V_y depends on the sign of N. When N = 0 we always have $V_y \sim V_x$. This explains why the motion of the normal regions was observed in [8] in indium ($N \neq 0$) but not in tin (N = 0).

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1) For long waves $\zeta(x = a_n/2) \approx \zeta(x = -a_n/2)$.

PRESSURE DEPENDENCE OF ELECTRON EFFECTIVE MASS IN INDIUM ANTIMONIDE

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Submitted April 26, 1966
ZhETF Pis'ma 4, No. 1, 11-15, 1 July 1966

By changing the width of the forbidden band in a crystal, hydrostatic compression affects the value of the effective mass of the carriers. If the band is non-parabolic, as is the conduction band in InSb, the effective mass varies with the filling of the band, and the pressure can have diverse influences on the effective mass.

According to Kane's theory, the electron effective mass m_n at the bottom of the conduction band is given by the formula

$$\frac{m_0}{m_n} = 1 + \frac{\epsilon_p}{3}\left(\frac{2}{\epsilon_g} + \frac{1}{\epsilon_g + \Delta}\right) , \qquad (1)$$

where $\epsilon_g = \epsilon(\Gamma_{1c}) - \epsilon(\Gamma_{15v})$ is the width of the forbidden band at the Γ-point, Δ the energy of spin-orbit splitting of the valence band at the Γ-point, $\epsilon_p = (2m_0/\hbar^2)\mathcal{P}^2$, $\mathcal{P}$ the matrix element of the interaction between the states Γ_{1c} (conduction band) and Γ_{15v} (valence band), and m_0 the mass of the free electron.

Under the usual assumption that ϵ_p is not affected by small changes of the lattice period, the mass m_n is almost proportional, according to (1), to ϵ_g, which in InSb increases linearly with the pressure P with a coefficient $d\epsilon_g/dP = 1.6 \times 10^{-5}$ eV/atm.

We undertook an experimental study of the influence of hydrostatic pressure up to 16.5 katm on the effective mass m_n of the electrons in InSb at T = 96°K. To this end, we measured the thermal emf and the Hall effect in classically strong magnetic fields, when neither depends on the electron scattering. From the Hall effect we determine the electron density n, and then the effective mass is determined from the value of the thermal emf α_∞ in the saturation region.

The measurements were made at temperature gradients 3 - 6 deg/cm; the difference in the temperature drops on opposite faces of the sample did not exceed 2%. The termocouples were introduced in the high-pressure chamber without a break of the continuity. The pressures were produced at nitrogen temperatures by a method proposed by Itskevich [1]. The investigated samples measured 10 x 3 x 2 mm.

In the absence of degeneracy and at low value of the non-parabolicity parameter $\gamma = kT/\epsilon_g$, the thermal emf in the saturation region equals, accurate to terms $\sim\gamma^2$,

$$\alpha_\infty = \frac{k}{e}\left(\frac{5}{2} + \frac{15}{2}b\gamma - \frac{45}{4}a\gamma^2 - \mu_0^*\right), \qquad (2)$$

where $\mu_0^* = -\ln\,[2(2\pi m_n kT)^{3/2}/nh^3]$, and a and b are some simple functions of ϵ_g and Δ (see [2]).

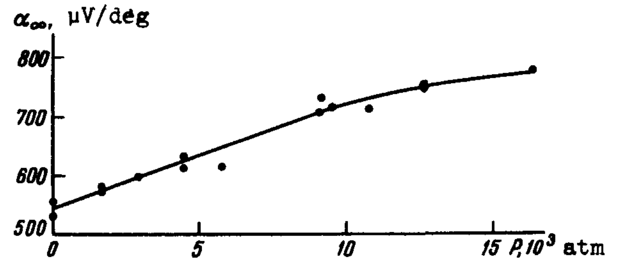

Fig.1

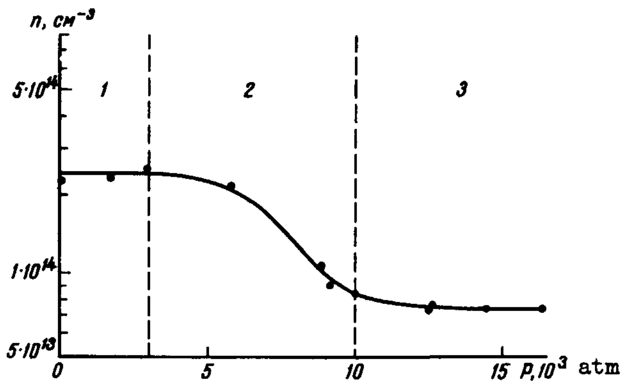

Fig. 2

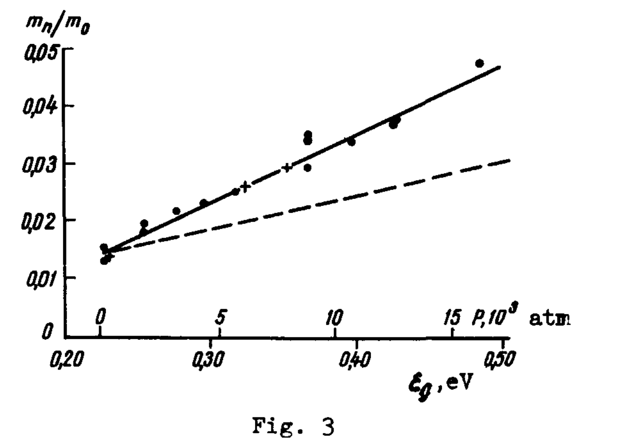

Fig. 3

Figure 1 shows a plot of α_∞ against P for two samples of n-InSb with $n \simeq 2.2 \times 10^{14}$ cm^{-3}, and Fig. 2 shows a plot of n vs. P. The pressure dependence of the effective mass m_n, calculated in accordance with (2), is shown in Fig. 3 for samples with $n \simeq 2.2 \times 10^{14}$ (•) and $n = 4.7 \times 10^{13}$ cm^{-3} (+). The same figure shows a theoretical plot of $m_n(P)$ calculated from (1) (dashed curve). In the calculations we assumed that $\epsilon_g(96°K) = 0.226$ eV, $\Delta = 0.9$ eV, and $\epsilon_p = 23$ eV. With increasing pressure, the disparity between the experimental and theoretical curves increases and appreciably exceeds the experimental errors. Two possible causes of this disparity have been considered: (i) change of matrix element $\mathcal{P}^2$ with pressure, (ii) change of perturbation of the mass m_n by the remote bands with changing pressure.

To reconcile the experimental and theoretical values of m_n it must be assumed that $\mathcal{P}^2$ increases by 20% at P = 5 katm and by 35% at 16.5 katm. If we recognize that at 16.5 katm the lattice constant increases by ≈1.5%, then admittedly such large changes in $\mathcal{P}^2$ are unlikely.

According to Cardona [3], allowance for the influence of the Γ_{15c} band leads to the following expression for the mass at the bottom of the Γ_{1c} band:

$$\frac{1}{m_n^*} \simeq \frac{1}{2m_n}\left(1 + \frac{\epsilon_{g1}}{\epsilon_{g2}}\right), \tag{3}$$

where m_n is determined by formula (1), $\epsilon_{g2} = \epsilon(\Gamma_{15c}) - \epsilon(\Gamma_{15v})$ and equals 3.4 eV for InSb, and $\epsilon_{g1} = \epsilon(\Gamma_{15}) - \epsilon(\Gamma_{2'})$ is the energy difference of the corresponding levels in gray tin and equals 2.8 eV. We have neglected in (3) terms that introduce an error not larger than 3%. Since there are no data on the variation of ϵ_{g1} and ϵ_{g2} with pressure in InSb and α-Sn, we can assume for estimating purposes that they are the same as in Si and GaP [4]. Assuming $d\epsilon_{g1}/dP = (+6 \times 10^{-6}$ to $+1 \times 10^{-5})$ eV/atm, we found no noticeable deviation from the theoretical curve of Fig. 3. To reconcile the calculated and experimental curves we must assume that the baric coefficients $d\epsilon_{gi}/dP$ greatly exceed the experimental values. The smallest values of $d\epsilon_{gi}/dP$ for which agreement can be obtained are:

$$d\epsilon_{g2}/dP \simeq 68 \times 10^{-5} \text{ eV/atm}, \quad d\epsilon_{g1}/dP \simeq -8 \times 10^{-5} \text{ eV/atm}$$

or

$$d\epsilon_{g2}/dP \simeq +3 \times 10^{-4} \text{ eV/atm}, \quad d\epsilon_{g1}/dP \simeq +1 \times 10^{-5} \text{ eV/atm}.$$

It is therefore unclear whether the influence of hydrostatic pressure on the electron effective mass can be explained within the framework of Kane's theory, since there are no experimental data on $d\epsilon_{gi}/dP$. The accepted semi-empirical approach to the analysis of the effect of pressure on the main characteristics of semiconductors is patently inadequate, and a quantitative theory that takes into account the dependence of the band structure on the lattice period.

In conclusion we note that the pressure dependence of the electron density (Fig. 2) likewise remains without adequate explanation.

We can propose the following: In region 1 the change of ϵ_g with pressure is slightly larger than 2kT, and the density n remains practically constant. With further increase in pressure (region 2) the lower edge of the conduction band moves more rapidly (i.e., with a larger baric coefficient) than the donor level [5] and the gap between them increases, while n decreases. At $P \approx 10^4$ atm the distance between the impurity level and the bottom of the conduction band reaches ~0.15 eV, i.e., the shallow level becomes essentially deep. In this connection it can be assumed, in analogy with Ge and Si [5], that the speed of the level increases appreciably with pressure: in region 3 it becomes equal to the speed of the bottom of the conduction band.

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INSTABILITY OF PLASMA ON TRAPPED PARTICLES

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Submitted 27 April 1966
ZhETF Pis'ma 4, No. 1, 15-19, 1 July 1966

As is well known, flute instability can set in in magnetic-mirror traps and in toroidal traps with closed force lines [1,2]. This instability is the result of the oppositely directed drifts of the electrons and ions in the inhomogeneous electric field (force-line curvature radius R) with velocities $v_j = cT_j/e_jHR$ (H = magnetic field, T_j = temperature, e_j = charge of particle of species j). This drift causes polarization of the plasma in perturbations of the "tongue" type on the plasma surface, and if the magnetic field decreases in the outward direction, then the electric field due to the polarization causes plasma to be ejected. Flute instability was investigated experimentally by Ioffe et al. [3,4].

Analysis in the hydrodynamic approximation [5,6] has shown that in toroidal systems with open force lines lying on toroidal surfaces the plasma should be stable if the pressure is low enough. In a rarefied plasma, however, when the hydrodynamic approximation does not hold, the deduction that flute instability is stabilized by the crossing of the force lines is no longer valid. In fact, this deduction is based on the notion that in a high-temperature plasma the charges due to the magnetic drift should cancel each other as a result of flow along the force lines. Such a cancellation effect does indeed occur, but it applies only to the so-called transit particles, which move freely along the force line. If the magnetic field varies along the force lines, then there is present, besides the transit particles, also a group of "trapped" particles, which oscillate between the magnetic mirrors, i.e., between regions with stronger magnetic fields. If the change of the magnetic field along the force lines is small, then the mirror ratio $P = H_{max}/H_{min}$ is close to unity, and the fraction of particles trapped between mirrors is small, $\epsilon \cong \sqrt{P - 1}$. Since the captured particles are trapped between the mirrors, they cannot move freely along the force lines, and consequently an instability of the flute type can develop on them, and can be naturally called trapped-particle instability.

It differs from flute instability in that the charges due to the trapped particles are cancelled out to a considerable degree by the transit particles. Owing to this effect, the growth increment becomes small, and the transit particles have time to acquire a Boltzmann distribution, i.e., the perturbation of their density is equal to $-e_j\varphi n/T_j$, where φ is the electric field potential and n the unperturbed density. For the perturbation n_t' of the trapped-particle density we have in the quasiclassical approximation, when the perturbation of the potential is in the form of a plane wave, and neglecting the scatter in the drift velocities, the following continuity equation: