owing to symmetry violation [8]. This makes it possible to obtain, in principle, generation without an external field to induce the transitions.

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INFLUENCE OF CARRIER DRIFT ON THE PROPAGATION OF ELECTROMAGNETIC WAVE IN A SOLID-STATE PLASMA

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Moss and co-workers [1], using a ring laser, succeeded in measuring with high accuracy the difference between the refractive indices of light propagating in a semiconductor parallel and antiparallel to the current lines. They associated the effect observed by them with the Fresnel dragging of light by the electrons of the semiconductor. By the same token, according to [1], this was the first observation of a relativistic effect caused by carriers moving relatively slowly in the direction of the field. A phenomenological theory of the Fresnel dragging, with the Doppler-effect contribution taken into account, was presented in [1] in accord with the assumed model.

It seems to us, however, that a more rigorous analysis should be based on the kinetic equation and Maxwell's equations, without using any model concepts for dragging. Indeed, if spatial disperson and drift of the electron gas are taken into account, the dielectric tensor $\varepsilon_{\mu\nu}(\omega, \vec{k})$ of the semiconductor does not satisfy the Onsager symmetry relation $\varepsilon_{\mu\nu}(\omega, \vec{k}) = \varepsilon_{\mu\nu}(\omega, -\vec{k})$. In the case of weak spatial dispersion, the expansion of $\varepsilon_{\mu\nu}(\omega, \vec{k})$ in powers of $\vec{k}$ in the presence of a constant electric field $\vec{F}$ contains linear terms (cf., e.g., [2]). The presence of these terms determines the difference between the conditions of electromagnetic wave propagation for $\vec{k} \uparrow\downarrow F$ and $\vec{k} \uparrow\uparrow \vec{F}$.

To estimate this difference quantitatively, we consider the linear response of an electron gas with an isotropic energy dispersion law $\varepsilon(\vec{p})$ to a perturbation due to an electromagnetic wave $\vec{E}(\vec{r}, t)$, $\vec{H}(\vec{r}, t) \sim \exp[i\omega t - \vec{k}\cdot\vec{r})]$. The kinetic equation of the problem is

$$i(\omega - \mathbf{kv})\phi + \left(e\mathbf{F} + \frac{e}{c}[\mathbf{v} \times \mathbf{H}_0]\right)\frac{\partial\phi}{\partial\mathbf{p}} + \left(\frac{\partial\phi}{\partial t}\right)_c = -\left(e\mathbf{E} + \frac{e}{c}[\mathbf{v} \times \mathbf{H}]\right)\frac{\partial f}{\partial\mathbf{p}},$$

where $\phi(p, \vec{k}, \omega)$ is the sought increment to the stationary distribution function $f(\vec{p})$; $\vec{F}$ and $\vec{H}_0$ are constant fields; $(\partial\phi/\partial t)_c$ is the collision integral, $\vec{v}(\vec{p}) = (\partial\varepsilon(\vec{p})/\partial\vec{p}) \equiv \vec{p}/m(\varepsilon)$ is the velocity of an electron with momentum $\vec{p}$.

The solution of Eq. (1) in a diffusion approximation with respect to the field F can be obtained by the method [3], assuming $|\vec{k}\cdot\vec{v}|/\omega \ll 1$. We write out the real part of the tensor $\varepsilon_{\mu\nu}(\omega, \vec{k})$, assuming also that the frequency ω greatly exceeds the mean reciprocal relaxation time $\tau^{-1}(\varepsilon)$ and the mean cyclotron frequency $\omega_c(\varepsilon) = -eH_0/cm(\varepsilon)$

$$\mathrm{Re}\,\epsilon_{\mu\nu}(\omega, \mathbf{k}) = \mathrm{Re}\,\epsilon_{\mu\nu}(\omega) + \hat{a}(\mathbf{k}\mathbf{G}(\epsilon))\delta_{\mu\nu} + \hat{\beta}(k_\mu G_\nu(\epsilon) + k_\nu G_\mu(\epsilon)). \tag{2}$$

The tensor $\mathrm{Re}\,\varepsilon_{\mu\nu}(\omega)$ is given in explicit form in [3]. If $\vec{H}_0 = \{0, 0, H_0\}$, then the vector $\vec{G}(\omega)$ in (2) is defined by the components

$$\mathbf{G}(\epsilon) = \left\{ \frac{F_x - \omega_c\tau F_y}{1 + (\omega_c\tau)^2}\,;\quad \frac{\omega_c\tau F_x + F_y}{1 + (\omega_c\tau)^2}\,;\quad F_z \right\}, \tag{3}$$

and the action of the operators $\hat{\alpha}$ and $\hat{\beta}$ on the arbitrary function $\psi(\varepsilon)$ is given by the following integrals:

$$\hat{a}\,\psi(\epsilon) = \frac{32\pi^2 e^3}{15\omega^3}\int_0^\infty d\epsilon\,\psi(\epsilon)\frac{df_0}{d\epsilon}\,\frac{p^5\tau(\epsilon)}{m^2(\epsilon)}\,\frac{dm^{-1}(\epsilon)}{d\epsilon}, \tag{4}$$

$$\hat{\beta}\psi(\epsilon) = \frac{16\pi^2 e^3}{15\omega^3}\int_0^\infty d\epsilon\,\psi(\epsilon)\frac{df_0}{d\epsilon}\,\frac{\tau(\epsilon)}{m(\epsilon)}\,\frac{d}{d\epsilon}\left(\frac{p^5}{m^2(\epsilon)}\right) \tag{5}$$

$f_0(\varepsilon)$ is the symmetrical part of the distribution function $f(\vec{p})$.

Substitution of (2) in the dispersion equation shows that in the general case the transverse and longitudinal waves do not separate. In a number of particular cases, however, to which we confine ourselves for simplicity, propagation of transverse waves is possible. The condition $(\vec{k}\cdot\vec{E}) = 0$ is satisfied, for example, for waves with

I) $\mathbf{E} \parallel \mathbf{H}_0,\ \mathbf{k} \perp \mathbf{H}_0,\ \mathbf{k} \parallel \mathbf{j}_0,\ \mathbf{F} \perp \mathbf{H}_0$; [1)]

II) $\mathbf{E} \parallel \mathbf{H}_0,\ \mathbf{k} \perp \mathbf{H}_0,\ \mathbf{k} \perp \mathbf{j}_0,\ \mathbf{F} \perp \mathbf{H}_0$;

III) at the Faraday configuration ($\vec{k} \parallel \vec{H}_0$, $\vec{E} \perp \vec{H}_0$), if $\vec{F} \parallel \vec{H}_0$.

In each of the cases (I) and (II), the dispersion equation has two solutions corresponding to linearly polarized waves propagating in opposite directions (arbitrarily "+" and "-" waves). The refractive indices for these waves differ by an amount

$$\Delta n_{\mathrm{I,II}} \equiv n^{(+)}_{\mathrm{I,II}} - n^{(-)}_{\mathrm{I,II}} = \frac{\omega}{c}\,\hat{a}\,G_{x,y}(\epsilon). \tag{6}$$

In case (III) there are four solutions ($n^{(1\pm)}$ and $n^{(2\pm)}$) for waves with opposite circular polarizations (1 and 2) and for opposite propagation directions ("+" and "-"). The difference between the refractive indices for the waves "1+" and "2-" is

$$\Delta n_{\mathrm{III}} = n^{(1+)} - n^{(2-)} = 2\frac{c}{\omega}\theta_0 + \frac{\omega}{c}\hat{a}F_z, \tag{7}$$

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1) $\vec{j}_0$ is the vector of the stationary current density.

where θ_0 is the specific (per unit length) Faraday rotation angle [3].

Thus, we see from (6), (7), and (4) that, unlike in [1], the dependence of the refractive index on the propagation direction of an electromagnetic wave relative to the current $\vec{j}_0$ is determined by the nonparabolicity of the energy band. If there is no magnetic field, then $\Delta n_{II} = 0$ and $\Delta n_I = \Delta n_{III} = (\omega/c)\hat{\alpha}F = \Delta n$. In the extremely degenerate case, Δn can be represented in the form

$$\Delta n = j_0 \frac{8\pi e p^2(\epsilon)}{5\omega^2 c m^3(\epsilon)} \frac{dm(\epsilon)}{d\epsilon}\Bigg|_{\epsilon = \epsilon_F} . \qquad (8)$$

An estimate based on (8) gives for the experimental conditions of [1] (n-InAs, T = 80°K, carrier density 5.5×10^{17} cm^{-3} and $\lambda = 3.39\ \mu$) a value of Δn which is smaller by a factor 2 - 3 than the observed value. This is apparently due to the fact that interband effects play an important role in this experiment (the energy of the light quantum is close to the width of the forbidden band). On the other hand, in the region of longer wavelengths, where interband effects are insignificant, measurements of Δn make it possible, in principle, to determine a number of important plasma parameters in the conduction band (see formula (8)).

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STATIONARY TURBULENCE OF A PARAMETRICALLY UNSTABLE PLASMA

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The theoretically predicted [1] phenomenon wherein the energy of a powerful radiation flux is transferred to a plasma at anomalously high speed is due to the onset of parametric instability of the plasma. Of great importance to the full understanding of this phenomenon is the development of concepts concerning the turbulent state produced in the plasma during the evolution of the parametric instability. The results obtained to date in the theory of stationary turbulence are connected with allowance for the nonlinear interaction of the growing plasma perturbations, this interaction being due to stimulated scattering of the waves by the particles. It is precisely the scattering of ion-acoustic waves by ions which made it possible to obtain a stationary level of fluctuations of a turbulent non-isothermal plasma in the field of a strong pump wave [2] (see also [3]). A stationary turbulence level was subsequently obtained under somewhat different conditions [4, 5] via scattering of waves by particles.

The stimulated scattering of waves by particles, which were taken into account in [2 - 5], are by far not the only nonlinear processes of interaction between waves and plasma particles in the same approximation of the expansion in powers of the plasma-fluctuation intensity. We wish to call attention here to the role that the nonlinear frequency shift of the plasma oscillations plays in the establishment of the stationary turbulence level of a parametrically unstable plasma. The stabilizing effect of such a frequency shift is