Anomalous galvanomagnetic properties of In-doped $Pb_{1-x}Sn_xTe$ at low temperatures

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Quantum oscillations of the Hall coefficient have been observed at 77 K in the narrow-gap semiconductor $Pb_{1-x}Sn_xTe(In)$, and a sign-varying magnetoresistance has been observed at 4.2 K. These results are explained in terms of special properties of a surface layer of the bulk semiconductor.

Experiments have shown[1,2] that the surfaces of the narrow-gap semiconductors $Pb_{1-x}Sn_xTe$, PbS_xSe_{1-x}, $Pb_{1-x}Sn_xSe$, and $Cd_xHg_{1-x}Te$ have some special properties which distinguish them from the bulk material. In this letter we are reporting a study of the properties of the surface layer in $Pb_{1-x}Sn_xTe(In)$ at low temperature. Specifically, we measured the Hall effect and the magnetoresistance in $Pb_{1-x}Sn_xTe$ doped with In ($C_{In}=0.5$ at. % $x=0.22$, 0.25, and 0.3) at liquid-nitrogen and liquid-helium temperatures in darkness. The magnitude and direction of the magnetic field were varied. Experiments were carried out in magnetic fields from 40 to 9000 Oe. We used the standard arrangement for measuring the resistivity and the Hall effect. The direction of the magnetic field was perpendicular to the direction of the current through the sample in all cases. The measurements were carried out point by point at steps of 2 Oe for two directions of the field and of the current through the sample. The dimensions of the sample were $1\times1.5\times5$ mm.

The most obvious result of the measurements of the Hall effect at $T=77$ K is the observation of oscillations in the Hall coefficient which are periodic in $(1/H)^k$, where $k\sim1$, in weak magnetic fields (Fig. 1). Oscillations of the Hall coefficient were observed in all the samples studied. The amplitude of these oscillations was 1.8–6.5% of the nonoscillatory part of R_x. The error of the R_x measurements was 1.3%. The largest oscillation amplitude was found in the samples with the composition $x=0.3$. We believe that these oscillations are magnetophonon oscillations of the Hall coefficient, i.e., a scattering of electrons by longitudinal optical phonons.[3] Working from the period of the oscillations in the reciprocal field, and knowing the frequency of a longitudinal optical phonon,[4] ω_{LO}, we calculated the effective mass of the current carriers involved in the resonance. It turned out to be $m^*=10^{-4}m_0$, i.e., exceedingly small and in agreement with effective mass found at liquid-helium temperature in Ref. 5. Lashkarev *et al.*[5] observed oscillations in the Hall coefficient in $Pb_{0.82}Sn_{0.18}Te$ at 4.2 K in very weak magnetic fields (beginning at 10 Oe). They attributed the effect to the complex shape of the constant-energy surface of the bulk spectrum—a shape which leads to the appearance of groups of carriers in "pockets" with a small effective mass $m^*=10^{-4}m_0$. However, the current models of the band structure of IV–VI materials rule out any pockets on the Fermi surface. We would like to suggest that the reason

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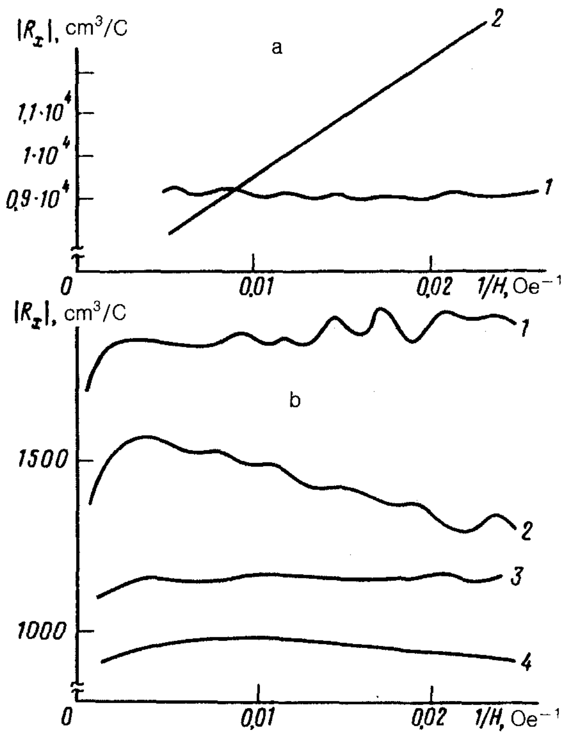

FIG. 1. Hall coefficient versus the reciprocal of the magnetic field (a) for various surface states of a sample and (b) for various angles between H and the surface of the sample (on which there were no contacts). a: *1*—Polished surface; *2*—ground surface. b: *1*—$\alpha=0$; *2*—$\alpha=30°$; *3*—$\alpha=45°$; *4*—$\alpha=60°$.

why such small effective masses are seen experimentally is a distortion of the band structure at the surface, where even a gap-free state may be reached.

When the surface quality is altered by coarse grinding (at $T=77$ K), the non-oscillatory part of the Hall coefficient decreases sharply with increasing magnetic field; the oscillations in the coefficient disappear completely. Grinding the surface means increasing the number of mechanical lattice defects of various sizes. These defects may limit the dimensions of the orbits of the free electrons, giving rise to an inhomogeneous broadening of the resonant levels. A further treatment of the sample in a polishing etchant restores the oscillations in the Hall coefficient (Fig. 1a).

Figure 1b shows the results of measurements of the Hall coefficient at $T=77$ K. These results were obtained by varying the angle between the magnetic field and the surface of the sample, on which there were no contacts. As $\alpha=0$ we took the position in which H was perpendicular to this surface of the sample. The angle was varied over the range $(0°\text{–}60°)\mp5°$. It can be seen from this figure that the period of the oscillations, $\delta(1/H)$, is proportional to $\cos\alpha$: $(\delta(1/H)_\alpha \approx \delta(1/H)_{\alpha=0}\cos\alpha)$. The oscillation amplitude decreases with increasing α, and the oscillations are no longer observed

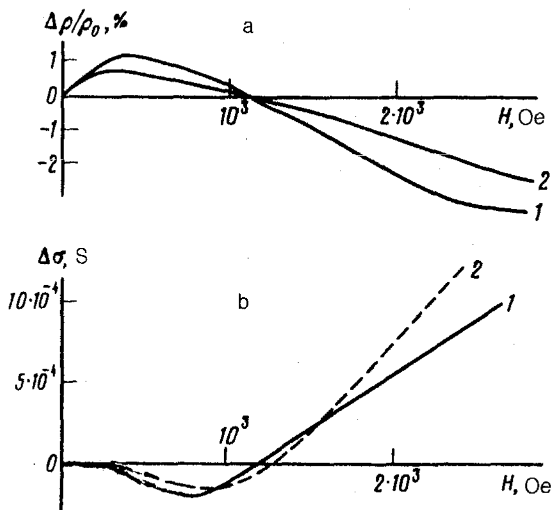

FIG. 2. Sign-varying magnetoresistance (a) and magnetoconductivity (b) versus the magnetic field. a: *1*—H is perpendicular to the surface of the sample; *2*—$\alpha=45°$. b: *1*—Experimental magnetoconductivity; *2*—theoretical.

at $\alpha=45°$. This behavior of $R_x=f(1/H)$ is evidence for the existence of a 2D electron gas in a surface layer of the sample.

Further evidence for a 2D nature of the electron conductivity comes from measurements of the magnetoresistance at $T=4.2$ K (Fig. 2). We see that the magnetoresistance varies in sign: It increases with the magnetic field, goes through a maximum, decreases, and goes negative. As the angle between the magnetic field and the normal to the main surface of the sample is varied, the curves of $\Delta\rho/\rho_0=f(H)$ shift up the magnetic-field scale, changing only slightly in shape (as in the Hall oscillations at $T=77$ K). The sign-varying magnetoresistance observed here is similar to that which was observed in Refs. 7 and 8 in studies of InGaAs/InP and *p*-InSb/*i*-GaAs heterostructures, in which the conductivity was mediated by a quasi-2D electron gas. For a quantitative explanation of the effects, we invoked the theory of quantum corrections to the conductivity, incorporating a weak localization of electrons with a spin–orbit coupling.

Adopting the hypothesis that a 2D electron gas exists in our material, and using the theory of weak localization with a spin–orbit coupling, we derived a theoretical magnetoconductivity from the equations given in Ref. 9:

$$\frac{\Delta\sigma(H)}{G_0}=\frac{3}{2}f(x_L^*)-\frac{1}{2}f(x_L). \tag{1}$$

Here $G_0=e^2/(2\pi^2/\hbar)$, $f(x)=\ln x+\psi(\frac{1}{2}+\frac{1}{x})$, ψ is the digamma function, and the notation is otherwise that of Ref. 9. From (1) we found a system of two equations for τ_φ and τ_{s0} (τ_φ is the relaxation time of the phase of the wave function, and τ_{s0} is the

spin–orbit interaction time) with a parameter D, which is the diffusion coefficient (or, respectively, the mobility μ). These equations are found by substituting the experimental values of H corresponding to $\Delta\sigma(H)=0$ ($H\neq0$) and $\partial\sigma/\partial H=0$ (the inflection point) into Eq. (1). In the interval $\mu=10^3$–10^7 (cm^2/(V·s), the system of equations had solutions for $\mu=10^4$ and 10^5 cm^2/(V·s). For $\mu=10^5$ cm^2(V·s) we find the values $\tau_\varphi=7.9\times10^{-13}$ s and $\tau_{s0}=4.0\times10^{-13}$ s. For $\mu=10^4$ cm^2/(V·s) we find $\tau_\varphi=4.6\times10^{-12}$ and $\tau_{s0}=2.5\times10^{-12}$ s. For the value $\mu=10^5$ cm^2/(V·s), we used the values found for τ_φ and τ_{s0} to plot a theoretical curve of $\Delta\sigma=f(H)$ (Fig. 2b). We see that the theoretical curve runs close to the experimental data in weak fields, while above $H\sim1500$ Oe the theory deviates from experiment. The probable reason for this deviation is a suppression of weak localization by a sufficiently strong magnetic field. Such a suppression of localization is part of the theory of Ref. 10.

The existence of carriers with a mass $m^*=10^{-4}m_0$, observed at liquid-nitrogen temperature, has yet to be explained. A necessary condition for observing quantum oscillations of any sort in kinetic coefficients in weak magnetic fields is that the direct gap $E_g=2\Delta$ be narrow. According to the results of the present study, a suitably narrow gap forms at the surface of the materials which we studied. Let us consider two microscopic reasons for the appearance of such a narrow-gap 2D spectrum. First, the width of the band gap decreases sharply or even changes sign (a band inversion occurs) near the surface, because of a local change in chemical composition or a deformation. As a result, coupled Tamm states arise at the surface. It was shown in Ref. 11 that a massless spectrum of the Weyl type generally arises in the case of a band inversion at a boundary. It was shown in Ref. 12 through a direct microscopic calculation that the appearance of Weyl states would be more likely in semiconductors with a 3D spectrum which is inverted in the interior ($x>0.33$), rather than in the semiconductors studied here.

We must therefore propose some other picture for the formation of a 2D spectrum with a narrow gap, in which we allow for the circumstance that the real surface of a semiconductor is charged and that the bands near the surface are greatly curved. A 2D electron gas then arises in this potential well. An important aspect of narrow-gap semiconductors is that since the gap is narrow these states can easily "seep" through the curved gap by a tunneling process, going into the interior; they accordingly become strongly hybridized with the bulk hole band. The scale depth of this seepage is $\kappa^{-1}=\hbar v/\Delta$, where $v\mu\sim10^8$ cm/s is a parameter of the bulk spectrum $E(k)$ of the semiconductor,

$$E(k)=\pm\sqrt{\Delta^2+\hbar^2v^2k^2}, \tag{2}$$

and $\hbar k$ is the quasimomentum. As a result, the gap "seen" by the 2D electrons ($2\Delta^*$) is much smaller than the bulk gap (2Δ). As the surface curvature increases, the level of the 2D electrons "runs into" the top of the bulk valence band, but because of the hybridization it is repelled from the bulk states, and it cannot go into the spectrum of the valence band. In this manner, an essentially gapless state is realized (Fig. 3).

The primary distinction between the state which arises according to this picture and a Weyl state is that the latter is not spin-degenerate. Accordingly, a spin flip due to a spin–orbit coupling is greatly hindered in this case, in contrast with the picture of

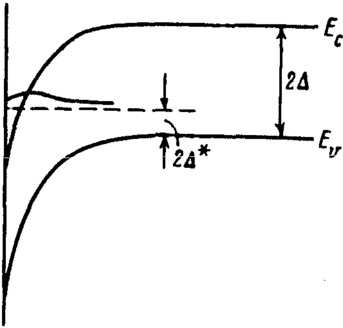

FIG. 3. Change in the energy gap at the surface.

band curvature in the Weyl case. There is thus the hope that the two pictures for the appearance of surface states described here can be distinguished by studying quantum corrections to kinetic effects.

We conclude by writing the spectrum of the 2D electrons in a magnetic field. In the customary case of a narrow gap, this spectrum is

$$E_n = \pm\sqrt{\Delta^{*2} + \left(\frac{\hbar v}{L_H}\right)^2 2n}, \tag{3}$$

where $L_H^{-2} = eH/\hbar c$, L_H is the magnetic length, H is the magnetic field, and $n = (0,1,...)$ is the index of the Landau level. When a tunneling-assisted seepage of 2D electron states occurs through the curved gap 2Δ near the surface, the spectrum of these states in the magnetic field retains the structure in (3), but we would have to insert the distance from the level of the surface electrons to the top of the valence band in the interior under the Δ^*, and we would have to replace the rate v in (3) by an effective $v^* = tv$, where $t < 1$ is a tunneling exponential with a decay scale κ.

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