Magnetoconductivity of Hot Electrons

HERBERT F. BUDD
Research and Engineering Center, Xerox Corporation, Webster, New York
(Received 25 June 1965)

A theoretical analysis of galvanomagnetic phenomena in an intense electric field is given. The Boltzmann transport equation is solved for the case of acoustic-phonon and impurity scattering and the magnetoconductivity tensor is calculated. The dependence of the Hall and conductivity mobilities on electric and magnetic fields is determined both for spherical constant-energy surfaces and a band structure of the many-valley type. Expressions are derived for both the Hall and conductivity mobilities in the limit of small magnetic fields, where one considers quadratic variations of mobility in the magnetic fields, as well as the case of large magnetic fields. Negative magnetoresistance obtains for the case of spherical constant-energy surfaces and predominant acoustic-phonon scattering, whereas the sign of the mobility variation is shown to depend sensitively on the many-valley band structure and intervalley transition processes.

IN this paper we consider the transport properties of charge carriers in the presence of electric and magnetic fields. We shall be particularly concerned with the calculation of the magnetoconductivity tensor for the case of intense electric fields.

Sodha and Eastman[1] have calculated the electric-field dependence of the small-magnetic-field Hall coefficient for a band structure with spherical constant-energy surfaces and scattering by acoustical phonons. Conwell[2] has derived general expressions for the magnetoconductivity tensor in a many-valley semiconductor in terms of the electron distribution function. As we shall discuss in Sec. I, the calculation of the distributlon function must be carried out in a self-consistent manner, in order to properly take account of the effect of the Hall field on the energy distribution. This complication does not arise for weak electric fields, since in this case the energy distribution, or the isotropic part of the distributlon function, is independent of both the electric and magnetic fields.

The author[3] has computed the small-magnetic-field Hall coefficient for hot electrons in a many-valley semiconductor as well as the distributlon function for arbitrary electric and magnetic fields for the case of acoustic-phonon interactions. The small-magnetic-field Hall mobility in n-type germanium has been studied by Das and Nag.[4] The effect of phonon non-equipartition on the low-magnetic galvanomagnetic properties has also been treated.[5,6] More recently, Matz and Garcia-Moliner have studied non-Ohmic transport phenomena in a magnetic field with emphasis on galvanomagnetic effects in the absence of a Hall field.[7]

In Sec. I of this paper, we shall first discuss the Boltzmann equation relevant to hot electrons and its solution for the case of acoustic-phonon and impurity scattering. Expressions for the field dependence of the conductivity and Hall mobilities are given in Sec. II and it is shown that under hot-electron conditions the conductivity mobility increases with increasing magnetic field, i.e., negative magnetoresistance. In Sec. IV we consider some of the effects associated with a many-valley band structure, in particular the dependence of the average electron energy in a given valley on *odd* powers of the magnetic field. The extreme sensitivity of the galvanomagnetic properties of hot electrons both to the band structure and to intervalley repopulation phenomena is discussed in this section.

1. THE TRANSPORT EQUATION

The time-independent Boltzmann equation for the distribution function $f(p)$ is given by:

$$e(\mathbf{E}+\mathbf{V}\times\mathbf{B})\cdot\nabla_p f=\hat{C}f, \tag{1}$$

where $\mathbf{E}$ and $\mathbf{B}$ are the electric- and magnetic-field vectors, C is the collision operator, and $\mathbf{p}$ and $\mathbf{V}$ are the momentum and velocity vectors, respectively. The distribution function f is written as the sum of an isotropic term $S(\epsilon)$, depending only on energy ϵ, and an anisotropic term A. We consider first the case of spherical constant-energy surfaces for which the following relation between S and A has been derived,[3]

$$A=-\tau\frac{dS}{d\epsilon}\mathbf{V}\cdot\left[\frac{e\mathbf{E}+(e^2\tau^2/m^2)\mathbf{B}\mathbf{B}\cdot e\mathbf{E}+(e\tau/m)e\mathbf{E}\times\mathbf{B}}{1+(e^2\tau^2/m^2)B^2}\right], \tag{2}$$

where m is the effective mass. In obtaining this result it has been assumed that the effect of collisions on the anisotropic part of the distribution function may be represented by an energy-dependent relaxation time $\tau(\epsilon)$

$$\hat{C}A=-A/\tau(\epsilon).$$

The isotropic part of the distribution function S satisfies

[1] M. S. Sodha and P. C. Eastman, Phys. Rev. **110**, 1314 (1958).
[2] E. M. Conwell, Phys. Rev. **123**, 454 (1961).
[3] H. F. Budd, Phys. Rev. **131**, 1520 (1963).
[4] P. Das and B. R. Nag, Proc. Phys. Soc. (London) **82**, 923 (1963).
[5] H. F. Budd, Phys. Rev. **134**, A1281 (1964).
[6] H. F. Budd, in *Proceedings of the International Conference on the Physics of Semiconductors, Paris*, 1964 (Dunod Cie., Paris, 1964), T2-4.
[7] D. Matz and F. Garcia-Moliner, Phys. Status Solidi **5**, 495 (1964).

the following equation[3]:

$$-\frac{2(eE)^2}{3m\epsilon^{1/2}}\frac{d}{d\epsilon}\left[\tau\epsilon^{3/2}\frac{dS}{d\epsilon}\left(\frac{1+\gamma_0^2\omega^2\tau^2}{1+\omega^2\tau^2}\right)\right]=\hat{C}S, \qquad (3)$$

where γ_0 is the direction cosine of $\mathbf{B}$ with respect to $\mathbf{E}$ and $\omega=eB/m$.

It is important to realize that the electric field which appears in these equations is not simply the applied field, but is the sum of the applied and Hall fields. This total field is determined from the condition that the current be constrained to flow in the direction parallel to the applied electric field. The self-consistent determination of the total electric field is complicated by the fact that the isotropic part of the distribution function is field dependent, whereas it is merely the thermal equilibrium distribution in the case of weak electric fields, i.e., Ohmic conductivity.

We take the magnetic field to be in the z direction of a Cartesian coordinate system and the applied electric field to be in the x direction. Thus, the electric and magnetic fields have the following components:

$$\mathbf{E}=(E_x,E_y,0), \quad \mathbf{B}=(0,0,B), \qquad (4)$$

where E_x and E_y are the applied and Hall fields, respectively. For this configuration γ_0 vanishes. The current density is simply given by

$$\mathbf{J}=ne\langle\mathbf{V}\rangle=ne\int\mathbf{V}A\,d\mathbf{p}\bigg/\int S\,d\mathbf{p}. \qquad (5)$$

Inserting Eq. (2) in Eq. (5) we obtain the following expressions for the average velocity:

$$\langle\mathbf{V}\rangle=\boldsymbol{\mu}\cdot\mathbf{E}, \quad \boldsymbol{\mu}=\begin{pmatrix}\mu_{11} & \mu_{12} & 0\\ -\mu_{12} & \mu_{11} & 0\\ 0 & 0 & \mu_{11}\end{pmatrix} \qquad (6)$$

$$\mu_{11}=\frac{e}{m}\left\langle\frac{\tau}{1+\omega^2\tau^2}\right\rangle, \quad \mu_{12}=\frac{e}{m}\omega\left\langle\frac{\tau^2}{1+\omega^2\tau^2}\right\rangle, \qquad (7)$$

where

$$\langle x\rangle\equiv\int\epsilon^{3/2}(dS/d\epsilon)X(\epsilon)d\epsilon\bigg/\int\epsilon^{3/2}(dS/d\epsilon)d\epsilon. \qquad (8)$$

The constraint of zero current in the y direction requires:

$$E_y/E_x=\mu_{12}/\mu_{11}. \qquad (9)$$

Since from Eq. (3) S depends on the magnitude of the total electric field $E^2=E_x^2+E_y^2$, Eq. (9) is an implicit equation which must be solved in order to determine E_y.

The qualitative effect of the Hall field may be seen from the following simple picture. The Lorentz force arising from the magnetic field deflects the carriers away from the x direction and thus diminishes the power input due to the applied field. We note in fact, from Eq. (3) that for very large magnetic fields ($\omega\tau\gg1$) the electric and magnetic fields appear only in the ratio $(E/B)^2$ for our configuration ($\gamma_0=0$). If $\mathbf{E}$ were simply the applied field, i.e., if there were no Hall field, this ratio would tend to zero for large B, and the isotropic part of the energy distribution would approach the field-independent thermal-equilibrium distribution. This corresponds to the charge carriers drifting in the $\mathbf{E}\times\mathbf{B}$ direction and thus receiving no power from the electric field. The Hall field counteracts the Lorentz field on the average and assures that there be no net current normal to the applied electric field, thus allowing the carriers to receive power from the applied field.

II. THE DISTRIBUTION FUNCTION

Let us first consider the case where the principle energy losses are due to acoustic-phonon interactions and further assume that the phonon distribution remains in equilibrium. Hot-electron phenomena for nonequilibrium distributions have been recently discussed by Conwell and Zylbersztejn.[8,9]

For the case of phonon equipartition, the author[3] has obtained the following solution to Eq. (3).

$$S=\exp\left[-\int_0^\epsilon\frac{(d\epsilon/kT)}{1+[2(eE)^2\tau/3mC_1\epsilon^{1/2}]((1+\gamma_0^2\omega^2\tau^2)/(1+\omega^2\tau^2))}\right], \qquad (10)$$

where C_1 is related to the electron-phonon coupling constants and is detailed in Ref. 3. We consider only the configuration discussed above, where $\gamma_0=0$.

Momentum relaxation by acoustic phonons and impurities will be treated, the latter process by the Conwell-Weisskopf[10] model. The energy-dependent relaxation time is then given by

$$1/\tau=1/\tau_a+1/\tau_I; \quad \tau_a=C_2/\epsilon^{1/2}, \quad \tau_I=I\epsilon^{3/2}, \qquad (11)$$

where C_2 and I are given by the usual expressions.

We shall not consider impurity scattering in detail since we wish to emphasize the negative magnetoresistive effects associated with the hot-electron regime. As the electric field is increased, and correspondingly the mean electron energy and relaxation time for impurity scattering, the role of acoustic phonons rapidly in-

[8] A. Zylbersztejn and E. M. Conwell, Phys. Rev. Letters **11**, 417 (1964).
[9] E. M. Conwell, Phys. Rev. **135**, A814 (1964).
[10] Ester M. Conwell and V. F. Weisskopf, Phys. Rev. **77**, 388 (1950).

creases. Conwell[11] has treated the role of impurity scattering in warm-electron phenomena at low temperatures and modest electric fields where it is of primary importance.

Upon inserting Eq. (11) into Eq. (10) and performing the necessary integration, one obtains a rather unwieldy expression for the energy distributlon function. We shall not follow this procedure here, but shall limit our considerations to several special cases of interest.

A. Acoustic Scattering Only

In this case the integration of Eq. (10) yields:

$$S=e^{-\epsilon/kT}[\epsilon/kT+p+\omega^2C_2^2/kT]^p, \qquad (12)$$

where $p=(3\pi/16)(E\mu_a/s)^2$. Here μ_a is the low-field mobility for acoustic phonon scattering and s is the sound velocity.

This is the magnetic-field generalization of the Yamashita-Watanabe[12] distribution and reduces to the latter when $\omega=0$. The distribution function for an arbitrary angle between the electric- and magnetic-field vector has been derived by the author.[3]

It should be re-emphasized here that p is a function of the magnetic field by virtue of the Hall field:

$$p=p_0[1+(E_y/E_x)^2]; \quad p_0=(3\pi/16)(E_x\mu_a/s)^2. \qquad (13)$$

B. Impurity and Acoustic-Phonon Scattering

We first consider the case of zero magnetic field, and obtain the following distribution function from Eq. (10):

$$S=e^{-x}[x^2+px+n]^{p/2}e^{-(p^2/2)\phi},$$
$$x=\epsilon/kT, \quad n=C_2/I(kT)^2, \qquad (14)$$

where

$$\phi=\frac{2}{[4n-p^2]^{1/2}}\tan^{-1}\left(\frac{2x+p}{[4n-p^2]^{1/2}}\right), \qquad 4n>p^2$$
$$=\frac{1}{[p^2-4n]^{1/2}}\ln\left[\frac{2x+p-[p^2-4n]^{1/2}}{2x+p+[p^2-4n]^{1/2}}\right], \quad 4n<p^2.$$

When there is no impurity scattering $n=0$ and Eq. (14) reduces to the Yamashita-Watanabe distribution. For large electric fields we neglect the constant term in the denominator of the integrand in Eq. (10) and we obtain the following result for nonzero magnetic field;

$$S=\exp\left[-\frac{(x+\omega^2C_2^2/kT)^2}{2p}\right]\times\left[\frac{\exp([\omega^2C_2^2/pkT]n^{1/2}\tan^{-1}[x/n^{1/2}])}{x^{n/p}}\right]. \qquad (15)$$

The first term in Eq. (15) is the high-electric-field distribution function in the absence of impurity scattering and the second term takes account of this additional scattering mechanism. We note that for large electric fields the contribution of impurity scattering diminishes rapidly.

C. Large Magnetic Fields

When $\omega\tau\gg1$ Eq. (10) may be simply evaluated for combined impurity and phonon scattering; we obtain

$$S=\exp(-x/(1+\alpha))\exp(a/(1+\alpha))(\tan^{-1}x/a), \qquad (16)$$

where

$$a=[n\alpha/(1+\alpha)]^{1/2}, \quad \alpha=\tfrac{1}{3}(E/Bs)^2.$$

In the absence of impurity scattering ($n=0$) the energy distribution approaches a simple Maxwellian with an effective temperature $T^*=T[1+\alpha]$.

III. THE MAGNETOCONDUCTIVITY TENSOR

In this section we shall compute the conductivity and Hall mobilities for hot electrons and acoustic-phonon scattering. Let us first consider the case of small magnetic fields and discard all but the lowest order terms in ω. We shall write the conductivity and Hall mobilities in the following form for small magnetic fields:

$$\mu_c=\langle V_x\rangle/E_x=\mu_c^0[1+a(\mu_c^0B)^2],$$
$$\mu_H=E_y/BE_x=\mu_H^0[1+h(\mu_H^0B)^2]. \qquad (17)$$

For the case of intense electric fields and spherical constant-energy surfaces we merely insert Eq. (15) into Eq. (6) and obtain the following results:

$$\mu_H^0=e\bar{\tau}\Gamma(5/4)/m(2p_0)^{1/4}\Gamma(\tfrac{3}{2}),$$
$$\mu_c^0=2e\bar{\tau}\pi^{1/2}/3m(2p_0)^{1/4}\Gamma(\tfrac{3}{4}), \quad \bar{\tau}\equiv C_2/(kT)^{1/2},$$
$$h=[2(\Gamma(\tfrac{3}{2})/\Gamma(5/4))^2(1/\Gamma(\tfrac{3}{2})-3\Gamma(\tfrac{3}{4})/\Gamma(\tfrac{1}{4}))-\tfrac{1}{4}]=-0.031, \qquad (18)$$
$$a=9/8(\Gamma(\tfrac{3}{4})/\Gamma(\tfrac{3}{2}))^2(\Gamma(5/4)/\Gamma(\tfrac{3}{4})-1/\Gamma(\tfrac{3}{2}))+\tfrac{3}{4}(\Gamma(5/4)\Gamma(7/4)/\Gamma^2(\tfrac{3}{2}))^2=+0.0076.$$

The negative magnetoresistance character shown in Eq. (18) is comprehensible when one considers the "cooling" effect of the magnetic field. The Lorentz force results in the deflection of carriers from the X direction, and is only cancelled on the average by the Hall field; thus one has an increased resistance in the presence of a magnetic field. For hot electrons, however, one has, superimposed on this normal effect, a reduction of the average electron energy by the magnetic field and consequently an increase of the average relaxation time. It is this latter phenomenon which results in negative magnetoresistance.

In the limit of large magnetic fields ($\omega\tau\gg1$) and no impurity scattering, one obtains the following asymptotic results from Eqs. (16) and (6):

$$S=e^{-\epsilon/kT_e},$$
$$T_e/T=\tfrac{1}{2}\{1+[1+\tfrac{4}{3}(9\pi/32)^2(\mu_aE_x/s)^2]^{1/2}\}. \qquad (19)$$

[11] E. M. Conwell, Phys. Rev. **90**, 769 (1953).

[12] J. Yamashita and M. Watanabe, Progr. Theoret. Phys. (Kyoto) **12**, 443 (1954).

The approximate expressions for large and small electric fields are given below:

$$T_e=T[1+\tfrac{1}{3}(9\pi/32)^2(\mu_a E_x/s)^2],$$
$$\mu_c=\mu_H=(9\pi/32)\mu_a[1-\tfrac{1}{6}(9\pi/32)^2(\mu_a E_x/s)^2],\quad \mu_a E_x\ll s;\quad (20)$$
$$T_e/T=(1/3^{1/2})(9\pi/32)(\mu_a E_x/s),$$
$$\mu_c=\mu_H=3^{1/4}(9\pi/32)^{1/2}(s/\mu_a E_x)^{1/2}\mu_a,\quad \mu_a E_x\gg s.\quad (21)$$

We note again the negative magnetoresistance associated with the hot-electron regime:

$$\mu_c(B\rightarrow\infty)/\mu_c{}^0=(27/64)\pi^{1/2}\Gamma(\tfrac{3}{4})(32/9\pi)^{3/4}=1.0054.$$

Thus for large magnetic fields the distribution function and galvanomagnetic properties saturate, as is the case for Ohmic conductivity with the neglect of Landau quantization.

IV. MANY-VALLEY BAND STRUCTURE

The calculation of the magnetoconductivity tensor for the case of a many-valley band structure such as occurs in n-type germanium or silicon is extremely complicated and we consider only the simple case of silicon with the electric field in the $\bar{1}10$ direction and the magnetic field in the 111 direction. We shall include the possibility of repopulation effects among the valleys in a semiquantitative manner.

We merely sketch here the method of calculation since the general expressions are extremely unwieldy even for the simple case considered. After transforming the ellipsoidal constant-energy surfaces to spheres and introducing effective electric and magnetic fields, one expresses the total current in terms of an arbitrarily oriented electric field and the applied magnetic field. The constraint that current flow only in the direction of the applied electric field then provides the necessary relations between the applied field and the Hall and anisotropy fields. These must again be determined self-consistently as was discussed in Sec. I. In carrying out this last procedure, one expands all averages in powers of ω.

$$\langle\tau^i/(1+\omega^2\tau^2)\rangle=\bar{\tau}_0{}^i[1+a^i\omega+b^i\omega^2\cdots],$$
$$n^i/n^j=(n_0{}^i/n_0{}^j)[1+x^{ij}\omega+y^{ij}\omega^2\cdots],\quad \text{etc.}\quad (22)$$

where the superscript refers to the valley under consideration and the last equation describes the ratio of the number of carriers in the i and j valleys.

The presence of odd powers of ω in these expansions may be understood by considering the manner in which the energy distribution depends on the electric field in the many-valley case. For ellipsoidal constant-energy surfaces the energy distribution depends on the following mass-weighted electric field.[4]

$$(E^2/m)^*=(1/m_T)[(\mathbf{E}\cdot\mathbf{I}_L)^2(1/K-1)+E^2],\quad (23)$$
$$K\equiv m_L/m_T,$$

where $\mathbf{I}_L$ is a unit vector parallel to the major axes of the ellipsoid and m_L and m_T are the longitudinal and transverse effective masses.

Recalling that $\mathbf{E}$ is the sum of the applied field $\mathbf{E}_A$ plus the orthogonal Hall and anisotropy fields $\mathbf{E}'$ we obtain

$$(E^2/m)^*=(1/m_T)\{[(\mathbf{E}_A\cdot\mathbf{I}_L)^2+(\mathbf{E}'\cdot\mathbf{I}_L)^2+2(\mathbf{E}_A\cdot\mathbf{I}_L)\times(\mathbf{E}'\cdot\mathbf{I}_L)][1/K-1]+E_A{}^2+E'^2\}.\quad (24)$$

This quantity will generally contain both even and odd powers of ω due to the first term in the angular bracket of Eq. (24). Thus, it is only in the case of isotropic effective masses that one has a variation of the energy distribution function which is quadratic to lowest order in the magnetic field. Even in the high symmetry direction we are considering these odd terms in ω contribute significantly.

In our configuration the anisotropy field is zero in the absence of a magnetic field and is of the order of only 1–2% of the Hall field in the presence of a magnetic field and will therefore be neglected. Even with the neglect of the anisotropy field there still remain odd terms in the magnetic field in Eq. (24), namely the last term in the inner bracket.

When one computes the magnetoconductivity tensor one obtains the usual quadratic dependence on the magnetic field for small fields. This does not imply that the odd terms discussed above mutually cancel when one computes the galvanomagnetic coefficients; they merely combine in such a way as to result in the usual forms exhibited in Eq. (17).

Intervalley repopulation effects were treated simply by assuming intervalley transition rates proportional to ϵ^n, where $n=\pm\frac{1}{2}$ roughly characterize phonon and impurity-assisted intervalley processes, respectively. Although this is a great oversimplification we include these two cases for the purpose of illustrating the extreme sensitivity of the hot-electron galvanomagnetic properties to intervalley processes.

The results of this calculation are conveniently expressed in the following form:

$$\mu_c=\mu_c{}^0[1+a(\mu_c{}^0B)^2],\quad \mu_H=\mu_H{}^0[1+h(\mu_H{}^0B)^2],$$
$$\rho=\mu_c{}^0/\mu_H{}^0,\quad \mu_H{}^0=[e\bar{\tau}\Gamma(5/4)/m_\perp(2p_0)^{1/4}\Gamma(\tfrac{3}{2})]\mathcal{L},\quad (25)$$

where $\bar{\tau}=C_2/(kT)^{1/2}$ and p_0 is the value of p (Eq. 13) for the 100 valley in the absence of a magnetic field. The constants a, h, ρ, $\mathcal{L}$ are given in Table I.

TABLE I. Variation of Hall and conductivity mobilities for small magnetic fields.

n	h	a	ρ	$\mathcal{L}$
Single spherical band	−0.031	+0.0076	0.943	1
0	−0.073	−0.2176	1.04	0.628
$\frac{1}{2}$	+0.041	−0.3568	0.991	0.651
$-\frac{1}{2}$	−0.157	−0.1247	1.09	0.609

It is important to note the extreme sensitivity of the galvanomagnetic properties to the many-valley structure, not merely with respect to magnitudes but particularly with respect to the signs of the mobility variations. A realistic calculation of the galvanomagnetic properties of hot electrons in many-valley semiconductors would necessitate a rather accurate treatment of intervalley repopulation processes and could furnish considerable insight into the details of these processes.

ACKNOWLEDGMENTS

The author wishes to express his sincere gratitude to Dr. A. Zylberstjen for his stimulating discussions. It is also a pleasure to thank A. Wilson for his invaluable assistance in the numerical analysis.

PHYSICAL REVIEW VOLUME 140, NUMBER 6A 13 DECEMBER 1965

Mean Free Path of Electrons and Magnetomorphic Effects in Small Single Crystals of Gallium. II*

J. F. Cochran

Department of Physics and The Center for Materials Science and Engineering, Massachusetts Institute of Technology, Cambridge, Massachusetts

AND

M. Yaqub

Department of Physics, Ohio State University, Columbus, Ohio

(Received 22 July 1965)

The variation of the electrical resistivity of pure gallium has been investigated as a function of specimen size in oriented single crystals for current flow along the A and B axes. An analysis of the data, based on the free-electron model and assuming diffuse scattering at the boundaries, gives $(\rho_b l_b)_{A\ \mathrm{axis}}=2.30\times10^{-11}\ \Omega\ \mathrm{cm}^2$, and $(\rho_b l_b)_{B\ \mathrm{axis}}=8.17\times10^{-12}\ \Omega\ \mathrm{cm}^2$. The longitudinal magnetoresistance was measured for fields up to 1400 G and was found to be dependent upon the dimensions of the wires. The resistivity of the wires was also found to be a function of measuring current, and this variation is attributed to the perturbation of carrier trajectories by the magnetic field of the current.

I. INTRODUCTION

In a previous paper[1] (hereafter referred to as YC), we have investigated the low-temperature variation of electrical resistivity of 99.9999% pure gallium single crystals as a function of temperature and size for current flow along the C axis. The present paper is an extension of the same work for current flow along the A and B axes. As before, the crystals investigated are in the form of square wires and their dimensions vary from 0.5 mm to 0.1 mm in six steps, which are more or less uniformly spaced as a function of $1/d$, where d is the side of the square. In order to determine the mean free path of the charge carriers for conductivity along these axes, the results are analyzed in exactly the same way as for the C axis. A full discussion of the theory used is given in YC.

II. EXPERIMENTAL DETAILS AND THE PRECISION OF THE MEASUREMENTS

The method of preparing oriented single crystals, the determination of their dimensions in order to calculate the resistivities, the measuring technique employed for the electrical conductivity, and the various corrections applied are discussed in great detail in YC. The crystals were oriented to an accuracy of about 1°. The dimension of the smallest crystals could be ascertained to about 1%. The resistance measurements for the largest B-axis specimens which had the smallest resistance were accurate to about 2% for the lowest temperatures and the smaller specimens could be measured with a much greater precision. Since the greater accuracy of the conductivity measurements for the small crystals was offset by a less accurate determination of their dimensions and vice versa, the result is that for most specimens the accuracy with which the resistivity could be determined is a little better than 1%, except for the two largest specimens in which the errors may be considerably higher. As for the C axis, the resistance of all the crystals was found to be strongly dependent on the measuring current. It was therefore necessary to extrapolate the resistance to zero measuring current. Consequently, an additional error of about 1% could easily be introduced into the resistivities.

* This research was partly supported by the Advanced Research Projects Agency under Contract SD-90, and partly by the National Science Foundation.

[1] M. Yaqub and J. F. Cochran, Phys. Rev. 137, A1182 (1965).