frequency ($|\omega - \omega_0| \gg \omega_0$), the impedance of TCA is given by formula (34) in which

$$J_1 = \frac{\varepsilon_0 L}{1 - \eta^2 v} - \frac{i\mu k_0 A_m}{\eta}\left[\frac{1 - e^{i(k+\eta k_0)L}}{(k + \eta k_0)^3} - \frac{1 - e^{i(k-\eta k_0)L}}{(k - \eta k_0)^2}\right],$$

$$J_2 = \frac{\varepsilon_0 L}{1 - \eta^2 v} - \frac{i\mu k_0 A_m}{\eta}\left[\frac{1 - e^{i(k+\eta k_0)L}}{(k + \eta k_0)^2} e^{i(k+\eta k_0)L} - (1 - e^{i(k-\eta k_0)L})\right.$$

$$\left.\times (k - \eta k_0)^{-2} e^{i(k-\eta k_0)L}\right], \quad \varepsilon_0 = 1 - v, \; v = \omega_0^2/\omega^2, \; \eta = (1 - 2\mu A_m)^{1/2},$$
$$A_m = K_0(k_m a), \quad k_m = \max\{k; k_0\}. \tag{37}$$

Furthermore, if the condition

$$2k_m L \sim 1 \tag{38}$$

is satisfied, then passing on to the limit $\eta \to 0$ in (37) we find that

$$Z = (2iB/u)\{[\varepsilon_0\alpha + iv(2 - (2 - i\alpha)e^{i\alpha})]^{-1} + [\varepsilon_0\alpha + iv(2 - i\alpha - 2(1 - i\alpha)e^{i\alpha})e^{i\alpha}]^{-1}\}. \tag{39}$$

Formula (39) for the impedance and the corresponding charge distribution along the conductor

$$\sigma(z) = (\Psi_0/2B)\{\varepsilon_0 + v[1 - ik(L + z)]\, e^{ik(L+z)}\} \tag{40}$$

can be obtained from the results of [1] in the isotropic limit.

In conclusion, we point out that the formulas for the characteristics of an antenna in a moving plasma, which have been presented here, hold if the investigated system is in a stable region of its parameters [1]. The stability of forced quasistatic oscillations of the current in a thin conductor embedded in a moving plasma is the subject of special investigation.

We note that the expressions given here differ from the corresponding formulas of [2] in that a boundary condition other than in [2] has been used. For this reason the current distribution along the antenna corresponding to charge (40) differs from the "triangular" (compare [2]).

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EFFICIENCY OF ADAPTIVE OPTICAL SYSTEMS IN TURBULENT ATMOSPHERE

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The efficiency of adaptive optical systems is estimated based on the interference criterion, which permits the analysis of both receiving and transmitting systems from a single viewpoint. It is shown that Strehl number, the signal-to-noise ratio in optical communication systems with amplitude and phase–frequency modulation – the quality of operation of adaptive interferometers, can be expressed in terms of the inferference criterion. Computations are carried out for systems with modal corrector of the wave front. The problems of optimization of the corrector are discussed.

1. The efficiency of operation of adaptive systems in a turbulent atmosphere has been investigated in a number of articles [1-4]. However, in an overwhelming majority of these receiving adaptive optical systems have been discussed. The quality of operation of these systems has been traditionally estimated from Strehl number [5].

In the present study the interference criterion [6] is used for estimating the efficiency of adaptive optics systems. This enabies one to analyze both receiving and transmitting optical systems including focusing systems with reference source (beacon) from a single viewpoint. It is shown that Strehl number, the signal-to-noise ratio in optical communication systems with amplitude and phase–frequency modulation – the quality of operation of adaptic interferometers can be expressed in terms of the interference criterion. The computations are carried out for systems with modal corrector of the wave front. It is assumed that an accurate reproduction of the first M Zernike polynomials describing the simplest aberration of the wave front [5] can be obtained with the use of the correction. The variance distribution of the uncompensated phase over the circular aperture of the adaptive system is computed for the first 10 Zernike polynomials. The results obtained here can also be used for the computation of segmented correctors and systems combining both modal and zonal correction of the wave front under atmospheric turbulence conditions. The problems of optimization of the corrector are discussed.

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