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UNIAXIAL COMPRESSION EFFECTS ON CONDUCTION MECHANISMS IN P-TYPE GALLIUM ANTIMONIDE

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Résumé. — Nous avons mesuré entre 4,2 K et 300 K l'effet d'une compression uniaxe sur la résistivité du GaSb de type p. Nous avons déterminé les constantes de potentiel de déformation b et d du niveau Γ_8 ainsi que le rayon de Bohr effectif en faisant varier l'énergie de liaison de l'accepteur avec la contrainte. Nous avons comparé les résultats obtenus directement par spectroscopie de modulation avec ceux obtenus par cette méthode. Nous faisons également une étude des processus de conduction par impuretés dans ce type de composés III-V. Les résultats obtenus sont comparés aux résultats théoriques obtenus par différentes théories : théorie de Mikoshiba lorsque la conduction se fait par bande d'impureté, théorie de Abrahams et Miller et théorie de Shklovskii lorsque la conduction se fait par *Hopping*.

Abstract. — The effect of uniaxial compressional stress on the resistivity of p type GaSb has been measured between 4.2 K-300 K. We have determined the deformation potential constants b and d of the Γ_8 level and the effective Bohr radius by measuring the variation of the acceptor binding energy. We have compared data obtained directly by modulation spectroscopy with acceptor binding energy method under uniaxial stress. We have made a study of impurity conduction processes in this p type III-V compound and compared our experimental results with i) the theory of Mikoshiba for conduction by impurity band and ii) the models of Abrahams and Miller and Shklovskii for conduction by hopping.

1. Introduction. — Since the first investigation by Hall [1] on p type Ge, several investigations of the influence of uniaxial compression on the electrical properties of p type semiconductors have been carried out. These investigations were made on : germanium [2, 3, 4]; silicon [5, 6] and indium antimonide [7]. For p type GaSb, Tufte and Stelzer [8] and, second, Metzler and Becker [9] report results about high field coefficient behaviour in uniaxially stressed p-type samples. Besides, Averous *et al.* [10] studied the transport phenomena in the same conditions. Up to date, all these data on GaSb were obtained in the temperature range 50-300 K i.e., in a range where the conduction is mostly dependent of the carriers in the degenerate valence bands.

In this work, we give data in the temperature range 4.2-300 K, and we study three conduction processes; valence band conduction, impurity band conduction, and hopping conduction. For each process we determine the corresponding energy activation: ε_1 , ε_2 and ε_3 . Under uniaxial stress, the valence band

splits in two components : the heavy holes band and the light holes band. This splitting induces a change in the acceptor wave function which in normally built from an admixture of Bloch functions from the two valence bands. When these two bands split, it becomes more and more composed of Bloch functions from the upper valence band.

Studying the change in ε_1 , ε_2 and ε_3 under strain provides a very interesting and powerful Tool method for investigating some fundamental features of the band structure :

i) The deformation potential constants of the Γ_8 level : b and d in the Pikus-Bir notation [11].

ii) The effective Bohr radius, and its change under strain.

iii) It also permits us to test the different impurity conduction theories of Mycielski [12], Frood [13] and Mikoshiba [14].

iv) Lastly, we find from the variation of ε_3 (hopping activation energy) that the application of a pressure decreases the overlap of impurity waves function and the hopping process becomes more important.

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2. Experimental details. — Undoped single crystals of gallium antimonide were grown by the Czochralski method. (111) samples were cut perpendicular to the growth axis and (100) samples obliquely relative to the growth direction. X-rays orientation was better than one degree. Six Indium solder contacts were used in order to attach the four potential probes and two current leads. On good quality samples, without surface cracks, uniaxial compressional stresses as large as 10^9 dynes. cm^{-2} could be applied safely.

Our stress apparatus is shown in figure 1. It is made of two parts : the first one applies a static compressional stress, the second one superimposes a small alternating component.

FIG. 1. — The schematic figure of the apparatus used to apply the stress and to measure the change of valence bands, impurity and hopping conduction under stress.

The static part is made of a hydraulic generator which allows a soft transmission from an emitter piston to a receiver piston. The upper part of the receiver piston pushes against a frame and the lower part transmits the stress to a stirrup and, through a stainless steel rod, to the sample. The stainless steel rod is directed by axial ball bearings. The modulated part superimposes a small alternating stress on the static one. This is made through the stirrup : a motor applies a sine displacement at one end of a calibrated spring, the lower end of this spring applies consequently an alternating stress on the stirrup and hence on the sample. A compensation device ensures a frequency stability of the motor better than 3×10^{-4} , for any load.

The part of the stress apparatus holding the sample is shown in an enlarged scale in figure 1. The two

ends of the sample, which had a cross-section of about 1.5 mm^2 and a length of 8.1 mm , were cemented with epoxy resin into brass cups. The samples were placed between the stainless steel rod and the stress captor by two balls of glass which ensured both the alignment of the sample and the electrical insulation. The stress is measured by a stress captor which also measures the static and the alternating components. The alternating stress is about two hundred dyne- cm^{-2} . The d.c. current, in the sample is parallel to the axis of the sample and to the stress direction. The a.c. voltage induced in the sample by the alternating stress is measured by a lock-in amplifier with voltage obtained from the captor, used as reference. This modulation technique was used only for piezoresistance measurements.

The part of the stress apparatus shown in figure 1 under the plate is inside a regulated temperature cryostat.

3. Experimental results. — Figure 2 shows the change in resistivity versus $10^3/T$ for five samples

FIG. 2. — Resistivity as a function of $10^3/T$ for three samples in the (100) direction and two samples in the (111) direction. In the next figures, we keep the same symbols for each sample.

(three in the (100) direction and two in the (111) direction).

We distinguish three different domains of linearity for $\log \rho$ vs. $10^3/T$. We associate the two extreme regions with the valence band conduction and the hopping conduction processes respectively, and the intermediate region with an impurity band conduction. This assignment is supported by the Hall constant measurements shown in figure 3. In the regions corresponding to ε_1 and ε_2 , we could measure a Hall constant : R_H starts to increase with $10^3/T$ reaches a maximum and then decreases when $10^3/T$ increases. The maximum value is found when $p_h \mu_h = p_i \mu_i$ where p_h and p_i are the number of holes in the valence band and the impurity band respectively and μ_h and μ_i their mobility. R_H disappears at very low temperature, in the region ε_3 , so we conclude that it corresponds to a conduction by bound carriers with no free carriers contribution.

FIG. 3. — Hall constant as a function of $10^3/T$ for three samples, samples 1 and 2 : (100) direction; sample 5 : (111) direction.

We give next the results of ρ versus $10^3/T$ for different uniaxial stresses. The stress is applied along the axis of the sample and is parallel to the current. We find the same behaviour for the three samples in (100) direction and the two samples in (111) direction.

Figure 4 shows on an enlarged scale the evolution of ε_1 , ε_2 , ε_3 with different stresses for sample 1. ε_1 decreases when the stress increases, whereas ε_2 increases. ε_3 varies slowly for weak stresses and then starts to increase also.

FIG. 4. — Resistivity as a function of $10^3/T$ for sample 1, (100) direction, for different uniaxial compression stresses on an enlarged scale. The three regions are separated (Fig. 3a, 3b and 3c). The slopes give ε_1 , ε_2 and ε_3 . The current and the stresses are in all cases parallel to the sample axis.

FIG. 5. — The same as figure 4, but for sample 5 (111) direction.

FIG. 6. — Longitudinal piezoresistance as a function of $10^3/T$ superimposed on the resistivity. The piezoresistance data given for different uniaxial compression static stresses are measured by a modulation stress technique. The superposition of the resistivity shows the correspondence between the change of the slope of the piezoresistance and the differently activated regions of the resistivity.

Figure 5 shows the behaviours of ϵ_1 , ϵ_2 and ϵ_3 with stress for sample 5. The comments are the same except that ϵ_3 increases more rapidly than in the preceding case.

When $10^3/T < 100$, we remark in ϵ_2 region, that ρ starts to increase with stress, passes through a maximum and then decreases. This fact was observed previously by Komatsubara *et al.* [2] and F. Pollak [15] on p type Ge and by Galavanov *et al.* [7], on p type InSb.

Remark. — In figure 2, the region corresponding to ϵ_2 was less apparent than the others, and in figure 6 we report the change of the piezoresistance with $10^3/T$ for sample 1. The three conduction processes appear very clearly : the piezoresistance is superimposed on the resistivity and one can see that the application of stress increases the part corresponding to ϵ_2 . This also appears in figures 4b and 5b. The complete piezoresistance data were published already [16]. In figure 7, we see the change of ϵ_1 , ϵ_2 and ϵ_3 as a function of strain for all our samples : We report ϵ_1 for samples 1 and 5 versus $1/\epsilon$ in figure 8. The extrapolation of ϵ_1 for $1/\epsilon = 0$ is obtained by a least squares fit.

Finally, table I gives the mean characteristics of our samples at room temperature : resistivity, Hall constant, Hall mobility, and acceptors concentration : Na. Na is obtained from the assumption that all the

FIG. 7. — The different energies of ionisation and activation ϵ_1 , ϵ_2 , ϵ_3 as a function of the strain. We remark that ϵ_1 is a function of sample axis, whereas ϵ_2 is not dependent on the sample axis. We say that ϵ_1 is anisotropic and ϵ_2 is isotropic (see values for $\epsilon = 0$).

TABLE I

Main characteristics of the samples used in this work

Samples	sample stress current } directions	ρ 300 K (Ωcm)	R_H 300 K (cm^3/C)	μ_H 300 K ($\text{cm}^2/\text{V.s}$)	N_A 10^{17} cm^{-3}	R ($\text{\AA}$)
1	(100)	7×10^{-2}	48.2	685	1.30	122
2	(100)	6×10^{-2}	48	800	1.30	122
3	(100)	1.6×10^{-1}	35.8	223	1.74	110
4	(111)	6.5×10^{-2}	43.5	670	1.43	119
5	(111)	6.5×10^{-2}	45.6	701	1.37	120

FIG. 8. — ε_1 as a function of $1/\varepsilon$ for sample 1 and sample 5. For sample 3 and samples 2 and 4, the behaviour is similar respectively to sample 1 and 5.

impurities are ionised at 300 K and R , the average impurities separation is deduced from the relation $R = (\frac{4}{3} \pi N_A)^{-1/3}$.

4. Discussion. — **4.1 HIGH TEMPERATURE REGION** ($T > 25$ K). — The region ε_1 is obviously the best known, ε_1 is the classical ionisation energy of the acceptor level [17] to [23].

Without strain and neglecting the spin-orbit interaction, the valence band edge is a six-fold degenerate p-type multiplet. The spin-orbit interaction splits the six-fold degeneracy into a four-fold $p_{3/2}$ (Γ_8) and a two fold $p_{1/2}$ multiplet (Γ_7). The effect of a uniaxial strain on the Γ_8 level is to split it into a pair of degenerate Kramers doublets separated by an energy E_s , which is a linear homogeneous function of the strain amplitude. Pikus and Bir [11] have shown that, for hole energies in the large strain limit, the resulting energy surfaces are ellipsoids whose principal axes coincide with those of the strain tensor. In this limit, the effective mass tensors for the two bands are independent of the strain amplitude and are functions of the strain geometry only. The two energy bands are given by :

$$E(\varepsilon, k) = \mp 1/2 E_s + 1/2 \hbar^2 k \gamma_{1,2} k \quad (1)$$

with

+ and γ_2 for heavy holes
— and γ_1 for light holes

γ_1 and γ_2 are the reciprocal mass tensors, which depend only on the relative values of the strain components ε_{ij} , k is the wave vector.

We neglect the small terms linear in k in eq. (1), coming from the different symmetry of GaSb compared to Ge; these effects can only be observed at low temperatures. E_s may be expressed in terms of two deformation potential constants b and d ; where b is an uniaxial orbital deformation potential and d is a shear deformation potential appropriate to strain of tetragonal symmetry :

$$E_s = 2 \sqrt{\frac{b^2}{2} \sum_{i>j} (\varepsilon_{ii} - \varepsilon_{jj})^2 + d^2 \sum_{i>j} \varepsilon_{ij}}$$

again ε_{ij} are the components of the strain tensor. We omit the energy displacement due to the hydrostatic part of the strain which is the same for the two bands.

For an infinite strain, the wave function of the acceptor ground state would be the product of Bloch functions from the top of the upper band only and an envelope function $F(r, \varepsilon)$ [15]; $F(r, \varepsilon)$ satisfies the effective mass equation :

$$\left\{ - \left[\frac{\hbar^2}{2 m_{\parallel}^*}(\varepsilon) \right] \frac{\partial^2}{\partial z^2} - \left[\frac{\hbar^2}{2 m_{\perp}^*}(\varepsilon) \right] \times \right. \\ \left. \times \left(\frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial y^2} \right) - \frac{e^2}{\chi r} \right\} F(r, \varepsilon) = \varepsilon_1(\varepsilon) F(r, \varepsilon) \quad (2)$$

$m_{\parallel}^*(\varepsilon)$ and $m_{\perp}^*(\varepsilon)$ are the effective masses parallel and perpendicular to the strain direction respectively.

In the limit of infinite strain, near $k = 0$,

$$\frac{m_0}{m_{\parallel}^*(\infty)} = A + B \frac{m_0}{m_{\perp}^*(\infty)} = A - \frac{B}{2} \quad \text{for (100) direction}$$

$$\frac{m_0}{m_{\parallel}^*(\infty)} = 1 + \frac{D}{\sqrt{3}} \frac{m_0}{m_{\perp}^*(\infty)} = A - \frac{\sqrt{3} D}{2} \quad \text{for (111) direction}$$

where A , B and D are the inverse cyclotron mass parameters and m_0 is the free electron mass. When r is large, an approximate solution of eq. (2) is given by :

$$F(r, \varepsilon) = [\pi a_{\parallel}^*(\varepsilon) a^*(\varepsilon)^2]^{-1/2} \times \exp - \left[\frac{z^2}{a_{\parallel}^*(\varepsilon)^2} + \frac{y^2 + x^2}{a_{\perp}^*(\varepsilon)^2} \right]$$

where $a_{\parallel}^*$ and $a_{\perp}^*$ are the effective Bohr radii related to ε_1 by :

$$a_{\parallel}^*(\varepsilon) = \hbar [2 m_{\parallel}^*(\varepsilon) \varepsilon_1(\varepsilon)]^{-1/2}$$

$$a_{\perp}^*(\varepsilon) = \hbar [2 m_{\perp}^*(\varepsilon) \varepsilon_1(\varepsilon)]^{-1/2}.$$

When E_s is reduced, the ground state function contains an increasing proportion of states from the light holes band, and we are no longer in the large strain limit. Price [24] however, has shown that for ε_1 the theory of the approach to large strain limit is straightforward and quite tractable. At least on the basis of the effective mass approximation, the results can be used to derive the deformation potential constants b and d from suitable data. Let us, first, expand ε_1 in terms of the inverse strain amplitude $1/\varepsilon$:

$$\varepsilon_1(\varepsilon) = \varepsilon_1(\infty) + \frac{\varepsilon'_1}{\varepsilon} + \frac{\varepsilon''_1}{\varepsilon^2} + \dots \quad (3)$$

Let us, now, expand the dispersion relation given by eq. (1) from the large strain limit to first order in $1/E_s$

$$E(\varepsilon, k) = \mp \frac{1}{2} E_s + \frac{1}{2} \hbar k \gamma_{1,2} \mp \frac{Z(k)}{E_s} \quad (4)$$

with

$$Z(k) = 3 B^2 (k_y^2 + k_x^2) \left[x_z^2 + \frac{1}{4} (k_y^2 + k_x^2) \right] + C^2 (k_x^2 k_y^2 + k_z^2 k_y^2 + k_x^2 k_z^2).$$

In the effective mass approximation, $\varepsilon_1(\varepsilon)$ is the ground state of the operator $H(i\nabla) + U(r)$ where $U(r)$ is the acceptor ion potential [24] (1). If the wave function of the acceptor ground state for infinite strain is expanded in terms of Bloch states of the upper band edge :

$$F_0(r) = \sum_m a_m \varphi_m(r)$$

we have

$$Z_0 = \sum_m |a_m|^2 Z(k)$$

or, in terms of the envelope wave function $F_0(r)$:

$$Z_0 = \int F_0 Z(-i\nabla) F_0 d^3r$$

Z_0 is equal for (100) direction strain to :

$$Z_0 = 2 B^2 \left(\frac{1}{a_{\perp}^{*4}} + \frac{1}{a_{\parallel}^{*2} a_{\perp}^{*2}} \right) + \frac{C^2}{3} \left(\frac{1}{a_{\perp}^{*4}} + \frac{2}{a_{\parallel}^{*2} a_{\perp}^{*2}} \right)$$

and for the (111) direction :

$$Z_0 = \frac{2 B^2}{3} \left(\frac{1}{a_{\perp}^{*4}} + \frac{1}{a_{\parallel}^{*2} a_{\perp}^{*2}} \right) + \frac{2 D^2}{9} \left(\frac{2}{a_{\perp}^{*4}} + \frac{1}{a_{\parallel}^{*2} a_{\perp}^{*2}} \right).$$

A calculation by second order perturbation of the change in the acceptor binding energy shows that [14] :

$$\varepsilon_1(\varepsilon) = \varepsilon_1(\infty) + \frac{Z_0}{E_s} - \frac{\varepsilon_1(\infty) Z_0}{E_s}. \quad (5)$$

If we compare equations (3) and (5) we can write

$$\frac{\varepsilon'_1}{\varepsilon} = \frac{Z_0}{E_s} \quad \text{and} \quad \frac{\varepsilon''_1}{2} = - \frac{\varepsilon_1(\infty) Z_0}{E_s}$$

and we obtain :

$$Z_0 = - \frac{\varepsilon_1(\infty) (\varepsilon'_1)^2}{\varepsilon''_1} \quad \text{and} \quad \frac{E_s}{\varepsilon} = - \frac{\varepsilon(\infty) \varepsilon'_1}{\varepsilon''_1}. \quad (6)$$

For an uniaxial strain in the (100) direction, a suitable expansion of E_s gives :

$$E_s = - 2 b(1 + r) \varepsilon$$

where r is the Poisson coefficient $= - S_{12}/S_{11}$ and S_{ij} are the elastic compliances. Thus we deduce :

$$b = - \frac{Z_0}{2(1 + r) \varepsilon'_1} = \frac{\varepsilon_1(\infty) \varepsilon'_1}{2 \varepsilon''_1(1 + r)} \quad (7)$$

where ε'_1 , ε''_1 , $\varepsilon_1(\infty)$ are obtained from our experimental results.

In the case where the strain is applied along the (111) axis, these results become :

$$E_s = - 2 d(1 + r') \frac{\varepsilon}{\sqrt{3}}$$

with

$$r' = - \frac{S_{11} + 2 S_{12} - 1/2 S_{44}}{S_{11} + 2 S_{12} + S_{44}}$$

$$d = - \frac{\sqrt{3} Z_0}{2(1 + r') \varepsilon'_1} = \frac{\sqrt{3} \varepsilon_1(\infty) \varepsilon'_1}{2 \varepsilon''_1(1 + r')}. \quad (8)$$

From the curves of figure 7 we deduce ε'_1 , ε''_1 and $\varepsilon_1(\infty)$. The results are given in table II. Finally, from the change of ε_1 as a function of strain, for samples in (111) and (100) directions with the strain parallel to the crystallographic axis of the sample, we obtain the following values of b and d :

$$b = - 2.3 \pm 0.2 \text{ eV}$$

$$d = - 7.0 \pm 0.2 \text{ eV}.$$

In table III we list the different values of b and d obtained from different experiments on GaSb. This

TABLE II

Results of a least squares fit on the data of figures 3 and 4. This table was given in ref. [10] with a mistake ; for $\varepsilon_1(\infty)$ the two values 6.3 meV and 11 meV were inverted

Samples	sample stress current } directions	$\varepsilon_1(\infty)$ meV	ε'_1 meV	ε''_1 meV	$\varepsilon_1(0)$ meV
1-2-3	(100)	6.3	5.75×10^{-3}	-6.25×10^{-6}	13.5×10^{-3}
4-5	(111)	11	2.3×10^{-3}	-2.5×10^{-5}	18×10^{-3}

TABLE III

Deformation potential constants b and d determined from experiments on strained p type GaSb

Ref.	b (eV)	d (eV)	Technic used
[25]	-1.8 ± 0.1	-4.2 ± 0.2	Wavelength modulated (direct optical edge)
[26], [31]	-3.3 ± 0.6	-8.35 ± 1.7	Modulated piezoreflectance (direct optical edge)
[27]	-2.0	-4.6	Piezoemission (direct optical edge)
[28]	-2.2 ± 0.2	-3.5 ± 0.4	Unknown
This work	-2.3 ± 0.2	-7.0 ± 0.2	Acceptor binding energy

shows that our deformation potentials, determined from the strain induced change in the acceptor binding energy [31] are different from those previously obtained from direct modulation spectroscopy. Such differences have been seen in Ge [29]. However, no quantitative theory is at present available for these effects and our experimental results only indicate that such effects are also important in III-V compounds; we can only make a qualitative comment. The hole-phonon interaction is determined by four deformation potential constants a , b , d and d_0 . The first three, pertaining to the acoustic scattering, also describe the change of valance band structure under strain. Consequently, they enter the theory of a large number of phenomena. d_0 is connected with the optical phonon scattering and so, appears only in the theory of transport phenomena. For pure acoustic scattering, there is a good overall agreement between available galvanomagnetic data and deformation potential obtained directly from optical experiments on strained GaSb. A slight discrepancy is due to ionized impurity scattering. At higher temperatures as in our case for region I, where the optical phonon scattering is efficient, the deviations are appreciable. The same phenomenon appears on Ge and Tolpygo [30] and subsequently Lawaetz [29] reconsidered the deformation potential theory of phonon scattering. At high temperatures e.g. in the range where optical phonon scattering is important a non trivial correction is pointed-out. The aim of the Lawaetz' theory is to show that this correction in the deformation potential approach to long-wavelength one phonon scattering in non-polar semiconductors is due to long range electrostatic forces.

4.2 INTERMEDIATE TEMPERATURES REGION ($20 < T < 10$ K). — In this region the conduction is due to the impurity band. When we apply a strain, the range of temperature corresponding to this conduction process becomes larger, and the limits are quite well obtained from our piezoresistance results.

Figure 7 shows that ε_2 is isotropic, i.e. it is not a function of the samples orientation. This phenomenon was already observed by Pollak [15] and suggests that this conduction process is thermally activated. Several theories exist which can explain this conduction process. The number of different theories shows that the problem is quite unclear. Let us remember simply the Mycielski [12], Frood [13] and Mikoshiba theories [14].

Mycielski assumes that the conduction is due to a hopping over the Coulomb's barrier between an occupied majority ion and an empty one. The energy of a carrier in the ground state is given by :

$$\varepsilon_1 = -\varepsilon_1^0 - \frac{e^2}{\chi R} \quad (9)$$

Where ε_1^0 is the ionisation energy of the isolated impurity, χ is the static dielectric constant, R is the average impurities separation.

The energy at the top of the barrier is

$$= E_2 = -4e^2/\chi R.$$

We deduce from the two last equations that

$$\varepsilon_2 = E_2 - \varepsilon_1 = \varepsilon_1^0 - \frac{3e^2}{\chi R}. \quad (10)$$

One remarks that $3e^2/\chi R$ is practically stress independent, except through R . However, we see from our curves (Fig. 7) that ε_2 at small strains increases with strain, this is not in agreement with Mycielski's theory.

Frood suggests that the conduction takes place in the lower tail of the valence band which includes the delocalized excited impurity states. This theory might explain the increase of ε_2 at low strains. Moreover the mechanism proposed by Frood is also thermally activated and it is in agreement with the isotropic behaviour of ε_2 . Unfortunately, this model is not in agreement with Hall constant and mobility variations vs. T . If we calculate the number of acceptors needed to have a tail of band according to the Kane's model [32, 33], we find approximately $N_A = 10^{18} \text{ cm}^{-3}$. The concentration in our samples were only 10^{17} cm^{-3} .

The last theory is due to Mikoshiba. At low temperatures, most of the holes are in the acceptor ground state, due to lack of thermal excitation, and they start to form a positively charged hydrogen-like ion. Because of the greater extension of the hole wave functions in the ion state, a band is formed with a greater mobility than in the band formed by the neutral acceptor ground states. For only one positive charge state, its energy is nearly equal to the neutral acceptor ionization energy ε_1 ; for several positive charge states, the exchange interaction β decreases their energy. So ε_2 is considered as the gap between the neutral acceptor ground state and the bottom of this so formed impurity band. This model is very well illustrated in figure 9 of ref. [15]. Under some assumptions, one obtains :

$$\varepsilon_2 = \varepsilon_1 - \beta \quad (11)$$

Mikoshiba calculation depends on two approximations :

i) the wave function of the positively charged state is approximated by a screened (1s) hydrogen wave function,

ii) the interaction potential $U(r)$ is approximated by a screened Coulomb interaction :

$$\beta = \frac{ne^2 s^2}{\chi a^*} \left(1 + \frac{sR}{a^*} \right) \exp\left(-\frac{sR}{a^*}\right) \quad (12)$$

We can calculate β for different samples :

a) *Sample 3* in (100) direction. We have $R = 110 \text{ \AA}$. $a^*(\infty)$ deduced from $a_{\uparrow}^*(\infty)$ and $a_{\downarrow}^*(\infty)$ by the relation $a_{\infty}^* = (a_{\uparrow}^*(\infty) a_{\downarrow}^*(\infty))^{1/3}$ is equal to $65 \text{ \AA}$, n is taken as adjustable parameter, function of neighbourhoods.

The screening parameter calculated by Slater's rules is 0.7 [15] :

$$\beta(\infty) \text{ is found equal to } 5.26 \times 10^{-3} \cdot n \text{ (eV)}$$

b) *Sample 5* : (111) direction :

$$\beta(\infty) = 2.79 \times 10^{-3} \cdot n \text{ (eV)} .$$

Consequently we can, from the extrapolated values of $\varepsilon_1(\infty)$ and $\varepsilon_2(\infty)$ in each of the three above cases, find n . We obtain

Sample 3 $n = 0.2$

sample 4 $n = 1.31$

sample 5 $n = 1.11$

FIG. 9. — This figure is due to Shklovskii [35] and shows the dependence of the activation energy ε_3 of hopping conduction in p and n type germanium weakly compensated on the concentration of principal impurities in weakly compensated samples : * phosphorous-doped germanium ; ◆ gallium doped germanium ; + antimony-doped germanium ; ◇ two antimony-doped germanium. The continuous straight line is the theoretical dependence of Shklovskii for $k = 0$ (eq. (14)). The dashed line represents the Miller and Abrahams equation (eq. (13)) for $k = 0$.

Remark. — The values of $\beta(\infty)$ show that β depends on the strain symmetry. This fact is not surprising, because ε_1 is a function of the strain direction whereas ε_2 is isotropic.

From the relation $\varepsilon_1 = \frac{e^2}{2\chi a^*}$ we obtain $a^*(0)$ which is the effective Bohr radius at zero strain :

$$\frac{\varepsilon_1(\infty)}{\varepsilon_1(0)} = \frac{a^*(0)}{a^*(\infty)}.$$

For samples 4 and 5 we find :

$$a^*(0) = 28.6 \frac{11}{18} = 17.4 \text{ \AA}.$$

We can check this result from the definition of the effective Bohr radius

$$a^* = \frac{\chi \hbar}{m^* e}$$

with

$$\chi = 13.7 \chi_0; \quad \hbar = 6.63 \times 10^{-27}$$

and

$$m^* = 0.42 m_0$$

we find : $a^*(0) = 17.3 \text{ \AA}$ in excellent agreement with our value of $17.4 \text{ \AA}$. We choose samples in the (111) direction rather than in the (100) direction, because the heavy holes band, is displaced toward the (111) direction due to the presence of linear k terms in the expression of the energy. On the other hand the value of $m^* = 0.42 m_0$ is very close to the known value of the heavy holes mass.

For sample 3 in the (100) direction,

$$a^*(0) = \frac{\varepsilon_1(\infty)}{\varepsilon_1(0)} a^*(\infty)$$

which gives $a^*(0) = 30 \text{ \AA}$. But we cannot further check this result, because we have in this case a very large uncertainty in the choice of the mass : the admixture of light holes is more important than in the first case.

Remark. — We will not discuss sample 1 because $\varepsilon_2(\infty)$ is larger than $\varepsilon_1(\infty)$. This result is quite surprising and is not understood.

4.3 LOW TEMPERATURES REGION ($T < 10 \text{ K}$). —

At low temperatures our samples exhibit an activation energy ε_3 which corresponds to hopping conduction. This process was studied by Fritzche [4], Mott and Twose [34], Kohn [35], Miller and Abrahams [36] then later by Shklovskii [37].

At low concentration, bonding does not occur and the conduction takes place by hopping of holes from occupied to empty localised acceptor states. The concept of a hole bound to an acceptor centre is complementary to that of an electron bound to a donor.

Hence our discussion for p type donors is carried over from hopping in n type material. Mott [34] has shown that there is hopping conduction when R , the average acceptor separation, is larger than three times the effective Bohr radius a^* . This conduction process is a function of the compensation K . The samples studied in this work are not intentionally doped and we make the reasonable assumption that K is small.

Several theories are given for hopping conduction. The most important are due to Miller and Abrahams [36] in one hand and to Shklovskii [37] in other hand. The principal feature of Miller and Abrahams theory is that a phonon induced hole hops from acceptor site to acceptor site, where a fraction K of the total sites concentration is vacant due to compensation. To first order in the electric field, the solution to the steady state and current equations is equivalent to the solution of a linear resistance network. The resistance network is evaluated and the result shows that the T dependence of the resistivity is

$$\rho \propto \exp(\varepsilon_3/kT).$$

For small K , ε_3 is given by :

$$\varepsilon_3 = \frac{e^2}{\chi} \left(\frac{4 N_A \pi}{3} \right)^{1/3} (1 - 1.35 K). \quad (13)$$

The expression given by B. I. Shklovskii [37] differs quite strongly from equation (13). The difference arises from a second vacant acceptor located near the donor, whose potential strongly alters the energy of the first acceptor. We do not give the detail of the theory [37, 38]. The result is :

$$\varepsilon_3 = 0.99 \frac{e^2}{\chi} N_A^{1/3} (1 - 0.3 K^{1/4}). \quad (14)$$

At low K values, the difference between eq. (13) and (14) is as high as 60 %. Figure 9, due to Shklovskii, shows the dependence of the experimentally determined activation energy ε_3 on the concentration of the principal impurities in weakly compensated samples of n and p type germanium. The dashed line shows the dependence obtained by equation (13) for $K = 0$. The solid line shows the dependence obtained by equation (14) for $K = 0$.

We find that Shklovskii's law for small value of N_D or N_A is in good agreement with experimental results. When $N_D(N_A)$ increases the activation energy ε_3 deviates from Shklovskii's law. This can be explained as follows : at low values of K , the levels of most of the donors or acceptors are little perturbed and their energies are very close to one another, the overlap of the wave functions of the resonant states therefore leads to the formation of an impurity band. This modifies the description of the hopping process investigated in ref. [37].

For p type GaSb, on using Shklovskii model we find for samples 1 and 5 the values of ε_3 given in

TABLE IV

Change in ε_3 with the strain for the samples 1 and 5 where N_A is 1.3×10^{17} and $1.37 \times 10^{17} \text{ cm}^{-3}$ respectively. We remark that the application of the strain which decreases the volume of the wave function of the acceptor states, increases the value of ε_3 .

Experimental	Sample 1	1.8	2.0	2.4	2.7	2.8	5.5	calculated by eq. (14)
activation energy	Sample 5	1.8	2.25	2.6	3.0	3.2	5.2	calculated by eq. (14)
ε_3	Strain applied	0	1.48×10^{-3}	2.2×10^{-3}	2.8×10^{-3}	3×10^{-3}		

table IV. We assumed K to be zero and a pure hopping process. The experimental values are quite different. This fact is not surprising as R is $120 \text{ \AA}$, i.e. N_A is about 1.3×10^{17} . Hence we are in a region where the deviation from equation (14) is important. When we applied a strain, the spatial extension of the wave function decreases as it shown by the variation of a^* [4]. The overlap between the neighbouring acceptor states also decreases. The application of a pressure should therefore lead to a decrease of the difference between experimental values of ε_3 and the calculated values. This is shown in table IV : when ε increases ε_3 also increases. It seems evident that if N_A is small (10^{16} cm^{-3}) the influence of the pressure is very small but we had no sample to check this point. Finally we remark that on the basis of Shklovskii's theory it is possible to explain the large piezoresistance data on p type GaSb at very low temperatures. This part of our work has been published elsewhere [16].

5. Conclusion. — We have applied uniaxial compressional stresses in the (111) and the (100) directions

on p type GaSb. From the variation of the acceptor binding energy, we obtain the values of the deformation potential constants b and d . We find $b = -2.3 \text{ eV}$ and $d = 7.0 \text{ eV}$. We compare the data obtained directly by modulation spectroscopy with the change in acceptor binding energy under uniaxial stress.

From Kikoshiba's theory for the impurity band conduction region, we calculate the effective Bohr radius $a^*(0)$ for samples in the (100) direction. The calculated value of $a^*(0)$ is compared with the value obtained by the definition of the effective Bohr radius. The very good agreement between these values, is a strong indication of the validity of the Mikoshiba theory. Finally in the low temperature region, where the conduction takes place by hopping, when we apply an uniaxial compressional stress the activation energy increases i.e. the overlap between the wave functions of neighbour acceptor states decreases and we find that the theoretical model of Shklovskii offers a better description of the hopping process.

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- [31] *Note* : values given by Gavini and Cardona for GaSb appear slightly higher than other results. Also, if we look on table V of ref. [26], we see that in their set of d values, the one given for GaSb is higher than the others. If we apply the point-ion model proposed by these authors to take account of the partially ionic character of III-V compounds Δd is equal to -0.69 eV starting from an average value of d between germanium and grey tin we find a value smaller than -5 eV in good agreement with most of the experimental results of table I.
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