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OPTOELECTRICAL EFFECT AND OPTOKINETIC DIA- AND PARAMAGNETISM

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1. A flux of electromagnetic waves (which we shall call optical) produces in a conducting medium placed in a magnetic field $\vec{H}$, which is not parallel to this flux, effects similar to the thermomagnetic effects [1, 2]. In particular, an electric current may be produced, with a density $\vec{j}$ given by

$$\mathbf{j} = \chi \mathbf{I}_k + \chi_1 [\mathbf{I}_k \times \mathbf{H}] + \chi_2 \mathbf{H}(\mathbf{I}_k \cdot \mathbf{H}), \tag{1}$$

where $\vec{I}_k$ is the Poynting-vector component parallel to the wave vector $\vec{k}$. This current can enhance (paramagnetism) or weaken (diamagnetism) the magnetic field inside the medium (optokinetic dia- and paramagnetism).

2. In the presence of an electromagnetic wave, the distribution function is $F = f_0 + f_1 + f_2 + f_3 + \ldots$, where f_0 is the equilibrium distribution function, f_1 and f_3 are proportional to the frequency ω of the wave and to double the frequency, respectively, and f_2 is independent of the frequency (f_2 and f_3 are quadratic in the fields of the wave).

We are interested in the functions f_1 and f_2; we supplement the corresponding kinetic equations, which are written out in [3], with terms that depend on the external magnetic field. We consider the case of a weak magnetic field ($\mu H/c < 1$, where μ is the mobility), and therefore confine ourselves to the calculation of the coefficient χ_1 in the term that is linear in H. The increment proportional to H in the distribution function f_2 is given by

$$f_2^{(1)} = \frac{e^3 \tau^3 y}{2m^2 c^2} \frac{\partial f_0}{\partial \epsilon} \mathrm{Re}\Big(\{ y(1 - x^2 - 2x \,\mathrm{tg}\, \phi)(\mathbf{H}_1^* \mathbf{H})\mathbf{E}_1 - (1 + x \,\mathrm{tg}\, \phi) \times$$

$$\times [[\mathbf{E}_1 \mathbf{H}_1^*] \mathbf{H}] + c\tau_1 y [x + x_1 - (x x_1 - 1)\,\mathrm{tg}\, \phi](\mathbf{E}_1 \mathbf{E}_1^*)[\mathbf{k} \times \mathbf{H}]\} \mathbf{v}\Big) + \frac{e^3 \tau_2^2}{2mc} \times$$

$$\times \Big(\Big\{ 2(\mathbf{E}_1 \mathbf{v})(\mathbf{k}' \mathbf{v})[\mathbf{E}_1^* \times \mathbf{H}] + (\mathbf{E}_1 \mathbf{v})(\mathbf{E}_1 \mathbf{v})[\mathbf{k}' \times \mathbf{H}] + 2\frac{\tau}{\tau_3} v^2 (\mathbf{E}_1 \times \mathbf{E}_1^*)[\mathbf{k} \times \mathbf{H}] \Big\} \mathbf{v}\Big) \times$$

$$\times \frac{\partial}{\partial \epsilon} \left\{ \tau \tau_1 y y_1 [x + x_1 - (x x_1 - 1)\,\mathrm{tg}\, \phi] \frac{\partial f_0}{\partial \epsilon} \right\},$$

$$\mathbf{k}' = \mathrm{Re}\,\mathbf{k}, \quad x = \omega\tau, \quad y = (1 + \omega^2 \tau^2)^{-1},$$

x_1 and y_1 are the same but with τ replaced by τ_1, and $\vec{v}$ is the carrier velocity. By averaging $f_2^{(1)}$ over the polarization directions and calculating the current, we obtain the coefficient χ_1:

$$\chi_1 = B\left\{\frac{1}{2} < \tau^3 y^2 (3 + x^2 + x^3 \operatorname{tg}\phi)> + < \zeta \tau_1 \tau_2^2 x y y_1 [x + x_1 - (x x_1 - 1)\operatorname{tg}\phi]>\right\}, \tag{2}$$

where

$$\zeta(\epsilon) = 1 + 10\frac{\tau}{\tau_3} - \frac{\tau^2}{\tau_2^2} + \frac{4}{5}\frac{\partial \ln \tau_2}{\partial \ln \epsilon} + 4\frac{\tau}{\tau_3}\frac{\partial \ln\left(\tau_2 \frac{\tau}{\tau_3}\right)}{\partial \ln \epsilon},$$

$$< \Psi(\epsilon) > = \frac{\int_0^\infty \frac{\partial f_0}{\partial \epsilon}\epsilon^{3/2}\Psi(\epsilon)\,d\epsilon}{\int_0^\infty \frac{\partial f_0}{\partial \epsilon}\epsilon^{3/2}\,d\epsilon}, \qquad B = \frac{4\pi e^4 N}{m^3 c^3},$$

N is the carrier density, ϕ is the phase shift between $\vec{E}_1$ and $\vec{H}_1$, while tan ϕ and the coefficient χ were calculated in [3]. The relaxation times τ_1, τ_2, and τ_3 for certain scattering mechanisms are also given in [3]. It should be noted that in the presence of an external magnetic field, the coefficient χ calculated in [3] will also make a contribution $\chi_1^{(0)}$ to the coefficient χ_1. In the low-frequency region ($\omega\tau < 1$) this contribution equals $\chi_1^{(0)}$ if $\omega > \omega_0^2\tau$ and $\chi_1 = -(1/2)B<\tau^3>$ if $\omega < \omega_0^2\tau$, where $\omega_0 = (4\pi e^2 N/m)^{1/2}$ is the plasma frequency. In the low-frequency region ($\omega\tau > 1$) the value of χ_1 is smaller than in (2) by a factor $\omega\tau$.

In the high frequency case ($\omega\tau > 1$), χ_1 is given by

$$\chi_1 = \frac{B}{2\omega^2}\left\{< \tau > + \omega < \tau^2 > \operatorname{tg}\phi + 2 < \zeta(\epsilon)\tau_2^2 \frac{\tau + \tau_1}{\tau\tau_1} > - 2\omega < \zeta(\epsilon)\tau_2^2 > \operatorname{tg}\phi\right\}. \tag{3}$$

For the Coulomb scattering mechanism we have

$$\chi_1 \approx -\frac{B}{2\omega^2}\left\{3 < \tau > + \frac{1}{7}\omega < \tau^2 > \operatorname{tg}\phi\right\}. \tag{4}$$

From (4) we see that in this case the coefficient χ_1 reverses sign.

3. In such stars as the sun, the main magnetic field is toroidal. It is generated by convective motions due to the superadiabatic temperature gradient that occurs in the region of partial ionization of hydrogen. We then obtain approximately for such stars from Maxwell's equations (recognizing that for most stars $\mu H/c < 1$)

$$\frac{\partial \ln H}{\partial r} \approx \frac{4\pi}{c}\int_0^\infty \left[\chi_1(\omega) - \frac{\sigma_1}{\sigma}\chi(\omega)\right]\tilde{I}(\omega)\,d\omega, \tag{5}$$

where $\tilde{I}(\omega)$ is the spectral density of the energy flux, and σ and σ_1 are the ordinary and Hall conductivities. Equation (5) with such coefficients is applicable from the center of the star to the internal boundary r_1 of the convection region, since the generation of the magnetic field takes place further on. The magnetic field in the convection region is enhanced also by the optokinetic paramagnetism, but in accordance with a different law. For the Coulomb scattering mechanism, Eq. (5) takes the form

$$\frac{\partial \ln H}{\partial r} \approx -4\pi\left\{\frac{3}{2}\langle r \rangle + 14\frac{\langle r^2 \rangle}{\langle r \rangle}\right\}\frac{Bh^2}{cT^2}\frac{L}{4\pi r^2}, \tag{6}$$

where L is the luminosity. Substituting τ, we get

$$\frac{\partial \ln H}{\partial r} \approx -\frac{3\pi \cdot 10^4}{2\lambda}\frac{\hbar^2}{m^{5/2}c^4T^{1/2}}\frac{L}{4\pi r^2},$$

where λ is the Coulomb logarithm. Assuming that $T \approx T_0(r/r_1)^{-\alpha}$, where T_0 is the temperature of the internal convection region, in ergs, and $\alpha > 2$, we find that the magnetic field increases in the interior of the star in accordance with the law

$$H \sim \exp\left\{-\left(\frac{r}{\delta}\right)^{\frac{\alpha}{2}-1}\right\},$$

where the characteristic length δ is given by

$$\delta \approx \left(\frac{\alpha}{2} - 1\right)\left[\frac{2\cdot 10^{-4}\lambda}{3\pi}\frac{m^{5/2}c^4T_0^{1/2}}{\hbar^2}\frac{4\pi r_1}{L}\right]^{\frac{2}{\alpha-2}} r_1 .$$

Assuming that near r_1 we have $\alpha \approx 4$ and $T_0 \approx 5 \times 10^{-11}$ erg (values typical of the sun), we obtain the estimate $\delta \approx 10^{-1}R$, where R is the radius of the sun. For stars hotter than the sun, the characteristic length is even smaller.

This enhancement of the magnetic field is produced by currents that flow along meridians and return through the star's rotation axis. These currents produce their own magnetic field, having also a toroidal structure and coinciding in direction with the magnetic field produced by the convection. Its maximum lies approximately near the center of the meridian-axis current contour, as a result of which the main field is also enhanced in the interior of the star up to this line, beyond which the enhancement decreases.

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FLUCTUATION SPECTRUM IN SUPERCONDUCTING POINT JUNCTIONS

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1. The question of fluctuations in superconducting point junctions in which the Josephson effect takes place is of great interest, since in most applications, the limiting characteristics are determined precisely by the fluctuations. It is important here to determine the fluctuation spectrum, particularly the values of the spectral density at low frequencies and the Josephson-generation