where n_F is the concentration of the localization centers with energy close to ϵ_F. For an arbitrary electric conductivity $\sigma_{\parallel}(\omega)$, for which the polarization vector $\vec{e}$ is parallel to the field F, we obtain by integration, with exponential accuracy,

$$\sigma_{\parallel}(\omega) \approx \frac{\lambda}{2} x_o^3 e^{-2x_o}, \qquad \lambda = \frac{2\pi}{9} \frac{e^2}{m^2} n_F \frac{\hbar^2 a^3}{F\omega}. \tag{3}$$

Here x_0 is the only positive root of the equation

$$x^4 + ax^3 - d = 0, \quad a = \frac{\hbar\omega}{Fa_B}, \quad d = \frac{\bar{\epsilon}}{Fa_B}(a\bar{R})^3. \tag{4}$$

If $a^4 << d$, i.e., $\hbar\omega << [\bar{\epsilon}(FR)^3]^{1/4}$ and $x_0 \simeq d^{1/4} - a/4 >> 1$, we have

$$a_{\parallel}(\omega) \sim \exp\left\{-2\left[\frac{\bar{\epsilon}(a\bar{R})^3}{Fa_B}\right]^{1/4} + \frac{\hbar\omega}{2Fa_B}\right\}. \tag{5}$$

According to [2], the absorption coefficient in the absence of a constant field is $\alpha(\omega) \sim \exp(-\text{const}\cdot\omega^{-1/3})$. As seen from (5), the constant field alters this relation quite strongly.

For the transverse electric conductivity $\sigma_{\perp}(\omega)$ (the polarization vector $\vec{e}$ is perpendicular to the direction of the field $\vec{F}$) we obtain similarly

$$\sigma_{\perp}(\omega) \approx \frac{3}{4} \lambda x_o^2 e^{-2x_o} \tag{6}$$

with the same x_0 as in (3). The ratio of the longitudinal conductivity to the transverse one is

$$\sigma_{\parallel}/\sigma_{\perp} = (2/3)x_o >> 1. \tag{7}$$

In conclusion, I am grateful to V.L. Pokrovskii for a discussion of the work.

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QUASISTATIONARY NOISE IN THERMIONIC EMISSION OF SEMICONDUCTORS

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As is well known [1], the quasistationary thermoinic field at the surface of a heated body makes the main contribution to the density of the electromagnetic energy and to the Maxwell stresses, but does not take part in the production of the energy flux. In particular, the quasistationary thermionic field makes the main constribution to the adhesion forces between two closely located bodies [2].

If a uniform semiconductor with temperature T occupies the half-space $z < 0$ then the spectral density of the fluctuations of the normal component of the thermionic field over the surface is given by [1]

$$< E_z^2(\omega, z) > = \frac{kT}{\pi c}\left\{\frac{c\epsilon''}{2z^3\omega|\epsilon_1 + 1|^2} - \frac{1}{2z^2}\,\mathrm{Im}\,\frac{\epsilon_1^*\sqrt{1-\epsilon_1^*}}{(\epsilon_1^* + 1)^2}\right\}, \tag{1}$$

where $\varepsilon_1(\omega) = \varepsilon' - i\varepsilon''$ is the dielectric constant. From (1) we see that the spectral intensity of the field fluctuations increases noticeably at small z (within the limits of the macroscopic analysis $z \gtrsim a$, where a is the correlation radius).

The quasistationary field can influence the electron emission, producing a noise in the current, obviously in those cases when this current is determined by the field at the cathode surface.

Let us examine in greater detail this influence in thermionic emission of semiconductors. Assume that an external constant field E_0 is normal to the surface of the semiconductor and causes electron emission in the region of the Schottky effect. This field is weak, so that the theory of electric images is valid and the normal Schottky law is satisfied, i.e.,

$$z_m = \frac{1}{2}\sqrt{\frac{e}{E_0}} \gtrsim r, \tag{2}$$

where r is the Debye radius [3]. We neglect the space charge of the electrons and assume therefore that there is no mutual correlation of the electrons in the emission.

Under these conditions, Schottky shot noise should be observed if the emission of each electron is statistically independent [4]. However, the emission current should contain also a noise due to the fact that the "height" of the Schottky barrier fluctuates under the influence of the quasistationary field $E_z(z, t)$. Let us calculate the current fluctuations in this case.

We assume that $E_z(z, t) << E_0$. We are interested in frequencies $\omega < v/z_m$, where z_m is the position of the maximum of the Schottky barrier and v is the thermal velocity of the electrons. Thus, during the time of electron emission the height of the barrier remains constant and the inhomogeneity of the barrier, due to the quasistationary field at these frequencies, is small, $\lambda \sim c/\omega >> z_m$. We can therefore assume that the shot fluctuations are independent of the barrier fluctuations. Then the spectral intensity of the current fluctuations over the positive frequencies in the band $\omega < v/z_m$ can be calculated in analogy with [5], and we obtain

$$< \Delta I^2(\omega) > = 2Ie + \frac{I^2 e^3}{2}\,\frac{< E_z^2(\omega, z_m) >}{E_0 k^2 T^2}, \tag{3}$$

where I is the Schottky current. This formula is valid when $z_m \gtrsim a \sim r$, i.e., when the external field at the cathode is

$$E_0 \lesssim \frac{e}{4r^2}.$$

Expression (3) is a superposition of the shot and quasistationary fluctuations of the current. It is convenient here to investigate the ratio of the intensity of the quasistationary noise to the shot noise. To this end we put

$$\epsilon_1 = \epsilon'(\omega) - i\frac{4\pi\sigma(\omega)}{\omega} \text{ and } |\epsilon_1| >> 1, \tag{4}$$

where σ is the conductivity, and obtain

$$\gamma(\omega) = \frac{I\sqrt{eE_0}}{\pi k T\epsilon'}\frac{\tau}{1+\omega^2\tau^2}, \tag{5}$$

where $\tau = \epsilon'/4\pi\sigma$.

We see that the quasistationary noise in semiconductors has a characteristic frequency dependence. Moreover, if $\epsilon'(\omega)$ and $\sigma(\omega)$ depend little on the frequency, then this dependence is typical of noise in semiconductors; for example, it is characteristic of generation-recombination noise [4]. In semiconductors with not too high a conductivity, one can observe here a quasistationary noise which exceeds greatly the shot noise in intensity, at relatively small emission currents. Thus, at $\sigma \sim 10^{11}$ sec^{-1} ($\rho \sim 10$ ohm-cm), $E_0 = 10^2$ V/cm, $T \sim 10^3$ °K, I = 1 mA, $\gamma(0) \sim 70$. Whereas in semiconductors the main contribution to the frequency dependence of $\gamma(\omega)$ is made by the first term of (1), in metals the main contribution is made by the second term, while the first gives rise to an increment that is independent of the frequency [5].

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PARAMETRIC BACK SCATTERING OF A LINEAR ELECTROMAGNETIC WAVE IN A PLASMA

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Possible mechanisms of energy absorption from a powerful laser beam by the plasma corona of a contracting deuterium drop [4, 5] are connected with parametric instabilities [1 - 3]. Two parametric processes lead to the conversion of the energy of the incident electromagnetic wave into the energy of plasma oscillations, namely the "decay instabilities" - the pump wave (photon) "plasmon" + "phonon" (ion sound) and the pump wave "plasmon" + "plasmon." The first of these processes can occur in a certain layer of the corona, where the frequency of the incident electromagnetic wave ω_0 is approximately equal to the electron plasma frequency

$$\omega_0 \cong \omega_p = \sqrt{\frac{4\pi n_c e^2}{m}}.$$