d) in the analytic terms of the decay amplitude it is possible to confine oneself to the terms linear in the energies and in the masses;

e) non-analytic (singular) terms arise only as a result of the diagrams shown in Figs. 1 and 2; this is valid of the scattering lengths are not large at low energies $a_T \lesssim 1/m_\pi$;

f) the pion mass differences can be regarded in the non-analytic terms.

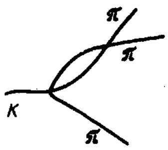

Fig. 1

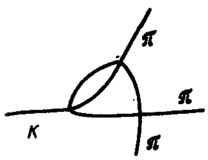

Fig. 2.

We note that the contribution of the diagrams (Fig.2) can be neglected for the following reasons: the contribution of these diagrams to the part of the amplitude corresponding to the symmetrical Young pattern is the same for all decays and therefore cancels out, and the interference of the symmetrical and asymmetrical Young patterns vanishes upon integration over the circle, while the square of the asymmetrical Young pattern is of the order of $m_\pi^2 E^2 a^4 < E^2/m_\pi^2$ and should be discarded.

It is seen from the table for $\Phi(R)$ that for the circle $R = 1/3$ the right-hand side of (1) does not exceed 1%, making it possible to check the assumption at the level of about one per cent.

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PHOTOCONDUCTIVITY OF SEMICONDUCTORS AT LARGE LIGHT INTENSITY

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1. It is known that the absorption coefficient of semiconductors decreases with increasing light intensity. According to Krokhin [1], the saturation effect consists of an equalization of the electron populations in the valence and conduction bands, so that the difference of the Fermi quasilevels of the electrons and holes becomes equal to the frequency of the light

$$\mu_n - \mu_p \approx \omega, \quad \hbar = 1. \tag{1}$$

Obviously, the concentration of the photoelectrons (holes) tends to the constant limit n.

2. We show in this paper that the photocurrent produced under the influence of an external electric field decreases with increasing light intensity at high intensities. The physical reason for this effect is that the anisotropy of the electron momentum distribution, appearing under the influence of the field E, leads to anisotropy of the absorption. The electrons are now produced predominantly with a momentum opposite that obtained from the field, since the absorption coefficient for such momenta becomes larger.

Thus, an electron redistribution takes place and decreases the total momentum.

3. We consider, for simplicity, the case of quadratic dispersion of electrons and holes with equal effective masses. The kinetic equations for the electron distributions functions $f_2(\vec{p})$ and $f_1(\vec{p})$ in the conduction and valence bands, respectively, are

$$-eE\frac{\partial f_2(p)}{\partial p} = \left(\frac{\partial f_2}{\partial t}\right)_{st} - \frac{f_2(p)}{\tau_R} + \frac{f_1(p) - f_2(p)}{\tau_I},$$
$$-eE\frac{\partial f_1(p)}{\partial p} = \left(\frac{\partial f_1}{\partial t}\right)_{st} - \frac{f_1(p)}{\tau_R} + \frac{f_2(p) - f_1(p)}{\tau_I}, \tag{2}$$

where $(\partial f_i/\partial t)_{st}$ are the integrals of electron collisions with the lattice, impurities, and electrons, $1/\tau_R$ is the recombination probability, $1/\tau_I = I(\omega)|M|^2$ is the probability of electron transition from the valence band to the conduction band under the influence of light of intensity $I(\omega) = Ig(\omega)$, $g(\omega) \sim 1/\delta\omega$ in the frequency band $\delta\omega$, with maximum at $(\omega - \Delta)/2$, and Δ is the width of the forbidden band. We note that the absorption coefficient is connected with τ_I by the relation

$$k(\omega) = \rho(\mu)/I(\omega)\tau_I,$$

where $\rho(\mu)$ is the density of states.

Only direct transitions between bands are taken into account in (2), and the photon momentum is neglected compared with the electron momentum. If the electric field is equal to zero, the functions f_1 and f_2 depend only on the energy and are Fermi functions [1], e. g.,

$$f_2(\epsilon) = \left[\exp\left(\frac{\epsilon - \mu_n}{kT}\right) + 1\right]^{-1}, \quad \mu_n \simeq \frac{\omega - \Delta}{2}. \tag{3}$$

In the presence of a weak field $\vec{E}$, we seek the solution in the form

$$f_i(p) = f_i(\epsilon) + f_i^{(1)}(p), \quad f_i^{(1)} << f_i(\epsilon). \tag{4}$$

Substituting (4) in (2), we get in the τ-approximation [2]

$$f_2^{(1)}(p) - f_1^{(1)}(p) = eEv_2\frac{\partial f_2(\epsilon)}{\partial \epsilon}\tau^*(\epsilon), \quad \tau^* = \frac{\tau\tau_I}{2\tau + \tau_I}, \tag{5}$$

where τ is the total momentum relaxation time, and account is taken of the fact that

$$v_2 = \partial\epsilon_2/\partial p - v_1 = \partial\epsilon_1/\partial p. \tag{6}$$

The expression for the current, with allowance for (4) - (6), becomes

$$j = -e\frac{2}{(2\pi)^3}\int(v_1 f_1(p) + v_2 f_2(p))d^3p =$$
$$= -e^2\frac{2}{(2\pi)^3}\int v_2(v_2 E)\frac{\partial f_2(\epsilon)}{\partial \epsilon}\tau^*(\epsilon)d^3p. \tag{7}$$

Assuming that $kT < \delta\omega << \mu$, we get after some calculations

$$j = E(e^2 n\tau^*(\mu)/m). \tag{8}$$

We see therefore that τ_1 plays the role of an addition momentum relaxation time, and as $I \to \infty$ we get $j \sim 1/I \to 0$.

4. To observe the effect experimentally, it is necessary to use thin samples with $k(\omega)\ell < 1$. The photocurrent begins to decrease at intensities

$$I > \frac{n}{k(\omega)\tau} \frac{\delta\omega}{\mu} \approx 10^{24} \text{ quanta/cm}^2\text{sec},$$

where $k(\omega) = 10^3$, $\tau = 10^{-11}$ sec, $n = 10^{18}$, and $\delta\omega/\mu \approx 10^{-2}$. Intensities of this order of magnitude are perfectly feasible.

In conclusion we note that a study of the photoconductivity at large light intensities can yield information on the saturation effect, and also on the values of the times τ_R, τ, and τ_1.

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