tensor in the case when $\theta_1 \gg m_e/\sqrt{s}$ and $\pi - \theta_1 \gg m_e/\sqrt{s}$ take the following form for the case of ρ-meson production

$$r_{ik} = \alpha_{ik} - D\beta_{ik}/(d\sigma/dx), \tag{5}$$

where $d\sigma/dx$ and D are given by formulas (2) and (3), and the tensors α_{ik} and β_{ik} are given by

$$\alpha_{11} = \alpha_{22} = 2/3; \quad \alpha_{12} = \alpha_{13} = \alpha_{23} = \beta_{12} = \beta_{23} = 0,$$

$$\beta_{11} = 2\frac{1 + x^2/y^2}{1 - x^2}; \quad \beta_{22} = 2\frac{1 + x^2y^2}{y^2(1 - x^2)}; \quad \beta_{13} = 4\frac{xm_\rho\sqrt{s}}{y^2(s + m_\rho^2)\sqrt{1 - x^2}} \tag{6}$$

with γ given by formula (3). In the liniting case $s \gg m_\rho^2$, we obtain from (5) and (6) $r_{11} - r_{22} = -1/3$ and $r_{13} = 0$. This means that the ρ-meson spin is fully aligned parallel and antiparallel to its momentum. Such a result can be attributed to the fact that when $s \gg m_\rho^2$ it is possible to neglect the mass of the ρ meson and its polarization state should coincide with the polarization state of the photon.

4. In view of the fact that the lifetime of the vector mesons is short, let us consider the annihilation of an electron and positron into two π or K mesons and a photon. If we disregard the hypothetical σ meson discussed in [5 - 7] and remain within the framework of the C-invariant theory, then it suffices, in the energy region under consideration, to take into account only the contributions of the diagrams in which the photon is emitted from the electron and positrol lines. We introduce the angle θ_2 between the momenta of the pion and the initial electron in the c.m.s. of the two pions. If we introduce $z = \cos\theta_2$ and designate by ω^2 the mass of the system of two pions, then it is easy to show that at all values of ω^2 the angular distributions of the pions is determined by the factor $1 - z^2$. Integrating the angular distribution with respect to z, we obtain the differential cross section in the approximation in which $\ln(s/m_e^2) \gg 1$, in the form

$$\frac{d\sigma}{d\omega^2} = \frac{\alpha^3(\omega^2 - 4m_\pi^2)^{3/2}(s^2 + \omega^4)m_\rho^4\ln(s/m_e^2)}{3s^2(s - \omega^2)\omega^4\sqrt{\omega^2}[(\omega^2 - m_\rho^2)^2 + m_\rho^2\Gamma_\rho^2]}. \tag{7}$$

If we take into account the relation $\omega^2 = s - 2\sqrt{s}\varepsilon_\gamma$, where ε_γ is the energy of the γ quantum in the c.m.s. of the initial particles, then it is easy to obtain from (7) the energy distribution of the γ quanta in the processes considered by us, observation of which would make it possible to investigate vector mesons with colliding beams at higher energies than used in [1, 2].

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NONLINEAR PASSAGE OF ELECTROMAGNETIC WAVES THROUGH A PURE METAL

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The purpose of the present work is to call attention to the possibility of a unique nonlinear passage of electromagnetic waves through a metal. This effect is connected with the free flight of electrons from one surface of a metal plate to the other, and is analogous to the plasma-echo effect investigated in a number of works. In particular, we have calculated [1] the plasma echo for the case of transverse electromagnetic waves in a plasma. The phenomenon dealt with in the present paper is analogous to that considered in [1]. It exists if the electron mean free path ℓ in the metal is much larger than the plate thickness d

$$\ell >> d \tag{1}$$

and if the reflection of the electrons from the metal surface is at least partly specular.

The experiment can be set up in a number of ways. For concreteness, we consider the following scheme: Assume that an electromagnetic wave $\vec{E}_1$ of frequency ω_1 is incident from the left on a metal plate. The field frequency is assumed high enough to produce the normal skin effect. More accurately speaking, we assume that the following inequalities hold:

$$\frac{v_0}{\omega_1} << \delta << d << \ell, \tag{2}$$

where δ is the depth of penetration of the field into the metal and v_0 is the electron velocity on the Fermi surface. If condition (2) is satisfied, we can write for δ

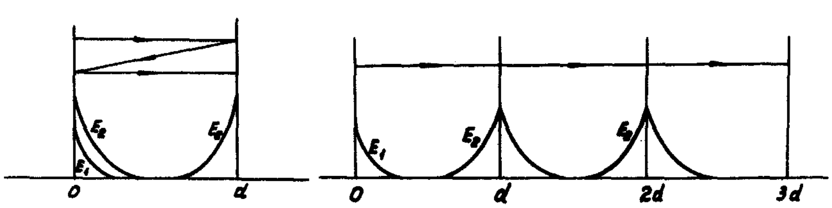
$$\delta \approx c/\omega_0,$$

where $\omega_0 = 4\pi Ne^2/m^*$ is plasma frequency of the metal and m* is the effective mass of the electrons. We neglect throughout the anisotropy of the crystal, and assume in particular that the Fermi surface is a sphere. Since $\delta << d$, the field E_1 does not penetrate through the plate. However, if we now apply on both sides of the plate a field E_2 of frequency ω_2, then a field E_e of frequency ω_1 will appear on the right side of the plate. The amplitude of the field E_e will be proportional to $E_1E_2^2$, and will have a sharp maximum at $\omega_2 \simeq 3\omega_1$. We note beforehand also that under condition (1) the field E_e will not decrease with increasing plate thickness d, but will increase.

We proceed to a quantitative calculation of the phenomenon. We must solve the collisionless kinetic equation for the electrons in the metal, accurate to terms of order $E_1E_2^2$. We ssume that the scattering of the electrons by the surface will have in the main a diffuse character with small but finite degree fo specularity $\beta << 1$.

The effect of interest to us results from electrons that are perturbed by the field E_1 at the left surface of the plate (surface 1), and are then specularly reflected from the right surface 2 and are again specularly reflected from surface 1. The number of such electrons is β^2 times the total number.

It is easily understood that for such electrons it is possible to replace their real motion with specular reflection by straight-line motion in fictitious (specularly continued) fields, as shown in the figure. Figure a shows the true field configuration, and Fig. b the configuration of the continued fields. The solution of the kinetic equation is then obtained in exactly the same manner as in [1].



a b

The resultant current density with frequency ω_1 near the surface 2 is equal to

$$J_y = i\beta^2 \frac{3}{2\pi^2} \frac{e^4 E_1 E_2^2}{\hbar^3} \frac{(2\omega_1 + \omega_2) d^2 v_0^3}{\omega_2^6 \delta^4} \times \tag{3}$$

$$\times \int_0^1 [1 - 2u^2 + \frac{3}{2} u^4 - u^4 \ln u^2] e^{i \frac{\omega_1}{v_0 u} x + i \frac{\omega_2}{v_0 u} d} \frac{du}{u}.$$

When $\omega_2 \simeq 3\omega_1$ the current density differs from zero only in a thin layer with thickness on the order of $r_1 = v_0/\omega_1 << \delta$ near the surface 2. Under this condition, the field E_e is determined by the formula

$$E_e = - \frac{2\pi\omega_1}{c^2} e^{-\frac{d-x}{\delta_1}} \delta_1 \int_{-\infty}^{d} i(x) dx. \tag{4}$$

Here we encounter another deviation from the results of [1], connected with the fact that the metal is bounded. In unbounded space the limits of the integral in (4) would be $-\infty$ and ∞. Such an integral vanishes and the approximation employed in (4) would not suffice.

As a result, at x - d, the amplitude of the transmitted field is given by

$$E_e = \frac{1}{243\pi} \beta^2 \frac{e^4 E_1 E_2^2}{c^2 \hbar^3} \frac{v_0^4 d^2}{\delta^3} \frac{1}{\omega_1^5} \times \tag{5}$$

$$\times \int_0^1 [1 - 2u^2 + \frac{3}{2} u^4 - u^4 \ln u^2] e^{i \frac{\omega_2 - 3\omega_1}{v_0 u} d} \frac{du}{u}.$$

Estimates show that the effect may not be small even at sufficiently small values of β. Thus, for example, for tin at $\omega_1 \sim v_0 \omega_0 / c$ we get $E_e \simeq 1.4 \times 10^3 \beta^2 E_2^2 d^2 E_1$, where d is in centimeters and E_2 in cgs esu. It is seen from (5) that the effect differs from zero when $|3\omega_1 - \omega_2| \lesssim v_0/d$. It is interesting that the nonlinear transparency of the plate increases with the plate thickness.

The authors thank A.F. Andreev for useful discussions.

[1] M.P. Kemoklidze and L.P. Pitaevskii, Zh. Eksp. Teor. Fiz. 58, 1853 (1970) [Sov. Phys.-JETP 31, No. 6 (1970)].

E R R A T A

In the article by V.V. Apollonov, Yu.A. Bykovskii, N.Y. Dgtyarenko et al., Vol. 11, No. 8, page 254, formula (1) should read

$$\epsilon = \left[\frac{3}{2} kT(1 + z) + \sum^{z} I_i \right] N = \frac{P}{\rho} \frac{1}{(\gamma' - 1)}.$$

In the same page, at the end of the next formula, read a $\lesssim$ 3 in place of a $\lesssim$ z.