$n_3 = 10^{15}$ cm^{-3}, I = 10^{16} photon/sec-cm^2 leads to a value $\sim 10^2$ sec.

IV. In principle, anomalous photoconductivity can be possessed not only by a three-layer system (e.g., n-p-n or p-n-p), but also by a system made up of two layers in contact, but of different materials, i.e., a heterojunction or even a junction of a semiconductor with a metal (Fig. 3). The role of the third (barrier) layer is played in such a "twolayer" system by the space-charge region at the boundary. It is possible that the material used in [1, 2] is a system of this kind (e.g., Se-HgSe heterojunction or mercury-semiconductor junction).

It follows from the foregoing, however, that the phenomenon of anomalous photoconductivity, first observed and known so far only in the Se + Hg system, can be deliberately produced also with other semiconducting materials.

We note in conclusion that, as shown by analysis, in the case of the usual n-p homojunction the absorption of light by free carriers in one of the regions and their transfer to the other region where they recombine with the majority carriers (a process inverse to that considered in detail in [3]) also leads to unique phenomena that are close to "anomalous" photoconductivity (inverse proportionality of to the light intensity and independence of the stationary photoconductivity of the intensity). In this case, however, the conductivity changes only on going to illumination with light of shorter wavelength, whereas the anomalous photoconductivity does not depend on the "direction" of the change of the wavelength.

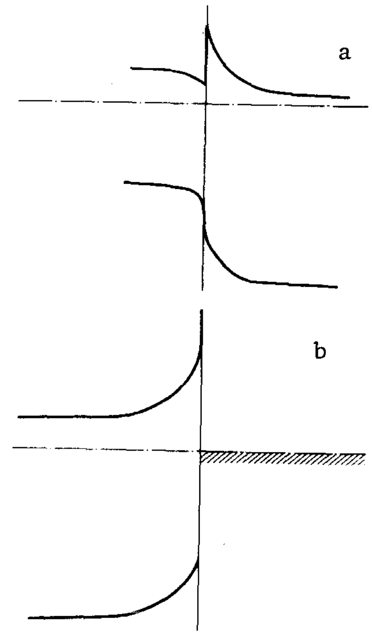

Fig. 3. Two-layer model of "anomalous" photoconductor (level scheme): a - semiconductor-semiconductor system, b - semiconductor-metal system.

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1) We assume for concreteness that the barrier region II is of the p-type, although it can in principle be also a high-resistance n- or i-region.

2) In analogy with the situation in field-effect transistors.

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EFFECTIVE INTERACTION OF ELECTRONS WITH SOUND ON CYLINDRICAL AND FLAT SECTIONS OF FERMI SURFACE

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Submitted 2 August 1973
ZhETF Pis. Red. 18, No. 6 381 - 384 (20 September 1973)

It is shown that absorption and dispersion of the spped of sound has a sharp maximum when the wave vector of the sound is inclined an angle $\sim s/v$ from the axis of the cylinder or from the plane. This phenomenon is analogous to the tilt effect.

In view of the large velocity difference, $s \ll v$, the electrons actively interacting with sound are those from a narrow strip on the Fermi surface, for which $\vec{k}\cdot\vec{v} = \omega$ (or $\vec{k}\cdot\vec{v} \simeq 0$ if $\omega \ll \nu$), and which move in phase with the wave. The fraction of these electrons is small, so that the relative absorption at $k\ell \gg 1$ reaches only a value $\sim s/v$, and the sound velocity dispersion is $\sim (s/v)^2$ [1].

On cylindrical and flat sections of the Fermi surface, however, if the sound wave vector $\vec{k}$ is directed along the cylinder axis or in the plane, all the electrons of these sections

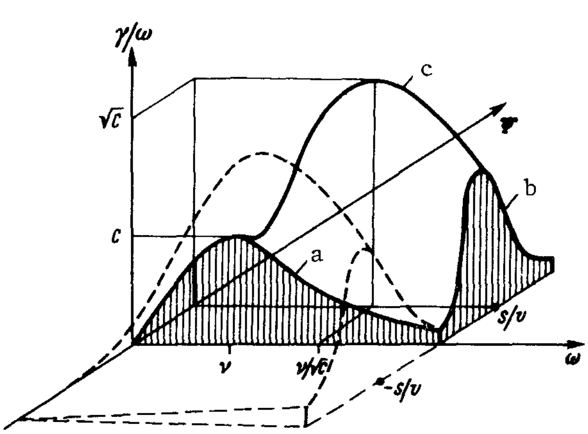

Fig. 1

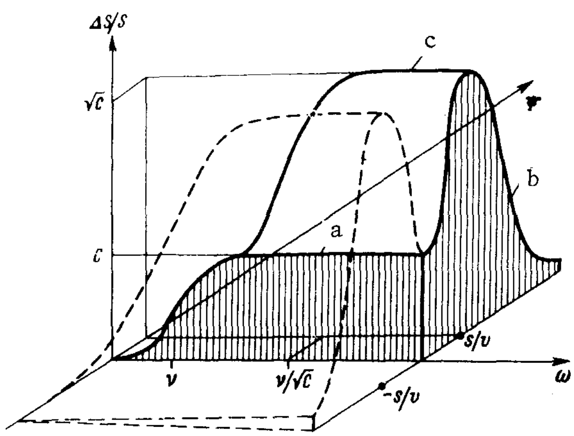

Fig. 2

Fig. 1. Sound absorption vs the inclination angle and frequency for a cylindrical Fermi surface.

Fig. 2. Dispersion of sound velocity vs inclination angle and frequency for a cylindrical Fermi surface.

interact effectively with the sound and move in the plane of the sound front ($\vec{k}\cdot\vec{v} = 0$). In real metals one frequently encounters flattened sections, and also sections close to cylindrical, such as corrugated cylinders or elongated ellipsoids. For example, in metals and alloys of the tungsten type there are octahedra with flat faces perpendicular to threefold axes (the causes of the octahedra can be understood from the strong-coupling model) [2]. A near-cylindrical surface is encountered for graphite. Under conditions of effective interaction, the absorption $\gamma/\omega \sim C[\omega\nu/(\omega^2 + \nu^2)]$ has at $\omega = \nu$ a maximum on the order of the electron-phonon coupling constant C, and the dispersion $\Delta s/s \sim C[\omega^2/(\omega^2 + \nu^2)]$ saturates at $\omega \gg \nu$ at the level of C [3] (Figs. 1a and 2a). A similar renormalization of sound velocity occurs in a strong magnetic field at $\vec{k} \perp \vec{H}$ [4].

At high frequencies $\omega \gg \nu$ the absorption decreases, since $\vec{k}\cdot\vec{v} = 0$ is no longer the resonance condition. It is possible, however, to satisfy the condition $\vec{k}\cdot\vec{v} = \omega$ by tilting $\vec{k}$ away from the cylinder axis (from the plane) through an angle $\psi_{max} \sim \pm s/v$. Both the absorption[1)] and the dispersion have a sharp maximum at $\psi = \psi_{max}$ in their angular dependence; this maximum is symmetrical with respect to the value at $\psi = 0$. This phenomenon is analogous to the tilt effect in a magnetic field [6].

We confine ourselves to a purely deformation-type interaction. Then, introducing the relaxation time ν and omitting the inessential subscripts of the deformation potential $\Lambda_{ik}(\vec{p})$, we obtain as a first approximation

$$\frac{\Delta\omega}{\omega_0} = -\frac{i\omega_0}{\rho s^2}\int\frac{ds}{v}\,\frac{\Lambda^2}{i(kv_K\sin\psi - \omega_0) + \nu}. \tag{1}$$

For a plane we have v = const, so that it follows directly from (1) that

$$\frac{\Delta\omega}{\omega_0} = \frac{C_{pl}}{i(kv_K\sin\psi - \omega_0) + \nu}, \tag{2}$$

where $C_{pl} = \int\Lambda^2(ds/v)/\rho s^2$. The absorption $\gamma/\omega \sim -\mathrm{Im}(\Delta\omega/\omega_0)$ and the dispersion $\Delta s/s \sim \mathrm{Re}(\Delta\omega/\omega_0)$ have maxima at $\sin\psi_{max} \sim \pm s/v$ (Figs. 1b, 2b).

Integration over the cylindrical surface yields

$$\frac{\Delta\omega}{\omega_0} = \frac{C_c}{\sqrt{(\nu - i\omega_0)^2 + k^2v_K^2\sin^2\psi}}, \tag{3}$$

where $C_s = (1/\rho s^2) \int m^* \Lambda^2(\vec{p}) dp_z$ (the integration is from $-p_0$ to p_0). At low frequencies $\omega << \nu$ both γ/ω and $\Delta s/s \sim (\omega/\nu)[1 + (kv_\kappa\psi/\nu)^2]^{-1/2}$ are maximal at $\psi = 0$, and at high frequencies they are maximal at $\sin\psi_{max} > s/v_{min}$. As seen from (3), the effect has a threshold and differs thus from the analogous effect on flat sections, when ψ_{max} is reached all the electrons of the plane resonate.

Maxima of γ/ω and $\Delta s/s$ are observed also when $\vec{k}$ is tilted relative to the axis of an ellipsoid with semiaxes $\sqrt{2m\varepsilon}$ and $\sqrt{2M\varepsilon}$ (M >> m) by an angle

$$\psi_{max} \approx \begin{cases} \sqrt{m/M}, & v_{min}/s >> 1 \\ s/v_{min}, & v_{min}/s << 1 \end{cases} .$$

The maxima of the absorption and dispersion, as follows from (1), increase with frequency, like ω_0/ν for a plane and like $\sqrt{\omega_0/\nu}$ for a cylinder. To determine the values to which the growth continues, we examine the dispersion equation in greater detail. We discuss a nondegenerate case - longitudinal sound:

$$\left(\frac{\omega}{\omega_0}\right)^2 \equiv x^2 = \eta - \frac{ix}{\rho s^2} \int \frac{ds}{v} \frac{\Lambda^2}{i\left(\frac{v_\kappa}{s}\sin\psi - x\right) + \frac{\nu}{\omega_0}}, \qquad (4)$$

where $\eta = \lambda_{\kappa\kappa\kappa\kappa}/\rho s^2$ and $\omega_0 = ks$.

In the case of a plane, (4) reduces to

$$x^2 - \eta = \frac{C_{pl}}{\left(x - \frac{v_\kappa}{s}\sin\psi\right) + i\frac{\nu}{\omega_0}} . \qquad (5)$$

Near resonance we have $(v_\kappa/s)\sin\psi \simeq \eta$, and the solution (5) takes the form

$$x_{1,2} = \eta + \frac{v_\kappa}{s}\sin\psi - i\frac{\nu}{\omega_0} \pm \sqrt{\left(\eta - \frac{v_\kappa}{s}\sin\psi + i\frac{\nu}{\omega_0}\right)^2 + 2C_{pl}} . \qquad (6)$$

It is seen from (6) that the growth at the maximum continues up to frequencies $\omega_0 \sim \nu/C_{pl}$, after which saturation of the absorption and a decrease of the absorption set in. The largest value at the maximum is $\sim \sqrt{C_{pl}}$ (Figs. 1c and 2c).

For a cylinder it follows from (4) that

$$x^2 - \eta = -\frac{C_c x}{\sqrt{\left(x - i\frac{\nu}{\omega_0}\right)^2 - \frac{v_\kappa^2}{s^2}\sin^2\psi}} . \qquad (7)$$

In the resonance approximation, Eq. (7) reduces to a cubic equation, solution of which shows that the largest value $\sim C_c^{2/3}$ is reached at $\omega \sim \nu/C_c$, after which the dispersion saturates and the absorption decreases. We note that C < 1. (At C $\sim$ 1 the maxima of the angular and freqeuency dependences are approximately the same.)

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1) In [5], the authors pointed out the presence of an absorption maximum when $\vec{k}$ is inclined to the axis of the cylindrical surface.

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MULTIPLICITY IN DIFFRACTION DISSOCIATION PROCESSES AT HIGH ENERGIES

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Submitted 15 August 1973
ZhETF Pis. Red. 18, No. 6, 385 - 387 (20 September 1973)

We study the contribution of the three-pomeron mechanism to the average multiplicity of hadrons produced at high energies. We show that this mechanism leads to a doubly-logarithmic contribution $\langle n\rangle_D \simeq 10^{-2} a \ln^2 s/s_0$ in addition to the contribution $\langle n\rangle \simeq a \ln s$ of the multiperipheral mechanism. The connection with other fragmentation models is discussed.

The pp → pK inclusive experiments at ISR [1] have revealed the presence of forward-scattering peaks corresponding to diffraction dissociation (DD) of one of the colliding particles on a beam with large invariant mass M.

In the three-reggeon (TR) scheme these peaks are described by the contribution of the three-pomeron (TP) vertex, and neither a vanishing nor nonvanishing TP vertex at zero momentum transfer is as yet in disagreement with experiment [2]. It is therefore of interest to examine the consequences of these alternatives for the average multiplicity of the particles (pions) produced in accomanïment with the observed proton p in the considered kinematic region.

If we interpret various fragmentation models in TR language, then the problem reduces to determining the dependence of the cross section $\sigma_{Pp}(t, M^2)$ for the scattering of a pomeron P by a proton, with allowance for the relative suppression of productions of beams with large M. Thus, the description of the inelastic diffraction peak by means of the non-scaling term PPR (R - secondary trajectory with $\alpha_R(0) = 0.5$) leads to a DD contribution $\sigma_D \sim$ const to the total cross section and to an average multiplicity $\langle n\rangle_D \sim \ln s$ [3]. Moreover in [4], where a zero TP vertex PPP was assumed and only trajectories with $\alpha_R(0) < 0.5$ were believed to contribute to PPR, it was concluded that the double DD cross section decreases, $\sigma_{DD} \sim \ln^{-1} s$, and the multiplicity $\langle n\rangle_{DD}$ is constant. In both cited papers it was assumed for saturation of $\langle n\rangle$ that the multiplicity in the beams is maximal, $\langle n(M^2)\rangle \sim M$. This question was considered in [5] for a TP vertex that does not vanish at t = 0, under the self-consistent multiperipheral assumption $\langle n(M^2)\rangle \sim \ln M$, and it was concluded that $\langle n\rangle_D$ increases logarithmically, like ln s. There however, the average multiplicity $\langle n\rangle_D'$ in DD processes was not quite correctly identified with the DD contribution $\langle n\rangle_D$ to the total average multiplicity.

It is shown in the present paper that in the pre-asymptotic region (the ISR energy region) there is besides the growth $\sigma_D \sim \ln s$ also a doubly-logarithmic growth $\langle n\rangle_D \sim \ln s$. The TP contribution to the pp → pX spectrum is ($\alpha_P(0) = 1$)

$$\frac{d\sigma}{dt\,dM^2} = \frac{G(t)}{M^2}\left(\frac{s}{M^2}\right)^{2\alpha_P' t}, \tag{1}$$

where for the effective TP vertex G(t) is chosen in the form $G(t) = -\tilde{G}t\exp(\tilde{R}^2 t)$ and $G(t) = G(0)\exp(R^2 t)$ for the case when it vanishes and does not vanishes, respectively, at zero t. It is assumed that the Pp → X process is similar to the usual hadron process, i.e., that its average multiplicity, normalized to σ_{Pp}, takes at large M^2 the form

$$\langle n(t, M^2)\rangle \simeq a \ln M^2, \tag{2}$$