of Terrestrial Magnetism. The protons scattered from a thin carbon target on a silver leaf backing were counted with an argon-filled proportional counter biased to count only those protons near the end of their range. The ratio of the number scattered by carbon to the number scattered by silver determined the ratio of the cross sections except for a constant multiplying factor due to the difference in the number of carbon nuclei and silver nuclei. Since the ratio of cross sections was determined the target thicknesses were not determined, but both targets were thin; and since both targets were bombarded at the same time the geometrical factors cancel. The protons scattered from the two different nuclei can be resolved because of the difference of the recoil energy of the nuclei.

Number vs. range curves were determined at 55°, 90°, 132.5°, and 160° in the laboratory system for bombarding energies of 2.5 and 2.75 Mev and at 90°, 132.5°, and 160° for energies of 2.0 and 2.25 Mev (Fig. 1). The area under a peak minus a fixed area determined by the channel width of the detecting system, which is a function of the effective counter and depth and the bias of the system, is proportional to the intensity. The channel width is defined here as the intercept of the curve of *widths of the proton peaks at half-maximum vs. range in air* at zero range. The slope of this curve is very nearly the slope of the range straggling curve computed for widths at half-maximum.[1] The area under the peaks was measured with a planimeter; several independent measurements were made for each area and the ratio of the average values was used.

The half-widths due to straggling and channel width were assumed to add as the sum of the squares then the ratio of areas was corrected for channel width by multiplying by

$$[(\sigma_{obs}^2-\sigma_{ch}^2/\sigma_{obs}^2)_c^{\frac{1}{4}}(\sigma_{obs}^2/\sigma_{obs}^2-\sigma_{ch}^2)_{Ag}^{\frac{1}{4}}]$$

where σ_{obs} is the observed width at half-maximum and σ_{ch} is the channel width. The ratio of the areas was also transformed to the carbon center-of-mass system, assuming that the silver scattering is classical.

The corrected area ratios vs. angle in the carbon center-of-mass system are shown in Fig. 2. The dotted line is[2] $R(\phi)$ at 2.5 Mev for S-wave scattering, with δ_0 at 125° as determined by Heitler[2] *et al.*, normalized to 0.1 at 60°. The error of ±10 percent consists principally of the statistical counting errors, errors in area measurement, and errors in measuring the widths at half-maximum. Except for the 2.0-Mev curve the curves all show a definite tendency to decrease at the larger angles and this probably
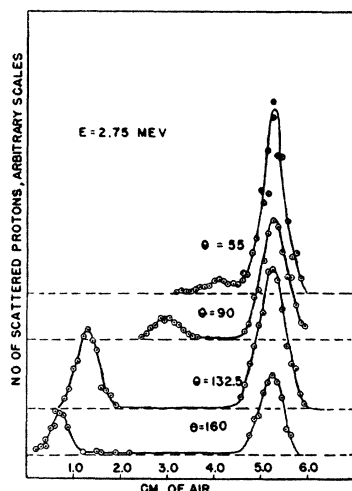

FIG. 1. Family of number vs. range curves for different angles for a thin carbon target on a silver leaf backing at 2.75 Mev. The long-range group is due to silver and the large scale has been shifted so that this group has the same range at all angles. The number scales are different for each angle so that the curves can be plotted on the same graph.
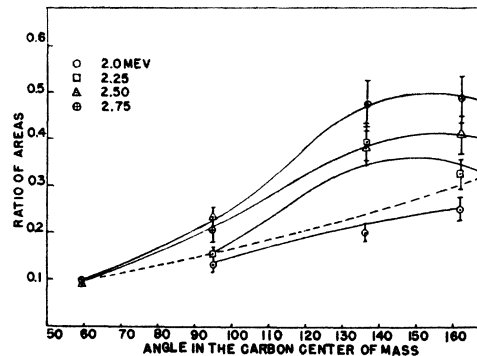

FIG. 2. The ratio S of the areas under the peaks vs. angle in the carbon center-of-mass system. Dotted line is $R(\phi)$ for 80 =125° at 2.5 Mev. See reference 2.

means that the P wave is becoming effective. More precise curves are necessary if the P-wave and S-wave components are to be determined accurately.

Note on the Hall Potential Across an Inhomogeneous Conductor

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IT seems that no formula has been given yet for the Hall potential as a function of the positional coordinates within a conductor of given shape. This can easily be done, however, with the help of potential theory, in the following case (which also may be of interest with respect to inhomogeneous conductors).

Let us consider a thin plate of *homogeneous* conducting material, perpendicular to B, the magnetic induction in the y direction. The relations between the components of E and i are:

$$E_x/\rho=(A_H\cdot B/\rho)i_z+i_x \quad \text{and} \quad E_z/\rho=i_z-(A_H\cdot B/\rho)i_x,$$

where ρ is the resistivity of the material and A_H its Hall coefficient. The y-components of E and i are ignored. We know E to be the negative gradient of a potential function φ, which must obey the Laplace equation $\Delta\varphi=0$. We assume that the solution of this equation can be written as a power series in $A_H B/\rho$, the zero-order term of course being the potential for $B=0$. The first-order term, which we shall call φ_{Hall} is obtained as a solution of the Laplace equation under the following conditions.

We suppose the plate to have the usual *rectangular* form, with current electrodes along two opposite sides ($x=0$ and $x=l$) and Hall probes facing each other somewhere along the other two sides ($z=-\frac{1}{2}b$ and $z=+\frac{1}{2}b$). Furthermore, we suppose the electrodes to be made of fairly well-conducting material, as is generally used, and we ignore the Hall effect in it. Then φ_{Hall} must be approximately zero along $x=0$ and $x=l$, respectively, and therefore the current must have a z-component there. We further assume, however, that in our sample the ratio $A_H B/\rho$ is very small and hence that the current still flows sufficiently homogeneously in the x direction, so that along $z=-\frac{1}{2}b$ and $z=+\frac{1}{2}b$ we have $\partial\varphi_{Hall}/\partial z=A_H JB/bd$, in which J is the total current through the sample, d its thickness in the y direction. $\partial\varphi_{Hall}/\partial z$, abruptly becoming zero at the edges, can be developed in a Fourier series in the usual way.
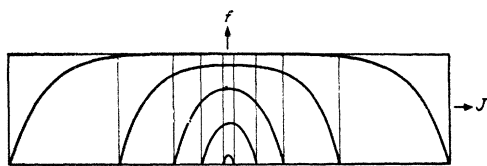
FIG. 1. f-values as a function of the position of the Hall sondes for some different shapes of the sample ("ideal" current electrodes).

The potential function satisfying these conditions is readily found to be

$$\varphi_{\text{Hall}} = A_H \frac{BJ}{d} \cdot \frac{4}{\pi} \sum_1^\infty \frac{1}{(2k-1)} \frac{\sinh[(2k-1)\pi z/l]}{\cosh[(2k-1)\pi b/2l]} \cdot \frac{\sin[(2k-1)\pi x/l]}{[(2k-1)\pi b/l]}.$$

The Hall potential between the probes A and B (on $z = +\frac{1}{2}b$ and $z = -\frac{1}{2}b$, respectively) is:

$$\underset{(A)}{\varphi_{\text{Hall}}} - \underset{(B)}{\varphi_{\text{Hall}}} = V_{\text{Hall}} = A_H(BJ/d)f(b/l, x/b),$$

in which f depends on the ratio of the length to the width of the sample and of the position of the probes along its sides. Some values of f are given in Table I. These results, which include

TABLE I. Values of the function $f(b/l, x/b)$.

x/l \ l/b	0.1	0.5	1.0	2.0	4.0	∞
00	0	0	0	0	0	0
0.125	0.04	0.20	0.39	0.62	0.83	1.00
0.250	0.06	0.30	0.56	0.82	0.98	1.00
0.500	0.075	0.37	0.68	0.94	1.00	1.00

those of Isenberg et al.[1] are plotted in Fig. 1. We determined experimentally the values of f for $x/l = 0.5$, using samples with l/b values of about 0.5, 1, 2, and 8. We found them to be in agreement with the theory. When the problem of the Hall potential is treated in this way, no so-called "einge-prägte Kräfte" nor electrostatic charges are introduced, as was done by Isenberg. It is sufficient to consider the resistivity of the material as a tensor.

Thus no appreciable Hall potential can be built up in the sample if the current electrodes are placed close to each other. Perrier,[2] in suggesting such an arrangement, has not taken into account this fact and therefore his suggestion pointing to very high Hall voltages is incorrect.

This point of view may be of some importance for the study of *inhomogeneous conductors*. Let us consider a material consisting of highly conducting grains separated by thin layers of lower conductivity, and let the picture in Fig. 2 be a consistent model of this situation. The macroscopic resistivity is to a good approximation: $\bar{\rho} = \rho_1 + (l_2/l_1)\rho_2$. The most interesting case occurs if the third term dominates the second. Let this be so. Then approximately $i_1 : i_2 = l_1 : l_2$, where i_1 and i_2 are the *mean* values of the components of the current densities in the direction of J. Now we
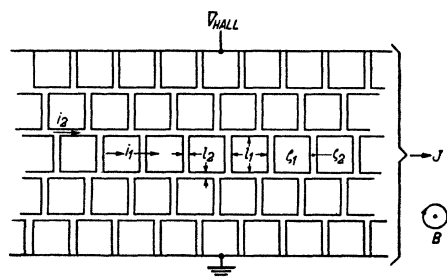

FIG. 2. A model for an inhomogeneous conductor.

have in each grain a Hall voltage of about $A_H^{(1)}l_1i_1B$ and in each longitudinal layer likewise $A_H^{(2)}l_2i_2B$, but in each perpendicular layer the Hall voltage cannot be higher than $0.75(l_2/l_1)A_H^{(2)}l_1i_1B$. We have used here the value of f for the middle of the layer and in the limit for small l_2, assuming that the conducting grains may be considered as electrodes for the layer, as was done above. We suppose $A_H^{(2)} > A_H^{(1)}$.

For the sample as a whole we obtain approximately (with c being of the order of magnitude unity):

$$\bar{A}_H = A_H^{(1)} + c(l_2/l_1)^2 \cdot A_H^{(2)}.$$

Generally $\bar{A}_H$ will not be very different from $A_H^{(1)}$; $\bar{\rho}$ on the other hand can be much greater than ρ_1. This means that, if the Hall effect and the resistivity of such an inhomogeneous conductor is to be explained in terms of electron density and mobility, no conclusion as to the mobility either in the highly conducting grains or in the badly conducting layers may be drawn.

We have in fact found in some samples of semiconducting materials very low values of the mobility, which, however, could be explained in the manner given.

[1] Isenberg, Russell, and Greene, Rev. Sci. Inst. 19, 685 (1948).
[2] A. Perrier, Helv. Phys. Acta 19, 410 (1946).

The Thermodynamics of Liquid Helium of He³—He⁴ Solutions

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THE thermodynamics of liquid helium on the basis of the two-fluid theory of He II has been considered independently by Gorter[1] and the writer.[2] Though using different methods these treatments are in many respects similar. We differ, however, in one rather fundamental point. I assumed zero enthalpy of mixing of normal and superfluid, which allowed me to set the partial molal enthalpy of superfluid[3] $\bar{H}_{4s}$ equal to zero. I then obtained an approximate expression for the partial molal entropy $\bar{S}_{4s}$, which arises from an entropy of mixing of normal and superfluid, to be expected from the picture I developed in which the superfluid is separated from the normal fluid in ordinary space as well as momentum space.

On the other hand, Gorter based all his applications on Tisza's[4] assumption that the total molal entropy is equal to the molal entropy of normal fluid times its mole fraction. This is equivalent to $\bar{S}_{4s} = 0$, or no entropy of mixing. On this basis $\bar{H}_{4s}$ must be finite; in fact, $-\bar{H}_{4s}/T$ will be given by an expression very similar to that found by me for $\bar{S}_{4s}$, and it plays much the same role in the theory, namely that of stabilizing the mixtures of normal and superfluid below the λ-point.

Either assumption is thermodynamically possible. We wish here to consider the possibility of distinguishing between them experimentally by means of the change of the λ-temperature T_λ with pressure or with mole fraction x_3 of He³ in a mixture of the isotopes.

If $\bar{S}_{4s} = 0$, we obtain instead of Eq. (25) of reference 2 (1949)

$$dp/dT_\lambda = S_\lambda/(V_\lambda - \bar{V}_{4s,\lambda}) \tag{1}$$

which differs from Eq. (25) only by lacking the factor $r/(r+1)$ (roughly 0.85) on the right-hand side $[\bar{V}_{4s,\lambda} \equiv \bar{V}_{2,\lambda}$ of Eq. (25)]. This would upset the previously found good agreement with experiment somewhat, but hardly enough to be conclusive in view of the uncertainties involved. It offers, however, a possibility for future investigation.

The effect of He³ on the λ-point of He⁴ has recently been investigated theoretically[5-7] and experimentally.[8,9] Assuming $\bar{H}_{4s} = 0$, and assuming that He³ and normal He⁴ form an ideal solution, with He³ insoluble in superfluid,[10] Engel and Rice showed that

$$dT_\lambda/dx_3 = -RT_\lambda^2/\bar{H}_{4n,\lambda}(1 - x_3) \tag{2}$$