Quantum oscillations of thermoelectric power under magnetic breakdown conditions

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Giant quantum oscillations of the thermoelectric power were observed in single-crystal samples of beryllium, zinc, magnesium, and aluminum in liquid helium under magnetic breakdown conditions when the samples were heated with electric current of ~0.1 A. The magnetic frequency corresponds to the extremal sections of the Fermi surface, and the phase of the oscillations is shifted by one quarter of a period relative to the Shubnikov–de Haas oscillations. The experimental results are compared with the theory of Blatt et al. (Phys. Status Solidi (a) **24**, 621, 1974).

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INTRODUCTION

The behavior of the thermoelectric power in metals at low temperatures in strong magnetic fields has so far attracted little attention. This question was considered in part by Blatt and co-authors,[1] who started with the general formulas of Mott and Jones for the thermoelectric power (see, e.g.,[2])

$$S=eLT\left(\frac{\partial \ln \sigma}{\partial \varepsilon}\right)_{\varepsilon=\varepsilon_F}, \tag{1}$$

where S is the differential thermoelectric power, e is the electron charge, L is the Lorentz number, T is the temperature, σ is the conductivity, and ε_F is the Fermi energy. It was shown that within the framework of the validity of the Wiedemann-Franz law it is possible, for the isotropic case, to transform this formula into

$$S=-eLT\left(\frac{\rho_{xx}'}{\rho_{xx}}\right)_{\varepsilon=\varepsilon_F}, \tag{2}$$

where ρ_{xx} is the diagonal component of the resistivity tensor. The final formula obtained for the change of the thermoelectric power in a magnetic field is

$$\Delta S(H,T)=-eLT\left[\frac{\Delta\rho/\rho}{1+\Delta\rho/\rho}\frac{\partial}{\partial\varepsilon}\ln\frac{\Delta\rho}{\rho}\right]. \tag{3}$$

Here, as usual, $\Delta\rho/\rho\equiv[\rho(H)-\rho(0)]/\rho(0)$. It follows from an analysis of this formula that despite the different possible asymptotic forms of the magnetoresistance in strong fields, the thermoelectric power $S(H)$ tends in any case to saturation. On the other hand, inasmuch as $\rho(H)$ is generally speaking an oscillating function (as a result of the Shubnikov–de Haas effect), it follows that quantum oscillations should be observed also in the dependence of the thermoelectric power on the magnetic field. Such a correlation in the quantum oscillations of the magnetoresistance and thermoelectric power was observed in tin.[3]

It is known that under magnetic-breakdown conditions the oscillations of the magnetoresistance become gigantic (see, e.g., the review[4] as well as[5–7]), owing to the oscillatory H-dependence of the "transparency" of the small orbit to the motion of the electron over an infinite network of magnetic-breakdown trajectories. The conductivity is determined in this case almost entirely by the electrons in the narrow magnetic-breakdown band, and, naturally, no less gigantic oscillations of the thermoelectric power are produced in the magnetic field. Observation of giant quantum oscillations of the thermoelectric power in beryllium were already briefly reported in[8], with an amplitude many times larger than the monotonic components. The present paper is devoted to a somewhat more detailed investigation of the behavior of the thermoelectric power in a magnetic field under magnetic-breakdown conditions. Measurements were made on single-crystal samples of beryllium, zinc, magnesium, and aluminum. It can be stated that the general character of the behavior of the thermoelectric power in a magnetic field is approximately the same for the investigated metals: as a rule, saturation in strong effective fields and oscillations are observed. Principal attention will be paid below to measurements made on beryllium, all the more because of the many published reports of observation of quantum oscillations of the thermoelectric power in aluminum[9] and magnesium.[10]

EXPERIMENT

The measurements were performed in a superconducting solenoid in magnetic fields up to $H\sim 80$ kOe. The investigated samples were cut out by the electric-spark method from already oriented crystallites whose orientation was checked by x-ray diffraction. We used in the experiment samples approximately 5–6 mm long. The sample was placed on a plate made of foil-wrapped fiberglass laminate, on which were mounted spring contacts of silvered beryllium-bronze wire of 0.2 mm diameter. Six contacts were placed over a length ~ 5 mm in an arrangement shown schematically in Fig. 3 below. When current was passed through the outermost contacts 1 and 2, heat was released and produced a small temperature gradient along the sample, which was located in liquid helium, and the contacts 3, 4, and 5 were then at different potentials relative to the remote contact 6; the potential difference was due to the integral thermoelectric power $U(H)=\int S(H,T)dT$. The heating of the region near the contacts was practically unchanged when the resistance of the sample was varied in a magnetic field, since the contact resistances, ~0.1 Ω, were much

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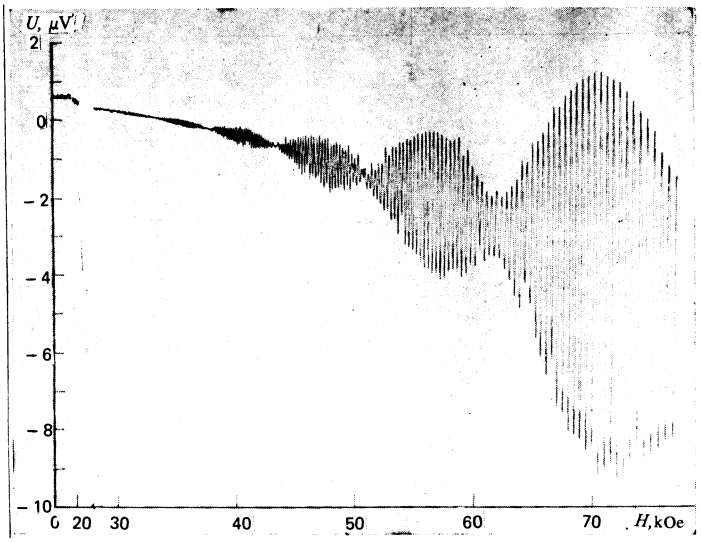

FIG. 1. Plot of the thermoelectric power $U(H)$ for Be-1 at $\mathbf{H} \| \mathbf{C}$ and contact temperatures ~ 8 and 4.2 K. Magnetic frequency $F = 9.4 \times 10^6$ Oe at $H > 60$ kOe and $F = 9.7 \times 10^6$ Oe at $H > 50$ kOe.

larger than the resistance of the sample itself. It can be assumed that the behavior of the observed potential difference in the magnetic field was determined entirely by the thermoelectric power produced in the investigated sample, and the possible thermoelectric power in the spring contact can be neglected. This pertains quite rigorously to the observed quantum oscillations, but the same statement, with slight stipulations, can be made also concerning the behavior of the monotonic part, in view of its dependence on the orientation in the magnetic field.

The contact plate with the sample could be rotated in the magnetic field about two independent and mutually perpendicular axes. The rotation around one axis was by means of a worm gear while the rotation about the other was either with the aid of a device of the "rocker arm" type in a range $\pm 10°$, or by a slight tilt of the principal axis of rotation through $\pm 4°$. The smallest angle-variation interval was approximately $0.2°$. Exact setting of the sample in the magnetic field was monitored mainly by the magnetoresistance effect, transverse or longitudinal, depending on the orientation of the crystallographic axes relative to the sample axes, the current being passed through contacts 1 and 6, and the potential difference measured between contacts 4 and 5. All the measurements were registered with an x-y recorder, the x coordinate being controlled by a voltage proportional to the current in a solenoid, while the y-coordinate was driven by a photoamplified voltage from the sample.

MEASUREMENT RESULTS

As already noted, magnetic breakdown leads to the appearance of giant oscillations of the magnetoresistance of the thermoelectric power in beryllium,[5,8] the magnetic frequency of the oscillations being determined by the central section of the cigar. We consider first the measurement results obtained in exceedingly weak effective fields at the beryllium sample (Be-1) with a resistivity ratio $\rho_{300}/\rho_{4.2} = 11.5$ (the transverse dimensions of the samples were 0.4×0.35 mm, the hexagonal axis of the crystal C was perpendicular to the sample axis; despite the large number of defects, good x-ray patterns were obtained with the single crystal). In this sample, the lifetime of the electrons is such that $\omega\tau \sim 1$ (ω is the cyclotron frequency on the central section of the cigar) only in a magnetic field $H \approx 50$ kOe. Consequently, magnetic breakdown cannot play an essential role in the electric conductivity at the employed magnetic field strengths. Indeed, the magnetoresistance of this sample in the $\mathbf{H} \| \mathbf{C}$ direction increases approximately quadratically with the field, and at $H \gtrsim 40$ kOe one can hardly notice (the relative amplitude is ~ 1%) the quantum oscillations of the resistance (the Shubnikov–de Haas effect), whose amplitude is not appreciable at $\mathbf{H} \| \mathbf{C}$. There is likewise no sharp anisotropy in the dependences of the thermoelectric power U on the magnetic field, a plot of which at $\mathbf{H} \| \mathbf{C}$ and $T = 4.2$ K is shown in Fig. 1. In weak magnetic fields, the monotonic part of the voltage U is positive (we have in mind the potential of the contact 3 relative to the most remote contact 6), and in a magnetic field $H \approx 35$ kOe the plot of U crosses zero and then becomes negative and increases with the magnetic field without exhibiting a tendency to saturation. Superimposed on the voltage oscillations, whose amplitude increases with the magnetic field, is a low-frequency modulation with the large period typical of beryllium, equal to 33 periods of the high frequency. This is evidence that the observed signal consists of a sum of two oscillating signals connected with the central and noncentral sections of the cigar and, just as in the de Haas–van Alphen effect, the contribution of the central section prevails. The magnetic frequency of the oscillations of F at $H < 50$ kOe, measured for the third, fourth, and fifth antinodes of the oscillations (from right to left in Fig. 1) turn out to be 9.7×10^6 Oe, with accuracy 0.5%, corresponding to the noncentral section of the cigar. However, in the section $H > 60$ kOe, which corresponds to the first antinode, the value of F turned out to be, at the same accuracy, equal to 9.4×10^6 Oe, corresponding already to the central section of the cigar. For the "intermediate" second antinode, F was measured with somewhat better accu-

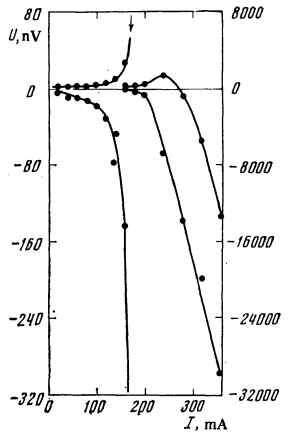

FIG. 2. Measurement of the thermoelectric power $U(H)$ in sample Be-1 with increasing heating. The upper and lower curves are the values of the corresponding maximum and minimum values of U in the magnetic field $H \approx 72$ kOe for different currents; the scale to the left of the arrow is magnified 100 times. Saturation of the oscillation amplitude is observed at $I = 250$ mA.

racy and turned out to be 9.495, 9.496, and 9.503×10^6 Oe at different sections. Such a restructuring of the fundamental frequency, which can be observed also directly in the variation of the depth of modulation, indicates that in strong magnetic fields, at $\omega\tau > 1$, magnetic breakdown comes into play, and it is typical here that the amplitude of the oscillations in these fields becomes larger than the monotonic component, while the voltage U becomes of alternating sign.

The increased heating as a result of the increase of the current through the contacts 1 and 2 leads to a change in the distribution of the temperature over the sample and accordingly to a change in the values of both the monotonic part of the thermoelectric power and of the amplitude of the oscillations. Figure 2 shows the dependence of the variation of the potential at the maximum (upper curve) and the neighboring minimum of the thermoelectric power in a magnetic field $H \approx 72$ kOe on the current. In the left side of the figure the scale is increased 100 times. It is seen that the amplitude of the oscillations first increases, then reaches saturation and remains subsequently unchanged, while the negative value of the monotonic component increases. This indicates that the sample is long enough and the temperature of the cold contact remains unchanged, i.e., $T_6 = 4.2$ K at all values of the current. (Experiments with a short sample[8] have revealed a "collapse" of the oscillations following an increase of the current, owing to the heating of the "cold" contact.) Assume that the temperature dependence of the differential thermoelectric power $S(H, T)$ has, as is usual for the case of Landau oscillations, a factor $\exp(-2\pi^2 kT/\hbar\omega)$ that decreases rapidly with temperature. Then, in the measured integral dependence of the thermoelectric power on the heating of the contact 3

$$U_{36}(H, T) = \int_{T_6}^{T_3} S(H, T)\, dT$$

(T_6 and T_3 are the temperatures of the cold and hot contacts 6 and 3) in the magnetic field $H \approx 70$ kOe, saturation should be observed at a temperature difference ~ 4 K (the ratio of the effective mass of the electron to the free mass for this section is $m^*/m = 0.17$ [11]). It can therefore be concluded that at a current 250–300 mA the heat rise amounts to ~ 4 K and $T_3 \sim 8$ K.

Comparison of the measurements of the thermoelectric power on different contacts has shown that its value decreases abruptly with increasing distance from the current contacts. Figure 3, the upper part of which shows schematically the arrangement of the current and the potential contacts, gives the values of the amplitudes of the voltage oscillations from peak to peak on the contacts 3, 4, and 5 relative to the contact 6 in a magnetic field $H \approx 72$ kOe for two values of the currents, $I = 250$ mA (upper curve) and $I = 140$ mA (lower curve). The voltage on contact 5 was in both cases at the noise level $\sim 2 \times 10^{-9}$ V. Of course, one cannot deduce at all from two points whether the potential drop along the sample is exponential, and the straight lines in the figure are drawn arbitrarily, but these results confirm that the contact 6 can actually be regarded as sufficiently "remote" and that the estimate of the temperature drop in this case is correctly performed.

It is of interest to trace for the qualitative and quantitative variations in the behavior of the thermoelectric power when measurements are made on a purer sample of beryllium (Be-2) with a resistivity ratio $\rho_{300}/\rho_{4.2} = 500$ (transverse dimensions 0.15×0.15 mm, C perpendicular to the sample axis). Despite the large ratio of the length to the transverse dimensions, the increase of the current in this sample in normal helium ($T = 4.2$ K) influences the temperature of the outermost contact 6, this being due to its higher thermal conductivity. In superfluid helium, however, the increase of the heat transfer coefficient already leads to a sufficiently abrupt decrease of the temperature along the sample, and the contact 6 could be regarded as sufficiently remote, just as in the case of sample Be-1 considered above. Therefore the principal measurements on sample Be-2 were made at $T \approx 2$ K. (The holder with the sample was placed in this case in additional cryostat inserted in the solenoid.)

The effective magnetic fields for this sample are much stronger than for Be-1, and magnetic breakdown appears already in relatively weak magnetic fields. This leads, as usual, to the onset of a narrow deep minimum of the magnetoresistance on the rotation diagram at $\mathbf{H} \parallel \mathbf{C}$. The angle oscillations with increasing angle ϑ between H and C, similar to those noted earlier,[5,12] are hardly observable. The dependence of the magneto-

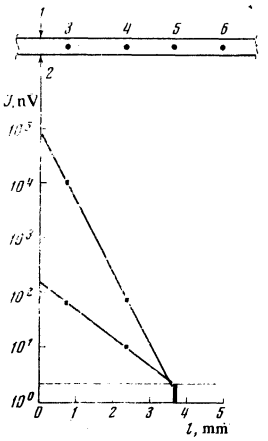

FIG. 3. Arrangement of the contacts (top) and variation of the oscillations of the voltage $U(H)$ over the Be-1 sample for two values of the current: $I = 250$ mA (upper curve) and 140 mA (lower curve) in a magnetic field $H \approx 72$ kOe.

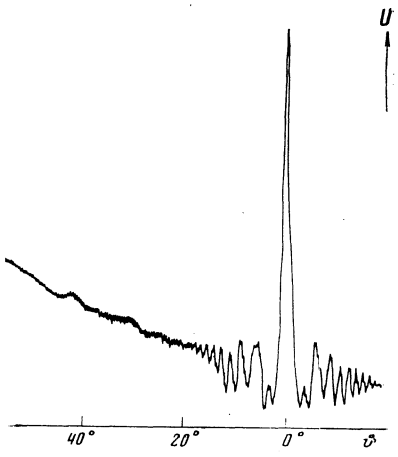

FIG. 4. Rotation diagram of thermoelectric power $U(H)$ in sample Be-2 at $T = 2$ K in a magnetic field $H = 48$ kOe.

resistance at the minimum $\vartheta = 0$ goes to a maximum at $H \approx 35$ kOe and then saturates, with a superposition of giant resistance oscillations whose magnetic frequency is determined by the central section of the cigar (see Fig. 7 below), and whose amplitude decreases sharply with increasing angle ϑ. A strong anisotropy is observed also in the dependence of the thermoelectric power on the angle ϑ, as shown in Fig. 4. The extremum at $\vartheta = 0$ is due to oscillations of the thermoelectric power, and when the magnetic field is shifted by half a period, a similar peak is oppositely directed. Angle oscillations are observed also at $\vartheta > 0$, but their amplitude is already one tenth as large as the amplitude of the central extremum. It must be stated, as already noted earlier,[12] that the relative magnitude of the amplitude of the quantum oscillations in beryllium at $\vartheta > 0$, due to the formation of elongated but nonetheless closed magnetic-breakdown trajectories, should decrease with increasing mean free path. Thus, for the sample Be-1 there is no abrupt change in the amplitude of the oscillations near $\vartheta = 0$; for the sample with the resistivity ration $\rho_{300}/\rho_{4.2} = 100$[8] the amplitude decreases to approximately one-half; for the measured sample (Be-2) it decreases by a factor of 10, only one deep extremum (a maximum or a minimum, depending on the magnetic-field phase) is observed at $\vartheta = 0$ for the sample Be-3 with the resistivity ratio $\rho_{300}/\rho_{4.2} \approx 1000$, and there are no angle oscillations at $\vartheta > 0$.

The dependence of the amplitude of the oscillations of the thermoelectric power and of its monotonic component in sample Be-2 on the current through contacts 1 and 2 is similar to the dependence shown in Fig. 2 for Be-1, but quantitatively the difference is very large. Figure 5 shows a plot of the thermoelectric power against the magnetic field for $\vartheta = 0$ and a current $I = 320$ mA, at which saturation of the oscillation amplitudes takes place with increasing heating, i.e., approximately at the same or somewhat smaller temperature of the hot contact 3 as in Fig. 1 ($T_3 \approx 7$–8 K). The amplitude of the oscillations increases with the magnetic field and reaches a value ~700 μV from peak to peak in a field H ~ 80 kOe. The low-frequency component is due to the Shoenberg effect[13] (see also the pertinent discussion below).

It is clearly seen that the monotonic part of the thermoelectric power is no longer dependent on the magnetic field in strong fields. Figure 6 shows plots of the monotonic part of the signal against the magnetic field for different values of currents through the contacts 1 and 2. The lower curve corresponds to the plot in Fig. 5, i.e., to a current 320 mA at a temperature $T_3 \approx 7$–8 K. The remaining curves correspond to smaller temperature drops, obtained in weaker currents: 220, 180, 100, and 60 mA (reading upward). On the lower left is shown, in magnified scale, a section of the lower curve in weak magnetic fields. Just as in Fig. 1, the signal reverses sign, but here this occurs in a field $H = 1$ kOe. It cannot be stated, however, that the reversal of the sign pertains to the thermoelectric power of the sample, since the observed voltage, as already noted above, corresponds to a difference in the thermoelectric power of the sample and of the measurement contacts, and therefore the true position of the zero of the thermoelectric power remains somewhat indeterminate.

In a magnetic field $H \approx 7$ kOe, all curves have a noticeable tendency to saturation, followed again by an increase and saturation in strong fields, and at $H \gtrsim 50$ kOe the plots are fully saturated, in agreement with the results obtained by Blatt *et al.*[11] It should be noted that for the $\mathbf{H} \perp \mathbf{C}$ direction, when there is no magnetic breakdown, similar $U(H)$ plots were obtained for the sample Be-3, with the saturation setting in already at $H \sim 5$ kOe. This may indicate that magnetic breakdown leads to a substantial change in the limiting value of the thermoelectric power $U(H)_{H\to\infty}$.

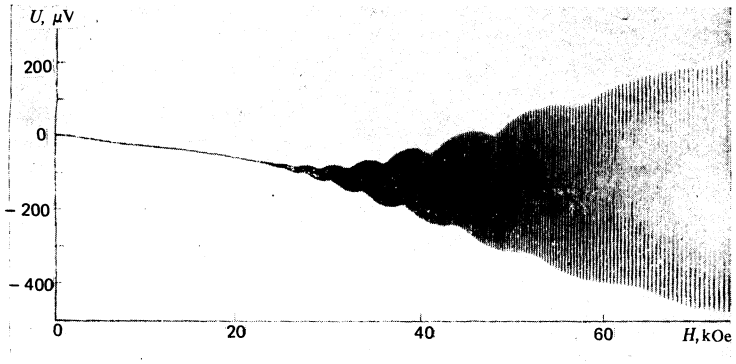

FIG. 5. Plot of the thermoelectric power $U(H)$ for Be-2, the contact temperatures are ~8 and 2 K, and the magnetic frequency is $F = 9.4 \times 10^6$ Oe.

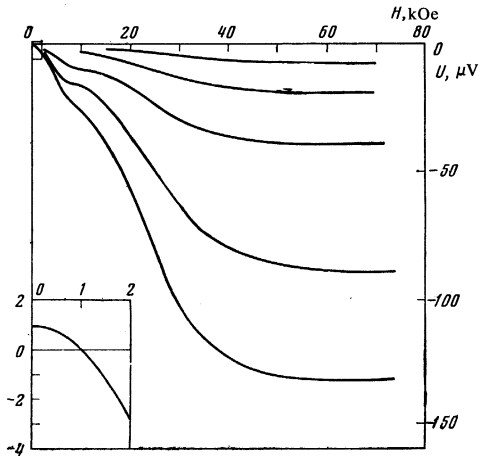

FIG. 6. Plots of the monotonic components of the thermoelectric power against the magnetic field. The lower curve corresponds to the plot of Fig. 5, and the others were obtained at smaller values of the heat rise. The insert in the lower left shows a section of the curve in magnified scale.

As already noted above (see Fig. 5), low-frequency envelopes are observed over the maxima and minima of the oscillations of the thermoelectric power, and are connected with the onset of the Shoenberg effect. The mechanism whereby similar envelopes are produced in the magnetoresistance was considered by Reed and Condon.[6] It makes sense to compare here the dependence of the magnetoresistance of this sample in the **H** ∥ **C** direction (see Fig. 7) at $T \approx 2$ K with the dependence of the thermoelectric power $U(H)$ in the case of a slight temperature rise ~0.1 K, when it can be assumed that the entire sample is also at $T \approx 2$ K and the Shoenberg effect manifests itself most strongly (Fig. 8). It is easily seen that the thermoelectric-power oscillations take an essentially different form and the low-frequency oscillations have different locations than in the oscillations of the magnetoresistance. The Shoenberg effect leads to induction jumps ΔB in the $B(H)$ plot of the sample; these jumps are located at the maxima of the differential susceptibilities dM/dH and are larger the larger the maximum. The oscillatory dependence of the susceptibility of beryllium on the magnetic field in the **H** ∥ **C** direction is amplitude-modulated as a result of the presence of contributions from the central and noncentral extremal sections of the cigar. The contribution from the noncentral section, which is 3% larger, prevails and since the oscillations of the magnetoresistance are determined by the central section, the minima and maxima of $\rho(B)$ coincide with the locations of the jumps ΔB in the antinodes and nodes, respectively, the values of ΔB being smaller at the nodes; the minima are in this case "singled out" more strongly, as it were, and the envelope over the minima has a noticeably larger amplitude than the envelope over the maxima. The oscillations of the thermoelectric power are also determined by the central section of the cigar, but they are shifted in phase by one-quarter period, therefore the extrema of the $U(B)$ plot coincide with the locations of the jumps ΔB approximately halfway between the antinodes and the nodes, and no difference is observed between the amplitudes of the envelopes over the maxima and minima. For the same reason, the positions of the extrema of these envelopes are shifted relative to the magnetic field, likewise by one quarter of a period of the difference frequency. This circumstance provides an addition possibility of determining, for beryllium, the phase of the measured oscillations relative to the oscillations of the susceptibility, since the values of the magnetic field in the antinodes and nodes are fully determined.

The Shoenberg effect leads also to a characteristic distortion of the waveform of the thermoelectric-power oscillations, which becomes noticeably non-sinusoidal. Figure 9 shows a stretched-out plot of the thermoelectric-power oscillations at $T \approx 2$ K. The start of the plot at $H = 46$ kOe corresponds to the location of the jump ΔB at the minima. With increasing magnetic field, the location of the jump relative to the phase of the oscillation changes (the jumps occur more frequently), and the picture repeats after 33 periods. The very same behavior was observed in the variation of the waveform of the magnetoresistance oscillations.

In zinc and magnesium, the situation is qualitatively close to that in beryllium, magnetic breakdown occurring in the basal plane in these metals in weaker magnetic fields. In addition, the amplitude of the oscillations of the magnetoresistance, both in zinc[4,14] and in magnesium,[4] is much less than in beryllium, even for samples of substantially better quality. Measurements of the thermoelectric power were made on a sample of magnesium ($\rho_{300}/\rho_{4.2} \sim 1300$, **C** perpendicular to the sample axis) and two zinc samples, for which $\rho_{300}/\rho_{4.2} \approx 3000$ and **C** is parallel to the axis in one sample (Zn-1) and perpendicular in the other (Zn-2). In all three samples oscillations of the thermoelectric power were observed at **H** ∥ **C**, with a magnetic frequency determined by the central section of the electron ellipsoid in the third zone: the cigar for magnesium ($F = 2.22 \times 10^6$ Oe) and the needle for zinc ($F = 1.5 \times 10^4$ Oe). Owing to the good thermal conductivity, it was impossible to obtain for these samples the dependence of the oscillation amplitude on the heat rise, which exhibited saturation just as in the case of beryllium. The amplitude increased

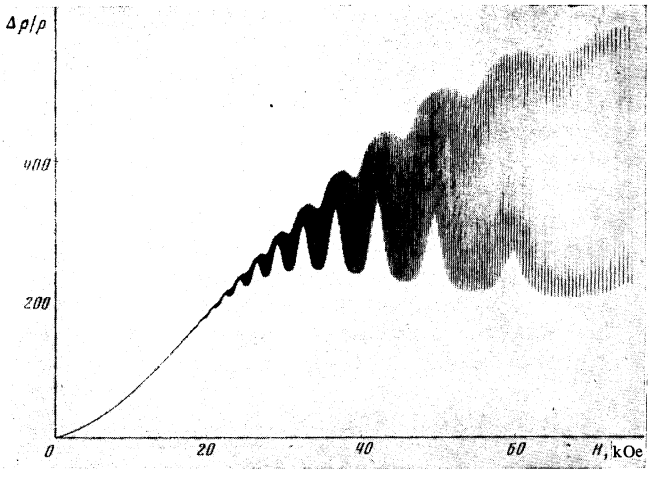

FIG. 7. Magnetoresistance of Be-2 at H ∥ C, $T \approx 2$ K.

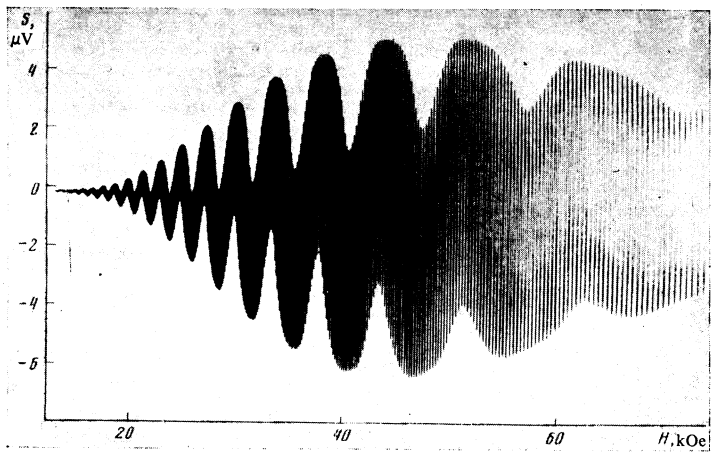

FIG. 8. Plot of the thermoelectric power $S(H)$ for Be-2 at $T \approx 2$ K and at a contact-temperature difference ~ 0.1 K. The low-frequency components appear as a result of the Shoenberg effect[13] (the same as in Fig. 7).

everywhere approximately quadratically with the current and, in addition, no abrupt decrease of the signal was observed likewise with increasing distance from the current contacts, so that the observed signal corresponded to a certain very small temperature drop. Figure 10 shows a plot of the thermoelectric power signal for sample Zn-2 at $\mathbf{H} \parallel \mathbf{C}$. In the analogous plots of Zn-1, when $\mathbf{H}$ is parallel to the sample axis, a much larger relative magnitude of the monotonic components of the signal is observed. This is apparently due to the very large thermal conductivity of the sample along the magnetic field. In both samples, a certain influence of the spin splitting on the shape of the oscillations peaks is observed, as was observed also in the magnetoresistance.[4]

In contrast to hexagonal metals, the magnetic breakdown in aluminum takes place at magnetic-field orientations relative to the crystal axis, a situation previously investigated in detail by measuring the magnetoresistance.[7,15] The thermoelectric-power measurements were carried out here only near the direction $\mathbf{H} \parallel [100]$, when a one-dimensional network of magnetic-breakdown trajectories is produced along [010] and giant oscillations of the magnetoresistance and of the thermoelectric power take place. The faces of the investigated sample were perpendicular to fourfold axes, and therefore the measurements of the sample were made in both a transverse and a longitudinal magnetic field. The results did not differ noticeably in the two cases. Figure 11 shows a plot of $U(H)$ for a magnetic field parallel to the sample axis. The amplitude of the oscillations greatly exceeds the monotonic component, and the magnetic frequency $F = 0.464 \times 10^6$ Oe corresponds to the β section of the sleeve in the third band.[7,15] Just as in the zinc and magnesium samples, no abrupt decrease of the voltage was observed with increasing distance from the current contacts, thus indicating a small temperature gradient. Comparison with the results of the measurements in[9] yields for the plot in Fig. 11 an estimate $\Delta T \sim 0.03$ K for the temperature difference.

DISCUSSION OF RESULTS

Our measurements show that in strong effective fields a saturation of the thermoelectric power $S(H, T)$ is observed as a function of H, in agreement with the results of Blatt *et al.*,[1] the monotonic component being $S(H, T)_{H \to \infty} \gg S(0, T)$, and under conditions of magnetic breakdown this manifests itself particularly strongly. In addition, the magnetic breakdown leads to the onset of giant oscillations of the thermoelectric power, with an amplitude that can exceed by many times the monotonic component. Taking into account the very large value of the amplitude of the thermoelectric power, it should be noted that in measurements, for example, of even a small magnetoresistance, the measurement currents in liquid helium, under conditions of magnetic breakdown, can lead to noticeable thermoelectric effects, despite the very small possible heat rise. On the other hand, it is much easier to observe the quantum oscillations by the described method than, for example, resistance oscillations, and the requirements with respect to the quality of the measured samples are much less stringent. Thus, voltage oscillations of noticeable amplitude (on the order of several microvolts) were observed in a magnesium sample of relatively low quality at $T = 4.2$ K, whereas oscillations of the resistance could be seen[4] only for samples with $\rho_{300}/\rho_T \gtrsim 10^5$. The foregoing seems to pertain generally to quantum oscillations of the thermoelectric power (not necessarily under magnetic-breakdown conditions), as is seen with the results of measurements of "dirty" beryllium as an example (Fig. 1) at $\omega\tau < 1$, where the role of the magnetic breakdown is negligible, and also of the oscillations of the thermoelectric power in beryllium at $\mathbf{H} \perp \mathbf{C}$, which were noted in[8].

In conclusion, notice should be taken of the large quantitative disparity between the observed amplitude of the oscillations of the thermoelectric power, espe-

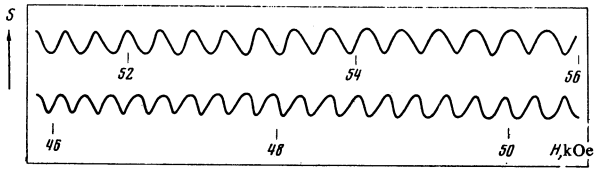

FIG. 9. Influence of the Shoenberg effect on the waveform of the oscillations. Stretched-out part of $S(H)$ at $T \approx 2$ K for Be-2.

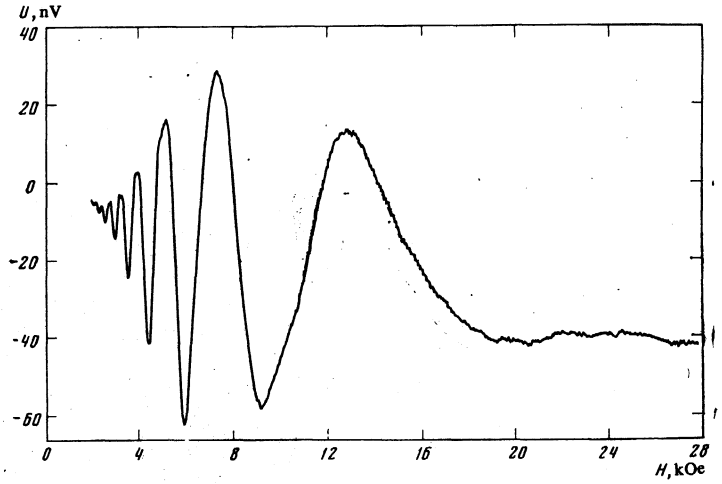

FIG. 10. Plot of the thermoelectric power $U(H)$ at $T = 4.2$ K for the sample Zn-2 at H ‖ C, $F = 1.5 \times 10^4$ Oe.

cially in sample Be-2 (Fig. 5), on the one hand, and formula (3) on the other. According to the plot of the magnetoresistance at $T \approx 2$ K (Fig. 7) in a magnetic field ~ 70 kOe, the relative change $\ln(\Delta\rho/\rho)$ in one half cycle is ~ 1, and the corresponding value is $\Delta\varepsilon = \hbar\omega/2 \approx 2.5 \cdot 10^{-3}$ eV. If we substitute these values in (3), then we obtain for the thermoelectric power $S \sim 20$ μV/K, which is smaller by almost one order of magnitude than that observed in experiment. It is difficult to attribute so large a discrepancy to the lattice contributions to the thermoelectric power. It was already noted earlier[8] that the motion of the electrons in a narrow layer of open magnetic-breakdown trajectories in a strong electric field can lead to the onset of a nonequilibrium situation. The reason is that only a very small group of electrons in the narrow magnetic-breakdown layer participate in the charge transfer and accordingly in the absorption of the energy of the electric field, and isotropization in momentum space, mainly as a result of elastic collisions with impurities, does not change within a time τ the non-equilibrium increment to the energy, which is dissipated within a much longer time. It is obvious that this "superheating" in the narrow layer of momentum space is larger the larger the mean free path and the stronger the electric field. In the measurements described above, the current through the contacts 1 and 2 leads to local heating near the contacts, and heat flow is produced along the sample, which is also transported by a small group of electrons in the narrow magnetic-breakdown layer. In this case the energy levels of the electrons in the magnetic field do not overlap in the entire layer of the open magnetic-breakdown trajectories, because of two circumstances that are typical of beryllium: narrowness of the layer, i.e., smallness of $(\Delta p_z)^2/m$ (Oz is the direction of the magnetic field) compared with $\hbar\omega_H$ all the way to fields $H \sim 5$ kOe, and the fact that the cigar is cylindrical, i.e., all the orbits in the layer have the same period. This leads, in a good crystal with large mean free path and at not too high a probability of magnetic breakdown, to a narrow magnetic breakdown band, the spectral width of which is comparable with kT. Since the entire charge and heat transfer is due to electrons of this band, whose position relative to the Fermi level is determined by the magnetic field, there occur on the Fermi level very large gradients of the conduction in energy ("transparency" of the magnetic-breakdown trajectory) of either sign. Accordingly, an anomalous electric field of the thermoelectric power is produced (on the order of 1 mV/mm for Be-2, corresponding to the electric field of a current on the order of 10^4 A/cm^2). In this case the result of the theory of[1] can apparently no longer be used, for the following reasons:

It is known[2] that from the generalized transport equations

$$I = e^2 K_0 E + \frac{e}{T} K_1 (-\nabla T), \tag{4a}$$

$$U = e K_1 E + \frac{K_2}{T} (-\nabla T) \tag{4b}$$

the electric field at $I = 0$ is

$$E = \frac{1}{eT} \frac{K_1}{K_0} \nabla T = S \nabla T$$

(S is the thermoelectric power) and the heat flux is

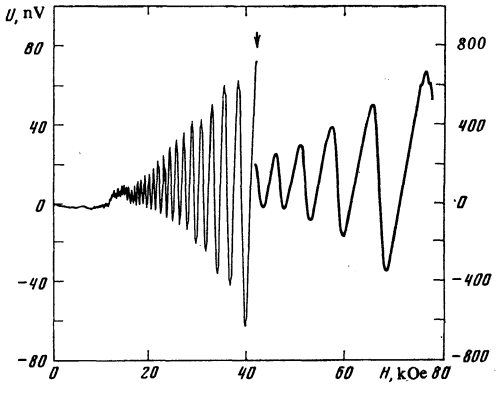

FIG. 11. Plot of the thermoelectric power $U(H)$ in aluminum, magnetic field parallel to the sample axis, $F = 4.64 \times 10^5$ Oe. The scale to the right of the arrow is decreased by a factor of 10.

$$U=\frac{1}{T}\left(K_2-\frac{K_1^2}{K_0}\right)(-\nabla T)=\varkappa(-\nabla T)$$

($\varkappa$ is the thermal conductivity). Generally speaking, K_n are tensors in a magnetic field, but for a compensated metal the problem can be regarded with high accuracy as one-dimensional, i.e., the electric field and the temperature gradient can be assumed directed along the sample. For metals one usually neglects the second term in the parentheses and the Wiedermann–Franz relation is obtained

$$\varkappa=\frac{\pi^2k^2}{3e^2}T\sigma, \quad \text{or} \quad \varkappa=\sigma TL, \tag{5}$$

where the Lorentz constant is $L=2.44\times10^{-8}$ V^2/K^2. In the case of large S, this neglect is incorrect and the substitution leads to the relation

$$\varkappa=\sigma T(L-S^2), \tag{6}$$

i.e., in order for the Wiedermann–Franz law to be satisfied it is necessary to have $S^2\ll L$, which is the case as a rule in metals.

In the present measurements on Be-2, the thermoelectric power was precisely such that $S^2\lesssim L$. Comparison with the earlier measurements[8] and with the measurements on Be-1, where S is much smaller and the discrepancy with formula (3) is smaller, gives ground for expecting to obtain in purer samples an even larger value of S. With increasing S, the determinant of the system (4) tends to zero, meaning that the system approaches an unstable state and the relaxation time increases. In this case formula (3), obtained within the framework of validity of the Wiedemann–Franz law, is no longer valid for a quantitative comparison with experiment, and the situation in the crystal can apparently no longer be described by the distribution of the local-equilibrium temperature. At $S^2=L$ the system (4) degenerates into $U=TSI$ and it can be assumed that either this situation cannot be attained, or in this case the transport processes can be described already by nonlinear equations even at large I and U.

The mutual relation of the nonequilibrium situation with the large value of S can be determined also from the following simple considerations. An estimate of the decrease of the temperature as a result of the Peltier effect yields, when (6) is taken into account, $\Delta T/T\sim S^2/4(L-S^2)$. It is clear that at sufficiently large S the crystal cannot be described by a gradient of a local-equilibrium temperature, and that this formula becomes incorrect.

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[1]F. J. Blatt, C. K. Chiang, and L. Smrcka, Phys. Status Solidi A **24**, 621 (1974).
[2]J. M. Ziman, Principles of the Theory of Solids, Cambridge, 1972.
[3]J. A. Woolam and P. A. Schroeder, Phys. Rev. Lett. **21**, 81 (1968).
[4]R. W. Stark and L. M. Falicov, Progr. Low Temp. Phys. Vol. 5, ed. C. J. Gorter, Amsterdam, 1967, p. 235.
[5]N. E. Alekseevskiĭ and V. S. Egorov, Zh. Eksp. Teor. Fiz. **55**, 1153 (1968) [Sov. Phys. JETP **28**, 601 (1969)].
[6]W. A. Reed and J. H. Condon, Phys. Rev. B **1**, 3504 (1970).
[7]B. J. Balcombe and R. A. Parker, Philos. Mag. **21**, 523 (1970).
[8]V. S. Egorov, Pis'ma Zh. Eksp. Teor. Fiz. **22**, 86 (1975) [JETP Lett. **22**, 38 (1975)].
[9]W. Kesternich and C. Papastaikoudis, Phys. Status Solidi B **64**, N1, K41 (1974).
[10]W. Kesternich and C. Papastaikoudis, J. Low Temp. Phys. **21**, 517 (1975).
[11]B. R. Watts, Proc. R. Soc. London Ser. A **282**, 521 (1964).
[12]V. S. Egorov, Zh. Eksp. Teor. Fiz. **69**, 2231 (1975) [Sov. Phys. JETP **42**, 1135 (1976)].
[13]D. Shoenberg, Proc. Tenth Intern. Conf. on Low Temperature Physics, LT-10, Moscow, 1966.
[14]Yu. P. Gaĭdukov and I. P. Krechetova, Zh. Eksp. Teor. Fiz. **49**, 1411 (1965) [Sov. Phys. JETP **22**, 971 (1966)].
[15]V. N. Morgun, V. I. Khotkevich, L. V. Sidorkina, and G. A. Zaĭtsev, Eighteenth All-Union Congress on Low Temperature Physics (NT-18), Kiev, 1974, p. 1975; V. N. Morgun, Dissertation, Khar'kov State University, 1975.

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