1964, edited by M. Hulin (Academic, New York, 1964), p. 423.

[12]M. H. Brodsky, R. S. Title, K. U. Weiser, and G. D. Pettit, Phys. Rev. B 1, 2632 (1970); M. Brodsky and R. S. Title, Phys. Rev. Lett. 23, 581 (1969).

[13]D. Beaglehole and M. Zavetova, in International Conference on Amorphous and Liquid Semiconductors, Cambridge, England, 24-27 September 1969 (to be published).

[14]F. G. Allen and G. W. Gobeli, Phys. Rev. 127, 150 (1962).

[15]F. G. Allen and C. W. Gobeli, J. Appl. Phys. 35, 597 (1964).

[16]H. Fritzsche, Bull Amer. Phys. Soc. 14, 342 (1969).

[17]S. C. Moss and J. F. Graczyk, Phys. Rev. Lett. 23, 1167 (1969).

[18]C. W. Peterson, J. H. Dinan, and T. E. Fischer, Bull. Amer. Phys. Soc. 15, 70 (1970).

MEASUREMENT OF n-TYPE GERMANIUM MICROWAVE CONDUCTIVITY DURING IMPACT IONIZATION OF IMPURITIES AT 4.2°K

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We have measured the microwave conductivity of n-type germanium with $N_D-N_A \sim 10^{14}$ cm^{-3} during impact ionization at 4.2°K, by means of a cavity perturbation method which requires no knowledge of the electronic spatial distribution. Values of the electronic temperature derived from the microwave conductivity are compared with theoretical predictions and Hall mobility data at higher temperatures.

We present a measurement of the complex microwave conductivity of n-type germanium at 4.2°K during impact ionization of neutral impurities (arsenic, N_D in the range 10^{14} cm^{-3}), using a dielectric cavity perturbation method. The cavity, containing dielectric of relative permittivity $K \sim 9$, is rectangular and operates at a central frequency $F_0 = 4065$ MHz in the TE_{011} mode. The germanium rod is introduced as shown in Fig. 1. The measurements can be performed either with stable dc current across the sample, with the dc electric field parallel to the rf electric field, or with short pulses to check whether or not heating occurs. rf measurements are made at a very low power level to avoid nonlinear effects. The apparatus is immersed in liquid helium and is protected from light. Variations of the resonant frequency F_0 and the quality

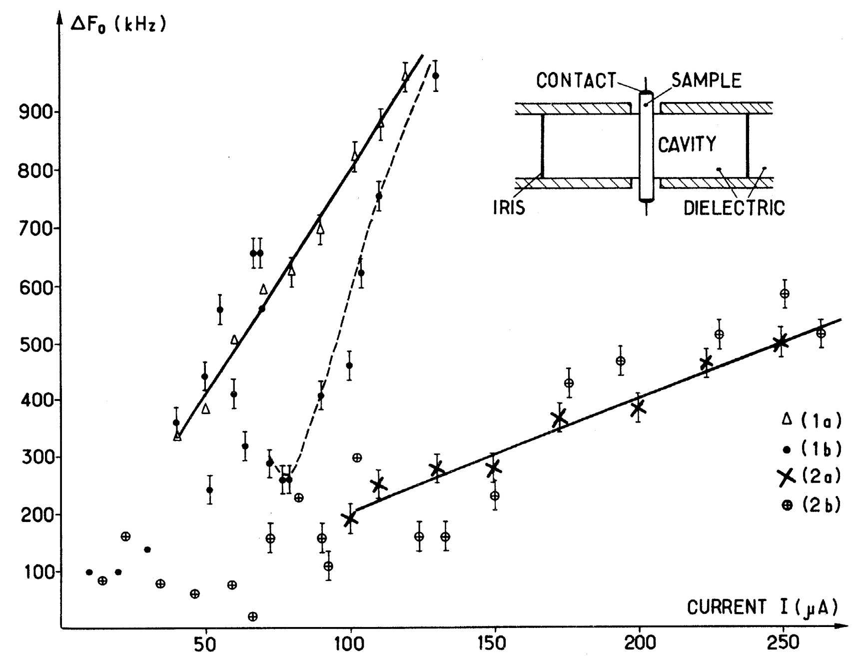

FIG. 1. Resonance-frequency shifts versus the current in the samples, and schematic drawing of the apparatus.

Table I. Derivation of electronic temperature ϵ/k_0 from the theoretical value of τ_p and the energy-balance equation (3). The difference between μ and μ_p is close to the experimental uncertainty.

Sample	N_D (cm^{-3})	N_A (cm^{-3})	E (V cm^{-1})	Z (av)	μ, expt. ($m^2\ V^{-1}\ sec^{-1}$)	ϵ/k_0 (°K)	$\mu_p(\epsilon/k_0)$ ($m^2\ V^{-1}\ sec^{-1}$)	$\tau_2(\epsilon)$ (10^{-9} sec)
1(a)	2×10^{14}	6×10^{13}	7.5	0.35	45	126	55	0.5
2(a)	4×10^{14}	10^{14}	15	0.20	24	210	32	0.4

factor Q are measured by two precision wavemeters incurring a frequency error of less than 100 kHz. These variations are given by the usual perturbation calculation,

$$\Delta F_0/F_0 = -\varphi\,\mathrm{Im}\hat{\sigma}/2,$$
$$\Delta(1/Q) = \varphi\,\mathrm{Re}\hat{\sigma}, \tag{1}$$

where Re$\hat{\sigma}$ and Im$\hat{\sigma}$ are the real and imaginary parts of the microwave conductivity, and φ is a shape factor depending on the dielectric constants involved and proportional to r^2, r being the plasma radius. Equations (1) are valid if the TE_{011} mode is weakly perturbed in the plasma, i.e., if the penetration depth of rf waves is much larger than $2r$. Such a condition provides a limit of current far above our range of measurements. We notice that even with no current in the sample the TE_{011} mode is perturbed, for the dielectric constant of germanium is different from that of the cavity. This fact does not change the validity of (1), considering only the second of two successive perturbations of the initially dielectric-filled cavity. The variations ΔF_0 and $F_0\Delta(1/Q)$ measured with samples cut in two differently doped ingots (see Table I for impurity concentrations) are shown in Figs. 1 and 2.

Samples are referred to as 1(a), 1(b), 2(a), and 2(b), the number indicating the ingot and the letter the type of contact: (a), 3% As- or Sb-doped lead-tin solder; (b), electroless plated nickel. Contact (a) provides the best results, i.e., quite linear variation of $F_0\Delta(1/Q)$ and ΔF_0 with I, but does not give reliable microwave data below a certain critical current I_s. Below I_s we observe a low-frequency instability related to a 5% negative drop of the voltage at the onset of breakdown; this drop cannot be significantly attributed to an intrinsic process such as in Melngailis and Milnes.[1] We do not understand clearly the (b) data and attribute the dispersion to the weakness of contact (b) at 4.2°K. We shall discuss hereafter the results obtained with contact (a) and $I > I_s$.

In order to interpret the results obtained for the shape-independent quantity $Z = \Delta F_0/F_0\Delta(1/Q)$, we derive the microwave conductivity from the usual method of calculation,[2] neglecting the ionization collisions which are supposed to be far less frequent than other scattering processes such as acoustic-phonon and ionized-impurity scattering. In the steady state, the mobility μ and mean energy ϵ are given by

$$\mu = q\tau_1/m_c = q\tau_p\tau_i/m_c(\tau_p+\tau_i), \tag{2}$$
$$q\mu E^2 = (\delta\epsilon)_p/\tau_p = \epsilon/\tau_2, \tag{3}$$

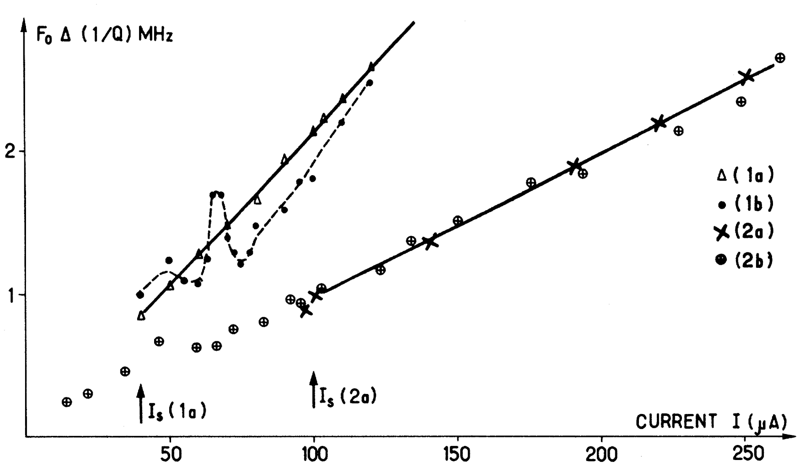

FIG. 2. Relative Q variations versus the current in the samples. I_s refers to the current below which instability occurs in the case (a), and no correct microwave measurement can be performed.

where τ_p and τ_i are momentum relaxation times related to acoustic-phonon and ionized-impurity scattering, m_c is the mobility effective mass ($m_c = 0.12 m_0$), and $(\delta\epsilon)_p$ is the mean energy loss of electrons during collisions with phonons. We moreover assume that $(\delta\epsilon)_p = (2m^*\epsilon)^{1/2}s$, where m^* is the density-of-states effective mass ($m^* \simeq 0.25 m_0$) and s the velocity of sound in germanium. Thus, in our frequency range, the complex conductivity $\hat{\sigma}$ is given by

$$\hat{\sigma} = \sigma_0[1.5 + i\omega\tau_2 - N(\epsilon)]/[N(\epsilon) + (1 + i\omega\tau_1)(1.5 + i\omega\tau_2)], \qquad (4)$$

where σ_0 is the steady-state conductivity and $N(\epsilon) = (\tau_i - 1.5\tau_p)/(\tau_p + \tau_i)$; τ_p varies as ϵ^{-1} [3,4] when $k_0 T < 0.1\epsilon$ (T is the lattice temperature), and τ_i varies as $\epsilon^{1.5}$. We have neglected, in deriving (4), the modulation of the electron density n by the rf field. This latter approximation is valid, for the values of F_0 considered, when the impurity concentration is less than 10^{15} cm^{-3} as we verified following theoretical and experimental studies.[5,6] Assuming that $\tau_2 \gg \tau_1$ and $(\omega\tau_2)^2 \gg 1$, Eqs. (1) and (4) lead to

$$2Z = 2\Delta F_0/F_0 \Delta(1/Q) = \omega\tau_1 - 2N(\epsilon)/\omega\tau_2. \qquad (5)$$

As we ignore a priori the value of $N(\epsilon)$ [+1 if phonon scattering predominates, -1.5 for ionized impurities, and, estimating roughly, $\omega^2\tau_1\tau_2 \sim 10$ from Eq. (3)], we drop the second term in the second member of (5); thus, experimental values of Z give $\mu = q\tau_1/m_c$. In order to relate the measured mobility to the electronic temperature ϵ/k_0 from Eq. (3), it is necessary to assume some theoretical dependence of τ_p or τ_2 on ϵ. Neglecting the heating of phonons by hot electrons,[3] which occurs in the avalanche saturation regime, we use the zero-point approximation[3,4] to obtain $\tau_p(\epsilon) = 4.6 \times 10^{-9}(\epsilon/k_0)^{-1}$. Then Eq. (3) leads to values of ϵ/k_0 listed in Table I.

To verify the consistency of the results we compare our experimental value of μ with $\mu_p = q\tau_p/m_c$ calculated from ϵ/k_0 deduced above; one can see in Table I that μ_p is close to μ, from which one infers that phonon scattering is predominant. From the Brooks-Herring formula, with the assumption of singly ionized sites $2N_A$ in number, we obtain the theoretical partial mobility related to ionized-impurity scattering, $\mu_i = q\tau_i/m_c = 200$ m^2 V^{-1} sec^{-1} with $\epsilon/k_0 = 100$°K and $N_A = 10^{14}$ cm^{-3}, in good agreement with $\mu_i \gg \mu_p$. We note that at $T = 10$°K the low-field mobility is somewhat lower than for hot electrons at 4.2°K; we controlled this in our auxiliary experiment [low-field dc Hall measurements as a function of T to obtain N_A as in Ref. 5: $\mu(\epsilon/k_0 = 10°\text{K}) = 7$ m^2 V^{-1} sec^{-1}]. We state also the deduced values of $\tau_2(\epsilon)$ in Table I; as far as we are aware, no direct measurement of $\tau_2(\epsilon)$ has been given.

To check the consistency of $\omega^2\tau_1\tau_2 \gg |2N|$, we have to calculate $\omega^2\tau_1\tau_2$ with our deduced values of ϵ; we see that the error is significant only for sample 2 (25%). For this sample, the electronic temperature is probably lower than the 210°K estimated; we could calculate ϵ/k_0 from the entire Eq. (5), neglecting the τ_i effect, but there is another source of error we have up to now neglected: neutral-impurity scattering. With 5×10^{14} As donors, the partial mobility relative to this scattering is 60 m^2 V^{-1} sec^{-1} only, which contributes to modify Eq. (5) in altering $N(\epsilon)$, and correction is quite ambiguous.

One can compare the mobilities obtained with Koenig's data[5] between 7 and 9°K; at least for sample 1, we obtain here similar values for $\mu(E)$: The mobility of hot electronics is independent of T when $\epsilon > 10k_0T$ and the agreement is good. Unfortunately we were not able to compare these results with prebreakdown Hall data at 4.2°K, as noise appeared in the Hall voltage below $T = 7$°K and remained up to saturation. We relate this transverse noise to former experiments[7] subject to the hypothesis of a great nonuniformity of the electronic spatial distribution, even though our samples are moderately compensated. This fact induced us to consider the more suitable microwave experiment presented here.

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[1] I. Melngailis and A. G. Milnes, J. Appl. Phys. 33, 3 (1962).

[2] A. F. Gibson, J. W. Granville, and E. G. S. Paige, J. Phys. Chem. Solids 19, 198 (1961).

[3] E. M. Conwell, *High Field Transport in Semiconductors*, Suppl. 9 to *Solid State Physics* (Academic, New York, 1967).

[4] E. M. Conwell, Phys. Rev. 135, A814 (1964).

[5] S. Koenig *et al.*, Phys. Rev. 128, 1668 (1962).

[6] V. A. Chuenkov, Fiz. Tekh. Poluprov. 2, 353 (1968) [Sov. Phys. Semicond. 2, 3 (1968)].

[7] K. Komatsubara *et al.*, Jap. J. Appl. Phys. 3, 2 (1964).