DISTORTION OF AN EXTERNAL MAGNETIC FIELD BY AN EXPANDING PLASMA SPHERE LOCATED IN A SLIGHTLY CONDUCTIVE SEMISPACE

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Expansion of an ionized gas in a cavity within a condensed medium located in a magnetic field leads to development of electromagnetic perturbations. Another source of electromagnetic fields is shock polarization of the material in the stress wave which develops within the condensed medium under the action of the expanding gas cavity [1]. It is of interest to clarify the roles of both these mechanisms and to compare their contributions to the total electromagnetic signal (amplitude, spectrum, polarization, etc.).

The present study will investigate a strongly heated plasma sphere, expanding within a weakly conductive condensed medium bounded by a vacuum. The effective magnetic moment of the plasma sphere will be calculated together with the electromagnetic field at the planar boundary with the vacuum. The dependence of the signal form and spectrum upon the plasma parameters and properties of the medium will be analyzed. The effect of the phenomenon studied will be compared to the contribution of shock polarization of the condensed medium.

We will consider a homogeneous plasma sphere located within a homogeneous magnetic field $\mathbf{H}_0$. At a time $t > 0$ let the radius of the sphere begin to increase by a law $R(t) = R_0\varphi(t)$, while the electrical conductivity of the plasma within the sphere changes by the law $\sigma = \sigma(t)$. If the sphere is located at a large distance from the surface $h \gg R$, then in calculating the field in the vicinity of the sphere we may neglect the effect of the free surface. We also then neglect the electrical conductivity σ_0 of the condensed medium ($\sigma \gg \sigma_0$) as well. The equations for the magnetic field in the quasisteady-state approximation (the medium is nonmagnetic, i.e., $\mu = 1$) have the form

$$\begin{gathered}\operatorname{rot}\mathbf{H} = 0,\ \operatorname{div}\mathbf{H} = 0,\ r > R;\\ \frac{\partial\mathbf{H}}{\partial t} - \operatorname{rot}[\mathbf{vH}] = \frac{c^2}{4\pi\sigma}\Delta\mathbf{H},\quad \operatorname{div}\mathbf{H} = 0,\quad 0 < r < R,\end{gathered}\tag{1}$$

where $\mathbf{v}$ is the velocity field within the plasma sphere. It is evident from the expressions presented below that as a solution of Eq. (1) for the external region $r > R$ it is sufficient to take the sum of the homogeneous field $\mathbf{H}_0$ and the magnetic moment field proportional to $\mathbf{H}_0$. Using an expression in reference vectors of a spherical coordinate system, we write this solution in the form

$$\mathbf{H} = H_0\left[\left(\frac{2\alpha R_0^3}{r^3} + 1\right)\cos\theta\mathbf{e}_r + \left(\frac{\alpha R_0^3}{r^3} - 1\right)\sin\theta\mathbf{e}_\theta\right],\quad r > R\tag{2}$$

(the angle θ is measured from the direction of $\mathbf{H}_0$). The form of the function $\alpha = \alpha(t)$ appearing in the magnetic moment $\alpha H_0 R_0^3$ will be determined below from the boundary conditions on the sphere surface.

Using an expression for the radius vector of a volume element of the homogeneously expanding plasma $\mathbf{r} = \mathbf{r}_0\varphi(t)$ (where $\mathbf{r}_0$ is the initial coordinate of the volume element), we write the plasma velocity as $\mathbf{v} = \mathbf{r}_0\dot{\varphi} = \mathbf{r}\dot{\varphi}/\varphi$. Substituting this expression in Eq. (1) and transforming to the Lagrangian variables $\mathbf{r}_0$, t, we obtain

$$\frac{\partial\mathbf{H}}{\partial t} + 2\frac{\dot{\varphi}}{\varphi}\mathbf{H} = \frac{c^2}{4\pi\sigma\varphi^2}\Delta_{\mathbf{r}_0}\mathbf{H},\quad \operatorname{div}_{\mathbf{r}_0}\mathbf{H} = 0.\tag{3}$$

We will omit the subscript on r_0, below, understanding by r the Lagrangian variable. We seek a solution of Eq. (3) in the form

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$$\mathbf{H} = H_1(r,\ t)\cos\theta\mathbf{e}_r - H_2(r,\ t)\sin\theta\mathbf{e}_\theta. \tag{4}$$

For H_1 and H_2 we find a system of equations

$$\begin{aligned}
&\dot{H}_1 + 2\frac{\dot{\varphi}}{\varphi}H_1 = \frac{c^2}{4\pi\sigma\varphi^2}\left[H_1'' + \frac{2H_1'}{r} - \frac{4(H_1 - H_2)}{r^2}\right],\\
&\dot{H}_2 + 2\frac{\dot{\varphi}}{\varphi}H_2 = \frac{c^2}{4\pi\sigma\varphi^2}\left[H_2'' + \frac{2H_2'}{r} + \frac{2(H_1 - H_2)}{r^2}\right],\\
&H_1' + 2(H_1 - H_2)/r = 0
\end{aligned} \tag{5}$$

(the prime denotes differentiation with respect to r, and the dot, differentiation with respect to time). In system (5) it will be convenient to transform to new unknown functions $f = (H_1 - H_2)/H_0$ and $g = (H_1 + 2H_2)/H_0 - 3$, which satisfy the new system of equations

$$\dot{f} + 2\frac{\dot{\varphi}}{\varphi}f = \frac{c^2}{4\pi\sigma\varphi^2}\left(f'' + \frac{2f'}{r} - \frac{6f}{r^2}\right); \tag{6}$$

$$\dot{g} + 2\frac{\dot{\varphi}}{\varphi}(g+3) = \frac{c^2}{4\pi\sigma\varphi^2}\left(g'' + \frac{2g'}{r}\right); \tag{7}$$

$$2f' + g' + 6f/r = 0. \tag{8}$$

Since at the initial moment a homogeneous magnetic field $\mathbf{H}_0$, existed everywhere, the initial conditions for f and g are as follows:

$$f(r,\ 0) = g(r,\ 0) = 0. \tag{9}$$

The normal and tangent components of $\mathbf{H}$ must be continuous on the sphere surface. Comparing Eqs. (2) and (4), we find the boundary conditions

$$f(R_0,\ t) = 3\alpha/\varphi^3,\ g(R_0,\ t) = 0. \tag{10}$$

We will show that the system of Eqs. (6)-(8) with boundary conditions (9), (10) has a unique solution. Expanding Eq. (8) in f and considering that f is finite as $r \to 0$, we obtain

$$f = -\frac{1}{2r^3}\int_0^r r_1^3\frac{\partial g(r_1,\ t)}{\partial r_1}\,dr_1. \tag{11}$$

We will prove that if g satisfies Eq. (7), then Eq. (11) satisfies Eq. (6). To do this we apply to Eq. (11) the operator $\partial/\partial t + 2\dot{\varphi}/\varphi$. Using Eq. (7), we obtain

$$\dot{f} + 2\frac{\dot{\varphi}}{\varphi}f = -\frac{c^2}{8\pi\sigma\varphi^2 r^3}\int_0^r r_1^3\frac{\partial}{\partial r_1}\left(g'' + \frac{2g'}{r_1}\right)dr_1.$$

Integrating by parts several times and applying Eq. (11), we can transform this expression to the right side of Eq. (6). Thus, having solved Eq. (7) for g with conditions (9), (10), and then substituting the solution in Eq. (11), we obtain the unknown functions satisfying the problem posed.

We write the solution of Eq. (7) in the form

$$g(r,\ t) = \frac{1}{r}\sum_{n=1}^{n}\gamma_n(t)\sin\frac{\pi n r}{R_0}. \tag{12}$$

The functions $\sin(\pi nr/R_0)$ satisfying Eq. (10) form a complete orthogonal system. Using Eq. (9) we have

$$\gamma_n(t) = \frac{6R_0(-1)^n}{\pi n\varphi^2(t)}\int_0^t\frac{d\varphi^2}{dt'}\exp\left(-\int_{t'}^t\frac{\pi n^2c^2}{4\sigma R_0^2\varphi^2}\,dt''\right)dt'. \tag{13}$$

Substituting Eqs. (12), (13) in Eq. (11) and integrating over r, we find

$$f(r,t) = -\frac{R_0^2}{2\pi^2 r^3} \sum_{n=1}^{\infty} \frac{\gamma_n(t)}{n^2} \left\{ \left[\left(\frac{\pi n r}{R_0} \right)^2 - 3 \right] \sin \frac{\pi n r}{R_0} + \frac{3\pi n r}{R_0} \cos \frac{\pi n r}{R_0} \right\}. \tag{14}$$

Now, using boundary conditions (10), we calculate the effective magnetic moment of the sphere

$$\alpha = -\frac{3\varphi(t)}{\pi^2} \sum_{n=1}^{\infty} \frac{1}{n^2} \int_0^t \frac{d\varphi^2}{dt'} \exp\left(-\int_{t'}^{t} \frac{\pi n^2 c^2}{4\sigma R_0^2 \varphi^2} dt'' \right) dt'. \tag{15}$$

Thus, Eqs. (2), (12)-(15) solve the problem posed.*

In the limit $\sigma \to \infty$ we obtain

$$\alpha = -\frac{3\zeta(2) R(t)}{\pi^2 R_0} \left[\frac{R^2(t)}{R_0^2} - 1 \right] = -\frac{R(t)}{2R_0} \left[\frac{R^2(t)}{R_0^2} - 1 \right].$$

This result corresponds to expulsion of the magnetic field under frozen conditions and complete absence of diffusion within the sphere. In the opposite limiting case $\sigma \to 0$ we have the obvious result $\alpha = 0$.

The low-frequency conductivity σ of a Lorentz plasma in which electron collisions predominate is given by the expression [2]

$$\sigma = \frac{4\sqrt{2}}{\pi^{3/2}} \frac{T^{3/2}}{Ze^2 L m^{1/2}}, \qquad r_D = \left(\frac{T}{4\pi n_e e^2} \right)^{1/2},$$
$$L = \begin{cases} \ln(r_D T / Ze^2), & Ze^2/\hbar u \gg 1, \\ \ln(r_D \sqrt{mT}/\hbar), & Ze^2/\hbar u \ll 1, \end{cases} \tag{16}$$

where Z_e is the ionic charge, L is the Coulomb logarithm, r_D is the Debye screening radius, m is the mass of the electron, n_e is the electron concentration, u is the mean relative velocity of electrons and protons, and the temperature T is measured in electrical units. If we consider the plasma an ideal gas expanding adiabatically with adiabatic index γ, then the expression $T\varphi^{3(\gamma-1)} = T_0$ is valid (T_0 is the initial plasma temperature). As a result the expression in the exponential within the integrand of Eq. (15) can be written in the form

$$\frac{\pi c^2 n^2}{4\sigma R_0^2 \varphi^2} = \frac{n^2}{\tau_d} \left(\frac{\varphi}{\varphi_m} \right)^{\nu}, \quad \tau_d = \frac{16\sqrt{2}}{\pi^{5/2}} \frac{R_0^2 T_0^{3/2}}{Ze^2 L c m^{1/2} \varphi_m^{\nu}}, \quad \nu = (9\gamma - 13)/2 \tag{17}$$

(φ_m is the maximum value of the function φ, which determines the final radius of the gas sphere). If we take $R_0 = 10^2$ cm, $T_0 = 100$ eV, Z = 2, L = 4, then the constant τ_d, which determines the diffusion time, proves equal to 0.5 sec. The characteristic sphere expansion time $\tau = 30$ msec $\ll \tau_d$. In the initial time period (i.e., $t \lesssim \tau$) when the electrical conductivity of the plasma is high, the magnetic field within the sphere is practically frozen. Therefore for rapid expansion of the sphere the field within decreaes and the plasma sphere takes on an effective magnetic moment. Thereupon ($t \sim \tau_d$) due to adiabatic cooling and retardation of the plasma motion the process of magnetic field diffusion back into the sphere from the external region becomes dominant, leading to relaxation of the magnetic moment.

To calculate the variable magnetic moment with consideration of the boundary between the two media we replace the plasma sphere by a variable magnetic dipole immersed to a depth h in the slightly conductive medium with electrical conductivity σ_0 (Fig. 1). The electromagnetic field components for the corresponding boundary conditions were obtained in [3] in Fourier representation. The presence in those expressions of a factor of the form

*V. I. Yakovlev has noted that after the substitution $\mathbf{H}_* = \varphi^2 \mathbf{H}$, $d\tau = dt(\sigma\varphi^2)^{-1}$ and introduction of the vector potential the original Eq. (3) can be solved with the aid of a Laplace transform. The quadratures obtained for α can be reduced to Eq. (15).

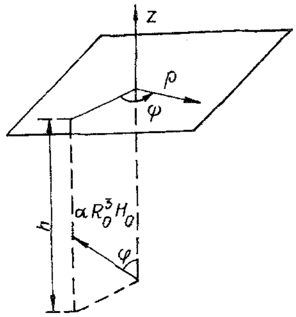

Fig. 1

$\exp(-h\sqrt{i\mu_0\sigma_0\omega})$ (where ω is the frequency) which considers decay of the electromagnetic field in the conductive medium shows that the low-frequency region of the spectrum is the most important. At the boundary with the vacuum (z = 0) in the near zone ($\omega\rho/c \ll 1$, $|\beta|\rho \gg 1$, $\beta = \sqrt{i\mu_0\sigma_0\omega}$) the field components have the form [4]

$$\begin{gathered}
E_\rho = -\frac{A}{2}\left(\frac{i\mu_0\omega}{\sigma_0}\right)^{1/2}\sin\varphi\sin\psi, \quad A = \frac{H_0R_0^3\alpha(\omega)}{\pi\rho^3}e^{-\beta h},\\
E_\varphi = A\left(\frac{i\mu_0\omega}{\sigma_0}\right)^{1/2}\left(\sin\varphi\cos\psi + \frac{3\cos\varphi}{2\beta\rho}\right),\\
\delta H_\rho = H_\rho - H_0\sin\varphi\cos\psi = A[\sin\varphi\cos\psi - 4\cos\varphi/(\beta\rho)],\\
\delta H_\varphi = H_\varphi + H_0\sin\varphi\sin\psi = A\sin\varphi\sin\psi/2,\\
\delta H_z = H_z - H_0\cos\varphi = A\left(\frac{9\cos\varphi}{2\beta^2\rho^2} - \frac{4\sin\varphi\cos\psi}{\beta\rho}\right).
\end{gathered} \tag{18}$$

When performing a reverse Fourier transform of Eq. (18) the integration range over frequency ω can be extended to infinity, since the contribution of high frequencies is insignificant in view of the factor $\exp(-\beta h)$ while the contribution of low frequencies is low, since $\alpha(\omega) \to 0$ as $\omega \to 0$. We will also neglect the field of the vertical component of the magnetic dipole, since the corresponding terms contain the small parameter $(\beta\rho)^{-1}$. Then the field components can be expressed in the following manner:

$$\begin{gathered}
\delta H_\varphi = HI_0, \ \delta H_\rho = 2\delta H_\varphi \operatorname{ctg}\psi,\\
E_\rho = -EI_1, \ E_\varphi = -2E_\rho \operatorname{ctg}\psi.
\end{gathered} \tag{19}$$

Here

$$\begin{gathered}
H = \frac{H_0R_0^3\sin\varphi\sin\psi}{(2\pi)^2\rho^3}, \quad E = H\left(\frac{\mu_0}{\sigma_0\tau_0}\right)^{1/2}, \quad \tau_0 = \frac{\mu_0\sigma_0h^2}{4},\\
I_0 = \int_{-\infty}^{\infty}\alpha(\omega)e^{-\beta h + i\omega t}d\omega, \quad \alpha(\omega) = \int_0^{\infty}\alpha(t')e^{-i\omega t'}dt',\\
I_1 = \sqrt{i\tau_0}\int_{-\infty}^{\infty}\sqrt{\omega}\,\alpha(\omega)e^{-\beta h + i\omega t}d\omega.
\end{gathered} \tag{20}$$

We extend a section from zero to infinity, with the integration contour lying on a band which satisfies the condition of attenuation at infinity and encloses the branching point ω = 0 from below. We substitute $\alpha(\omega)$ in the integrals I_0, I_1 and change the order of integration. By a replacement of variable the integrals over frequency in Eq. (20) can be reduced to the interval 0, ∞. Considering the signs of the integrands on the physical band we arrive at the integrals

$$I_0 = 2\int_0^{\infty}\alpha(t')\,dt'\int_0^{\infty}e^{-\sqrt{2\omega\tau_0}}\cos w\,d\omega, \ w = \sqrt{2\omega\tau_0} - \omega(t - t'), \tag{21}$$

$$I_1 = \sqrt{2\tau_0}\int_0^\infty \alpha(t')\,dt' \int_0^\infty e^{-\sqrt{2\omega\tau_0}}\sqrt{\omega}\,(\cos w + \sin w)\,d\omega. \tag{21}$$

Calculating the inner integrals in Eq. (21) [5] and considering that they are nonzero for $t' < t$, we obtain the expressions

$$I_0 = 2\sqrt{\pi\tau_0}\int_0^t \alpha(t-t')\,t'^{-3/2}\exp(-\tau_0/t')\,dt',$$

$$I_1 = -\sqrt{\pi\tau_0}\int_0^t \alpha(t-t')\,t'^{-5/2}(t'-2\tau_0)\exp(-\tau_0/t')\,dt'. \tag{22}$$

Equations (15), (19), (22) define the electromagnetic field of the plasma sphere on the boundary of the condensed medium with the vacuum.

For $t \gg \tau$ in the integral within the exponential of Eq. (15) we can substitute the maximum value $\varphi = \varphi_m$. Considering the rapid convergence of the series of Eq. (15) we will limit our examination to the first term. Then, integrating by parts, we obtain

$$\alpha(t) = -\frac{3\varphi(t)}{\pi^2}\left[\varphi^2(t) - e^{-t/\tau_d} - \frac{1}{\tau_d}\int_0^t \varphi^2 e^{(t'-t)/\tau_d}dt'\right]. \tag{23}$$

According to the estimate of Eq. (17), $\tau_d \gg \tau$, so that at $t \lesssim \tau$ the exponential in Eq. (15) is equal to unity to an accuracy $\sim\tau/\tau_d$. This approximation is also applicable to the representation of α in the form of Eq. (23), which is valid with the indicated limitation for all t. We approximate the function φ by the expression

$$\varphi = 1 + (\varphi_m - 1)[1 - \exp(-t/\tau)]. \tag{24}$$

Substituting Eq. (24) in Eq. (23) and neglecting factors $\sim\tau/\tau_d$ in the coefficients, we have

$$\alpha = -\frac{3(\varphi_m - 1)}{\pi^2}\left[\varphi_m(\varphi_m+1)e^{-t/\tau_d} - 2\varphi_m^2 e^{-t/\tau} + \right.$$
$$\left. + (1-\varphi_m^2)e^{-t/\tau_*} + 3\varphi_m(\varphi_m-1)e^{-2t/\tau} - (\varphi_m-1)^2 e^{-3t/\tau}\right],\quad \tau_*^{-1} = \tau^{-1} + \tau_d^{-1}. \tag{25}$$

We will study the behavior of the integrals I_0, I_1 of Eq. (22) for various ranges of t. If $t \ll \tau$, then, using the linear expansion of Eq. (25), we obtain

$$\delta H_\varphi = \frac{24H(\varphi_m-1)\tau_0}{\pi^{3/2}\tau}\left[\frac{(\varphi_m+1)\tau}{2\tau_d} - 1\right]\left[\frac{\sqrt{\pi}}{2}\left(2 + \frac{t}{\tau_0}\right)\operatorname{erfc}\left(\sqrt{\frac{\tau_0}{t}}\right) - \right.$$
$$\left. - \sqrt{\frac{t}{\tau_0}}\exp\left(-\frac{\tau_0}{t}\right)\right],\quad E_\rho = \frac{24E(\varphi_m-1)\tau_0}{\pi^{3/2}\tau}\left[1 - \right.$$
$$\left. - \frac{(\varphi_m+1)\tau}{2\tau_d}\right]\left[\sqrt{\frac{t}{\tau_0}}\exp\left(-\frac{\tau_0}{t}\right) - \sqrt{\pi}\operatorname{erfc}\left(\sqrt{\frac{\tau_0}{t}}\right)\right]. \tag{26}$$

The character of Eq. (26) is determined by the parameter τ_0. For example, at h = 300 m, $\sigma_0 = 10^{-2}\ \Omega^{-1}\cdot m^{-1}$ $\tau_0 \sim 3\cdot 10^{-4}$ sec $\ll \tau$. Analysis of the commencement of the signal ($t \ll \tau_0$) shows that the increase occurs by laws $\delta H_\varphi \sim t^{5/2}\exp(-\tau_0/t)$, $E_\rho \sim t^{3/2}\exp(-\tau_0/t)$, while at $t \gg \tau_0$ the character of the time dependence of δH_φ is practically linear, $E_\rho \sim \sqrt{t}$ (at $t \gg \tau_0$ the functions $\operatorname{erfc}(\sqrt{\tau_0/t})$ and $\exp(-\tau_0/t)$ can be taken equal to unity).

For the region $t \gg \tau$ substitution of Eq. (25) in Eq. (22) leads to typical integrals I_2, I_3 which can be evaluated in the following manner:

$$I_2 = \int_0^t \exp\left[-\left(\frac{\tau_0}{t'} + \frac{t-t'}{\tau}\right)\right]\frac{dt'}{t'^{3/2}} \approx \frac{\tau}{t^{3/2}} + \sqrt{\frac{\pi}{\tau_0}}\exp\left(-\frac{t}{\tau}\right), \tag{27}$$

$$I_3 = \int_0^t \exp\left[-\left(\frac{\tau_0}{t'} + \frac{t-t'}{\tau}\right)\right]\left(\frac{t'-2\tau_0}{t'^{5/2}}\right) dt' \approx \frac{\tau}{t^{3/2}}. \tag{27}$$

The first term on the right side of Eq. (27) considers the contribution of the region $\tau \lesssim t' < t$, while the second in the expression for I_2 defines the integration region $0 < t' \lesssim \tau$, which is small for I_3 and has thus been omitted. Using the estimate of Eq. (27), for the range $\tau \ll t \ll \tau_d$ we find the expressions

$$\begin{gathered} \delta H_\varphi = H\left[2\pi\alpha(t) + \frac{\tau\sqrt{\tau_0}}{(\pi t)^{3/2}}(\varphi_m - 1)(11\varphi_m^2 + 5\varphi_m - 4)\right], \\ E_\rho = \frac{(\varphi_m - 1)E}{2\pi^{3/2}}\sqrt{\frac{\tau_0}{t}}\left[12\varphi_m(\varphi_m + 1) + \frac{\tau}{t}(11\varphi_m^2 + 5\varphi_m - 4)\right] \end{gathered} \tag{28}$$

($\alpha(t)$ is determined by Eq. (25)). As is evident from Eq. (28) the signal begins to fall, so that the substitution $t \sim \tau$ in Eq. (28) gives a coarse estimate of the maximum value of the electromagnetic field components: $|\delta H_\varphi|_m \sim H_0 R_0^3 \varphi_m^3/(2\pi\rho)^3$, $|E_\rho|_m \sim |\delta H_\varphi|_m [\mu_0/(\sigma_0\tau)]^{1/2}$. Taking the parameter values indicated above and setting $H_0 \sim 50$ A/m, $\varphi_m \sim 30$, we have $|\delta H_\varphi|_m \sim 5 \cdot 10^3/\rho^3$ A/m, $|E_\rho|_m \sim 10^{10}/\rho^3$ μV/m (with ρ in meters).

The drop in signal at $t \gg \tau_d$ is given by

$$\begin{gathered} \delta H_\varphi = -\frac{6H}{\pi}\varphi_m(\varphi_m^2 - 1)\left[\sqrt{\frac{\tau_0}{\pi}}\frac{\tau_d}{t^{3/2}} + \exp(-t/\tau_d)\right], \\ E_\rho = -\frac{3E\varphi_m(\varphi_m^2 - 1)}{\pi^{3/2}}\frac{\sqrt{\tau_0}\,\tau_d}{t^{3/2}}. \end{gathered} \tag{29}$$

It is evident from comparison of Eqs. (28) and (29) that E_ρ changes polarity. The maximum modulus of reverse polarity E_ρ (at $t \sim \tau_d$) is less than $(\tau_d/\tau)^{1/2}$ times the initial peak. While the characteristic width of the first peak $\sim\tau$, the reverse polarity peak is significantly wider, since it depends on τ_d. The onset of the signal is determined by rapidly occurring plasma sphere expansion processes. The subsequent slower fall is caused by relaxation of the effective magnetic moment as a consequence of magnetic field diffusion into the gas cavity.

Results of electromagnetic signal calculations are shown in Fig. 2a, b for $\varphi = \pi/2$, $\psi = \pi/6$, $\rho = 3$ km, $h = 250$ m, $\varphi_m = 30$, with curves 1-4 corresponding to $\tau_d = 0.2$, 0.2, 0.5, 0.5 sec and $\tau = 0.03$, 0.02, 0.03, 0.02 sec; the dashed lines are calculations with the approximate analytic expressions (28), (29). The signal form and amplitude correspond to those observed in experiment [6]. At the same time in the experiments the polarization of $\mathbf{E}$ and $\delta\mathbf{H}$ in some cases corresponded to the field of a magnetic dipole, and in others, to the field of an effective electric dipole which develops upon shock electrical polarization of the condensed medium. This indicates that the effect is caused by both mechanisms. The analysis performed in [1] and the present study show that the amplitudes of E_ρ and δH_ρ caused by the vertical component of the electric dipole and the horizontal component of the magnetic dipole lie within the same order of magnitude, i.e., are comparable. The components E_z, δH_z are probably related to shock polarization of the medium, while E_φ, δH_φ are caused by disturbance of $\mathbf{H}_0$ by the plasma sphere. The orientation of the electric dipole related to asymmetry of the shock wave front is of a random character [1]. Therefore in some experiments the contribution of this mechanism is negligibly small. At the same time the weakness of the effect produced by the other mechanism (if the magnetic field is vertical) can be predicted a priori. The difference in these effects also manifests itself during the time of signal falloff. For shock polarization the signal duration is determined by the time required for development of the destruction wave or the characteristic time for relaxation of the condensed medium polarization. The perturbations produced by expansion of the plasma sphere have a duration of the order of magnitude of the time required for diffusion of the external magnetic field into the gas cavity. According to the estimates made above this time is the largest.

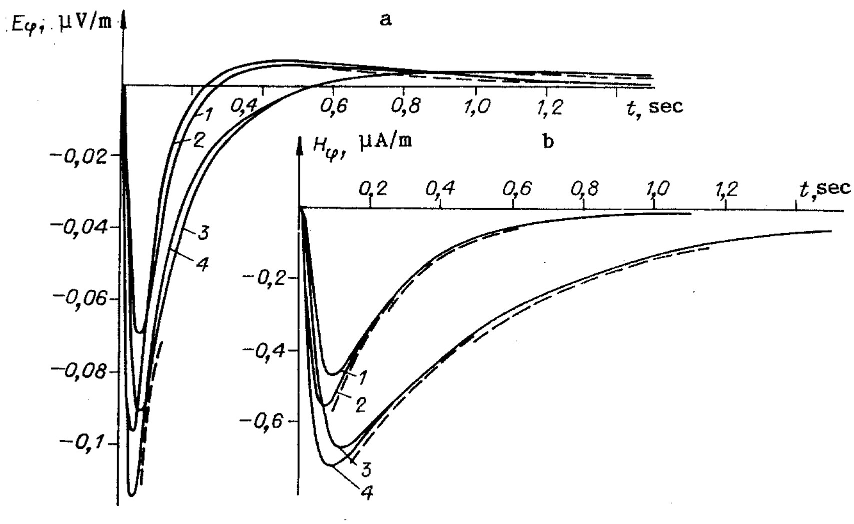

Fig. 2

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