Amplification of quantum oscillations in the polarization parameters of ultrasound as it interacts with a helical wave in tungsten

V. V. Gudkov, I. V. Zhevstovskikh, N. A. Zimbovskaya, and V. I. Okulov

Institute of Metal Physics, Ural Branch of the Academy of Sciences of the USSR
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Quantum oscillations in the polarization ellipticity of ultrasound, in the angle through which the polarization plane is rotated, and in the velocity and absorption of circularly polarized waves have been studied at a frequency of 196 MHz in a tungsten single crystal at 1.8 K in magnetic fields from 30 to 80 kOe. Beyond the edge of the Doppler-shifted cyclotron resonance, the amplitudes of the oscillations in the ellipticity and the rotation angle of the ultrasound vary nonmonotonically with the field. These amplitudes are so large that they cannot be described by the simple expressions for strong fields. The amplification of the oscillations in the polarization parameters is explained on the basis of a theory which incorporates the interaction of ultrasound with a helical wave of the type predicted by Kaner and Skobov. A detailed quantitative comparison reveals good agreement between the theoretical results and the experimental data and reveals numerical values of some new parameters which characterize the Fermi surface, the electron relaxation frequency, and the strain energy.

INTRODUCTION

Some new possibilities for acoustic spectroscopy of metals have recently been pointed out. They are associated with the observation of polarization phenomena accompanying the propagation of transverse ultrasonic waves. Characteristic in this regard are the first results of studies of the quantum oscillations in the polarization ellipticity and the angle through which the polarization plane of ultrasound is rotated in tungsten.[1–4] In particular, these studies have made it possible to observe the interaction of ultrasound with a helical wave in a compensated metal in a strong magnetic field. This wave was predicted in Ref. 5.

Quantum oscillations are of two basic types. The oscillations of one type are characteristic of the static elastic moduli and are analogous to the oscillations in thermodynamic quantities. They were studied in Refs. 6–8. The oscillations of the second type are giant quantum oscillations in the collisionless absorption. These oscillations were predicted in Ref. 9 and were first observed in zinc.[10] They were discussed in detail in a review.[11] Oscillations of the collisional absorption (oscillations of the Shubnikov-de Haas type) are also possible, but they are usually less important. The quantities which are measured in practice sometimes contain oscillatory contributions of both types.

For the transverse ultrasound, propagating in the direction of the magnetic field and a high-symmetry crystallographic axis, in which we are interested here, there are oscillations in the phase velocities and in the absorption coefficients for circularly polarized normal modes. If the amplitudes and phases of the oscillations depend on the polarization of the modes, there are also oscillations in the polarization parameters: the ellipticity and the angle through which the polarization plane of the ultrasound rotates. This effect has been observed in tungsten in the region of the doppleron-phonon resonance.[1] It reflects the role played by the interaction of ultrasound with an electromagnetic wave. In a strong field (beyond the cyclotron-absorption edge) in a compensated metal, on the other hand, we should find a situation in which the sound velocity undergoes oscillations of the first type, while the absorption undergoes oscillations of the second type. The amplitudes and phases of these oscillations, like the monotonic components, should be essentially the same for the two circularly polarized modes. The reason is that the orbits of the conduction electrons are small in comparison with the wavelength, and the electromagnetic field has no circularly polarized normal modes.

In other words, in the strong-field limit the ellipticity and the angle through which the polarization plane is rotated should be vanishingly small for ultrasound. The experimental results of Ref. 2 were accordingly unexpected. They demonstrated the existence of large-amplitude oscillations in the polarization parameters which did not decay with the field up to 80 kOe. Attempts to explain those results on the basis of corrections to the limiting values of the velocities and the absorption in a strong field did not result in a satisfactory agreement, even in order of magnitude. A fundamentally different interpretation[3,4] became necessary. That interpretation also made it possible to distinguish the interaction of the ultrasound with a helical wave, which had not previously been observed.[5] That interpretation yielded a qualitative explanation of the features in the oscillations of the polarization parameters. Specifically, they resulted from amplification of oscillations when a helical wave was excited.

Preliminary results of measurements in a narrow interval of magnetic fields have been reported in some brief communications.[2–4] No quantitative comparison of theory with experiment was made there. Our purposes in the present study were to learn about the quantum oscillations in the polarization parameters of ultrasound over a wide range of magnetic fields (30–80 kOe) and to make a detailed quantitative comparison of the experimental data with a theory incorporating the interaction of ultrasound with a helical wave in a compensated metal. The good agreement which we find lays a solid foundation for the existence of this effect and for determining some new parameters which characterize conduction electrons in tungsten.

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EXPERIMENTAL PROCEDURE AND RESULTS

The acoustic parameters were measured on the ultrasonic apparatus of Ref. 12 at a temperature of 1.8 K. The test sample was made of a tungsten single crystal with a resistance ratio $\rho_{300\,K}/\rho_{4.2\,K} \gtrsim 1.5\cdot 10^5$. It was approximately a cylinder about 10 mm in diameter and 4.57 mm long. Transverse ultrasonic waves with a frequency $\omega/2\pi = 196$ MHz were excited and detected by x-cut lithium niobate piezoelectric transducers. The wave propagation direction and the direction of the magnetic field coincided with the [001] crystallographic axis. A magnetic field was produced by a superconducting solenoid. This field was uniform within $5\cdot 10^{-4}$ in a volume of 1 cm^3.

The method described in Ref. 13 was used to determine the polarization parameters, the absorption, and the velocity of the ultrasonic waves. The amplitude changes $\Delta N_j = \ln[A_j(H)/A_j(0)]$ were measured, where A_j is the amplitude of the signal at the detector. The phases of the signal $\Delta\varphi_j = \varphi_j(H) - \varphi_j(0)$ were measured in the cases of parallel ($j = 1$) and antiparallel ($j = 2$) propagation of the waves with respect to the magnetic field **H**. The sensitive directions of the detecting and exciting transducers were rotated with respect to each other in the plane normal to the wave vector. This rotation was through an angle ψ which was not a multiple of $\pi/2$. Taking account of the excitation of two circularly polarized waves in the metal—these waves have identical amplitudes, identical group velocities, and identical phase velocities $s^\pm$ at $H = 0$—we can write the following relation under the assumption $[s^\pm(H) - s_0]/s_0 \ll 1$:

$$\Delta q^\pm L = -i\ln\{(\exp[\Delta N_2 + i(\Delta\varphi_2 \mp \psi)] - \exp[\Delta N_1 + i(\Delta\varphi_1 \pm \psi)])[2i\sin(\mp\psi)]^{-1}\}. \quad (1)$$

Here $\Delta q^\pm = q^\pm(H) - q$, $q^\pm$ are the complex wave vectors of the circularly polarized ultrasonic waves, $q = \omega/s_0$, s_0 is the phase velocity of the sound at $H = 0$, and L is the acoustic path length. Expression (1) can be used to work from the measured values of ΔN_1, ΔN_2, $\Delta\varphi_1$, and $\Delta\varphi_2$ to calculate the real and imaginary parts of the wave vectors of the modes or the magnetic-field dependent terms in the absorption coefficient $\Delta\Gamma^\pm = \mathrm{Im}\,\Delta q^\pm$ and the velocity $\Delta s^\pm = -(s_0^2/\omega)\,\mathrm{Re}\,\Delta q^\pm$. The ellipticity ε and the angle through which the polarization plane of the ultrasound is rotated, ϕ, are related to $\Delta\Gamma^\pm$ and $\Delta s^\pm$ by

$$\varepsilon = \mathrm{th}[L(\Delta\Gamma^+ - \Delta\Gamma^-)/2], \quad \phi = \omega L[\Delta s^+ - \Delta s^-]/2s_0^2. \quad (2)$$

In the present study we examined the quantum oscillations in the parameters ε and ϕ as a function of the magnetic field between 30 and 80 kOe. The observation of quantum oscillations in weaker fields was reported in Ref. 1. The boundary between the two characteristic field intervals is determined by the position of the doppleron-phonon resonance, which is, according to Ref. 14, the last resonance peak along the H scale in the absorption of ultrasound. Taking this circumstance into account, we can assume that at frequencies $\omega \sim 10^9$ s^{-1} the region of fields which are above the cyclotron-absorption edge begins at about 30 kOe.

Figures 1 and 2 shows the results of the measurements of the quantum oscillations in ε, $\Delta\Gamma^\pm$, ϕ, and $\Delta s^\pm$. According to Ref. 15, the oscillation period $\delta(1/H) = (0.113 \pm 0.002)\cdot 10^{-6}$ Oe^{-1} is evidence that these oscillations stem from charge carriers in the central cross sections of hole ellipsoids (see Ref. 16, for example, for the Fermi surface of tungsten). The centers of the ellipsoids coincide with the centers of the rhombi (points of the N type) which form the Brillouin zone. The axes are oriented along the $N\Gamma$, NP, and NH directions. Figure 3 shows three ellipsoids; the nine others have a corresponding arrangement and have been omitted to keep the figure simple. Ellipsoids with centers at points of the N' and N'' type (there are eight such ellipsoids) contribute to the oscillations. The four other ellipsoids, with centers at points of the N''' type, which lie in the plane perpendicular to **H** and which pass through the center of the Brillouin zone, Γ, do not contribute to the oscillations.

It can be seen from Fig. 3 that when the vector **H** deviates a few degrees from the [001] axis some beats may appear in the oscillations. These beats stem from the appearance of different periods from the ellipsoids at points N' and N'' (Fig. 4a). To improve the precision of the field orienta-

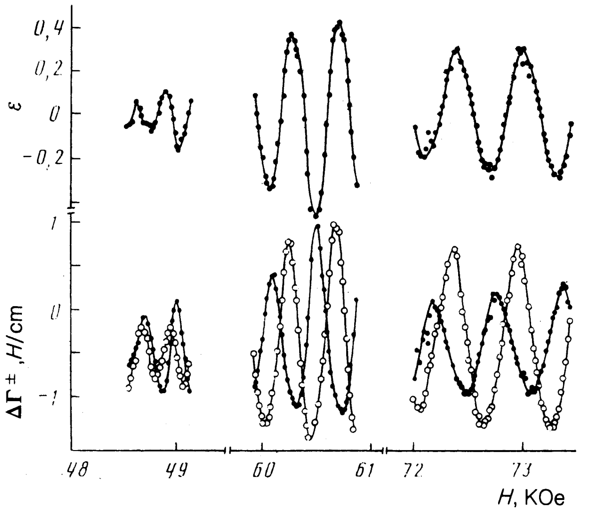

FIG. 1. Absorption of circularly polarized waves and ellipticity of the ultrasound versus the magnetic field. ●— – polarization; ○— + polarization.

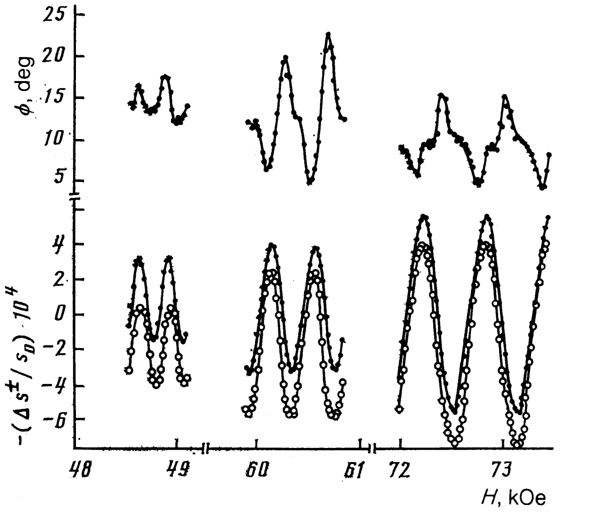

FIG. 2. Magnetic-field dependence of the velocities of circularly polarized waves and of the angle through which the polarization plane of the ultrasound is rotated. ●— – polarization; ○— + polarization.

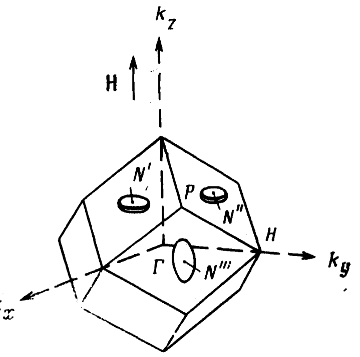

FIG. 3. Positions of the hole N ellipsoids with respect to the center of the Brillouin zone, Γ. The heavy lines on the ellipsoids centered at N' and N'' show the cross sections which cause the quantum oscillations in the ultrasonic parameters in tungsten for the given field configuration.

tion, we used a sample holder which made it possible to vary the angle between **H** and the [001] axis. We concluded that we had achieved the minimum angle for the given experimental conditions when there were no beats in the oscillations on the curve of the amplitude of the received signal versus the field over the entire range of **H** (Fig. 4b).

On the oscillatory curves of the ellipticity and the polarization-plane rotation angle which we found, we first note the large values of these parameters, which are characteristic of a situation in which sound is interacting with electromagnetic waves. A particularly large amplification of ε and ϕ is observed at about 60 kOe. It turns out that the oscillations in the velocities for the waves of different polarizations have different amplitudes, which increase with increasing field. Their phases, in contrast, are identical. In the absorption, the oscillation amplitudes vary nonmonotonically with increasing field. There is a large phase shift between the $\Delta\Gamma^+$ and $\Delta\Gamma^-$ oscillations. To explain this behavior, we take a detailed theoretical look at the situation.

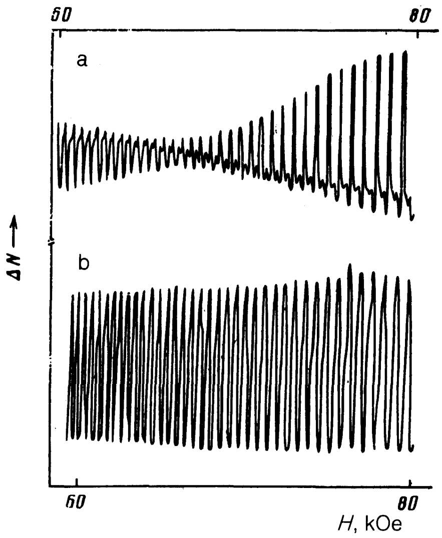

FIG. 4. Magnetic-field dependence of the amplitude of the signal at the piezoelectric transducer detecting the sound. a—Imprecise orientation of the field **H** with respect to the [001] crystallographic axis; b—with **H** parallel to the [001] axis as accurately as possible in these experiments.

THEORY

Using the equations given in Ref. 17, we can construct the following expression for the small increment in the wave vector due to the conduction electrons:

$$\Delta q^\pm = \frac{q}{2\lambda} i\omega\alpha^\pm + \frac{qH^2}{8\pi\lambda}\left[\frac{(1-4\pi i\omega\beta^\pm/cqH)^2}{1-4\pi i\omega\sigma^\pm/c^2q^2} - 1\right], \quad (3)$$

where λ is the elastic modulus; $\alpha^\pm$, $\beta^\pm$, and $\sigma^\pm$ are kinetic coefficients in the equations relating the Fourier transforms of the current density $J^\pm$ and the average strain energy $\overline{\Lambda}^\pm$, on the one hand, to the elastic displacement amplitudes $u^\pm$ and the electric field $E^\pm$, on the other:

$$J^\pm = \sigma^\pm\left(E^\pm \mp \omega\frac{H}{c}u^\pm\right) + \omega q\beta^\pm u^\pm, \quad (4)$$

$$\overline{\Lambda}^\pm = \beta^\pm\left(E^\pm \mp \omega\frac{H}{c}u^\pm\right) + \omega q\alpha^\pm u^\pm. \quad (5)$$

Expressions (3)–(5) have the same form in quantum field theory as in semiclassical theory.[18] There are differences in the specific expressions for the kinetic coefficients. As usual, the quantum effects lead to oscillatory corrections to the semiclassical values under these conditions. These corrections are the topic of interest here.

In strong fields, the characteristic displacements d of the current carriers along the field direction over a cyclotron period are small in comparison with the wavelength of the sound, so the kinetic coefficients can be expanded in series in qd. On the other hand, the mean free path l is long in comparison with the wavelength, so the parameter ql should be regarded as large. We consider oscillatory quantum increments under these conditions, but we will not go through the calculations here. These calculations will be the subject of a separate publication (see also Ref. 3).

Since the quantum oscillations observed in tungsten stem from the hole ellipsoids at the N points in the Brillouin zone, we need calculate only their contributions. The basic oscillatory corrections to the coefficients $\alpha^\pm$ and $\beta^\pm$ are proportional to the values of the off-diagonal components of the strain-energy tensor averaged over the time of the motion in the cyclotron orbit. These average values are nonzero for ellipsoids centered at the points N' and N''. Denoting by Λ_1, Λ_2, and Λ_3 the principal values of the strain tensor at point N, we can write this average value of an off-diagonal component, in the coordinate system of the principal axes of the ellipsoid, as $\Lambda_e = (\Lambda_2 - \Lambda_3)/2$.

The quantum correction in $\beta^\pm$ is also proportional to the average value of an off-diagonal component of the momentum flux tensor, $\Pi_e = \varepsilon_F(m_3 - m_2)/(m_3 + m_2)$, where ε_F is the magnitude of the Fermi energy, reckoned from the maximum in the dispersion law at the N point, and m_2 and m_3 are the principal values of the effective-mass tensor in the coordinate system outlined above. The quantity Π_e is nonzero by virtue of the fluctuations in the z component of the carrier velocity in the cyclotron orbit.

The results calculated for the quantum corrections to $\alpha^\pm$ and $\beta^\pm$ can be written

$$\alpha^\pm_{\text{kv}} = \frac{\lambda_e}{i\omega}\left(K + i\frac{\pi s_0}{2v_F}K_g\right), \quad (6)$$

$$\beta^\pm_{\text{kv}} = \pm\frac{icq}{\omega H}\lambda_e\varkappa\left(K + i\frac{\pi s_0}{2v_F}K_g\right), \quad (7)$$

where $\lambda_e = 2g\Lambda_e^2$, $g = |m| m_{zz} v_F / \pi^2 \hbar^3$ is the Fermi density of states in one ellipsoid, $v_F = (2\varepsilon_F / m_{zz})^{1/2}$, $m = [2m_1 m_2 m_3/(m_2 + m_3)]^{1/2}$ is the cyclotron mass, and $m_{zz} = (m_2 + m_3)/2$ is the longitudinal mass. The functions K and K_g describe oscillations of two types, respectively thermodynamic and giant:

$$K = \left(\frac{\hbar\Omega}{2\varepsilon_F}\right)^{1/2} \sum_{n=1}^{\infty} \frac{(-1)^n n\theta}{n^{1/2}\,\mathrm{sh}\,(n\theta)} \times \cos\left(\pi n \frac{2\varepsilon_F}{\hbar\Omega} - \frac{\pi}{4}\right) \cos\left(\pi n \frac{\Omega_0}{\Omega}\right), \quad (8)$$

$$K_g = 2 \sum_{n=1}^{\infty} \frac{(-1)^n n\theta}{\mathrm{sh}\,(n\theta)} \cos\left(\pi n \frac{2\varepsilon_F}{\hbar\Omega}\right) \cos\left(\pi n \frac{\Omega_0}{\Omega}\right), \quad (9)$$

where $\theta = 2\pi^2 T/\hbar\Omega$, Ω is the cyclotron frequency, T is the temperature in energy units, and $\hbar\Omega_0$ is the spin-splitting energy. Expression (9) is valid under the condition $ql \gg (\varepsilon_F/T)^{1/2}$.

The quantity $\varkappa = \Pi_e/\Lambda_e$ in (7) is small ($\sim 10^{-2}$), because the dimensions of the ellipsoids are relatively small. Accordingly, if we retain in $\Delta q^{\pm}$ only those corrections which are linear in $\varkappa$, we can ignore the quantum increments in the conductivity $\sigma^{\pm}$. The latter are proportional to Π_e^2 and describe the carrier component of the magnetic susceptibility and the giant oscillations in the absorption of energy of the electromagnetic field.

In a description of the oscillations in the polarization parameters, a matter of fundamental importance is the way in which the contribution of the semiclassical conductivity to $\Delta q^{\pm}$ is manifested. In the strong-field limit, the inequality $4\pi i\omega\sigma^{\pm}/c^2q^2 \ll 1$ holds because of the decay of the conductivity with increasing field, and we also have $4\pi i\omega\beta^{\pm}/cqH \ll 1$. The polarization parameters are small in this limit, so only the small correction terms in (3) depend on the polarization. The oscillatory increments contain some additional small amplitude factors, so it is not possible to explain the observed values of the amplitudes of the oscillations in ε and ϕ or the phase shift of the $\Delta\Gamma^+$ and $\Delta\Gamma^-$ oscillations in this case (corresponding quantitative estimates are given in the following section of this paper). We should thus assume that, despite the small value of qd, the quantities $4\pi i\omega\sigma^{\pm}/c^2q^2$ and $4\pi i\omega\beta^{\pm}/cqH$ are not small in that region of the fields and frequencies in which the oscillations are observed experimentally.

Examining the expansion of the conductivity $\sigma^{\pm}$ in powers of qd, we first note that formally the main contribution should be that which describes the collisionless absorption of waves (because the local Hall conductivity of a compensated metal is zero). If that were the case, we would find it difficult to explain the observed polarization effects. However, according to the results of Ref. 14, the collisionless electromagnetic absorption in tungsten is so weak that it is not noticeable against the background of the collisional absorption. The corresponding contribution to the conductivity from the hole ellipsoids, for example, is proportional to Π_e^2. It can thus be assumed that by virtue of certain properties of the Fermi surface of tungsten the collisionless dissipative contribution to the conductivity is unimportant. We can therefore write the asymptotic behavior of the conductivity as

$$\sigma^{\pm} = i\frac{c^2q^2}{4\pi}\left(\mp\frac{1}{\omega_s} - i\tau_s\right), \quad (10)$$

where ω_s and τ_s are proportional to H^3 and $1/H^2q^2$, respectively, with coefficients which depend on the shape of the Fermi surface. The quantity τ_s may have an imaginary part proportional to ω, but under our conditions we need consider only the real part, which contains the electron relaxation frequency. Substitution of expression (10) into the dispersion relation for electromagnetic waves, $1-4\pi i\omega\sigma^{\pm}/c^2q^2 = 0$, shows that ω_s is the eigenfrequency of a wave which is a helical wave in the compensated metal, as predicted in Ref. 5. The interaction of ultrasound with this wave can also explain the observed amplification of the polarization parameters beyond the cyclotron-absorption edge. Also of importance in describing this effect is the coefficient $\beta^{\pm}$. The leading classical term in the expansion of the latter coefficient can be written

$$\beta_{\mathrm{cl}}^{\pm} = \pm\frac{cqH}{4\pi i\omega}\frac{\omega}{\omega_s}\Delta, \quad (11)$$

where the dimensionless parameter Δ depends on the strain energy.

As a result we find the following expressions for the oscillating real and imaginary parts of the quantum increments in the wave vectors of the circularly polarized waves:

$$\mathrm{Re}(\Delta q_{\mathrm{kv}}^{\pm}) = \frac{q}{2\lambda}\lambda_e\left[(1+2\varkappa\Phi_{\pm}')K - 2\varkappa\frac{\pi s_0}{2v_F}\Phi_{\pm}''K_g\right], \quad (12)$$

$$\mathrm{Im}(\Delta q_{\mathrm{kv}}^{\pm}) = \frac{q}{2\lambda}\lambda_e\left[2\varkappa\Phi_{\pm}''K + (1+2\varkappa\Phi_{\pm}')\frac{\pi s_0}{2v_F}K_g\right], \quad (13)$$

where $\Phi_{\pm}'$ and $\Phi_{\pm}''$ are the real and imaginary parts, respectively, of the function

$$\Phi_{\pm} = \frac{1\mp\omega\Delta/\omega_s}{1\mp\omega/\omega_s - i\omega\tau_s}. \quad (14)$$

Expressions (12) and (13) are our starting point for the analysis of the experimental data below.

INTERPRETATION OF THE EXPERIMENTAL RESULTS

We will first give numerical values for the quantities which appear in the amplitude factors in (12) and (13) and which can be found from the existing literature. Among these quantities are the sound velocity $s_0 = 2.88 \cdot 10^5$ cm/s, the elastic modulus $\lambda = 1.5 \cdot 10^{12}$ dyn/cm^2 (Ref. 19), and the parameters of the hole ellipsoids: the strain component $\Lambda_e = (\Lambda_2 - \Lambda_3)/2 = -8 \cdot 10^{-12}$ erg, found with the help of the results of Ref. 20; the effective masses $m_1 = 0.54m_0$, $m_2 = 0.29m_0$, $m_3 = 0.22m_0$ (m_0 is the mass of a free electron); the energy $\varepsilon_F = 4.36 \cdot 10^{-13}$ erg; and the velocity $v_F = 6 \cdot 10^7$ cm/s, calculated from the data of Ref. 21. A value of the g-factor for the holes in the ellipsoids is given in Ref. 8: $g_e = 1.73$. We thus have $\Omega_0/\Omega \approx 0.3$. From the effective masses and the energy ε_F we find the constant $\varkappa = \varepsilon_F(m_3 - m_2)/\Lambda_e(m_3 + m_2) \approx 0.02$.

We should now discuss the shape of the oscillations in the functions K and K_g. Since the observed oscillations $\Delta q^{\pm}$ are approximately sinusoidal, it is clear that the first term is the most important one in the sums over n in (8) and (9).

For this reason, the temperature T, which appears in the amplitude of the harmonics, can be determined quite accurately. A detailed comparison of the observed and calculated oscillatory dependences over a broad field range shows that the value $T = 2.7$ K can be regarded as the optimum value. Along with the first term ($n = 1$) in the sums over n, we consider the correction term with $n = 2$. The difference between this value of T and the experimental temperature (1.8 K) stems from collisions (from the contribution of the Dingle temperature). The corresponding relaxation frequency, $1/\tau_D \sim 10^{11}$ s^{-1}, determines the width of the Landau levels in the hole ellipsoids. Since the relaxation time τ_D refers to a small group of current carriers which we have singled out, it differs slightly in magnitude (but not by more than an order of magnitude) from the transport relaxation time, which determines the resistance of the tungsten.

The next question which arises in our interpretation of the observed oscillations on the basis of expressions (12) and (13) is this: What are the values of the parameters Δ, ω/ω_s, and $\omega\tau_s$, which determine the functions $\Phi_\pm$ under these experimental conditions? We would first like to stress that it is not possible to explain the experimental data under the assumption that the strong-field asymptotic behavior prevails, with small values of ω/ω_s, $\omega\Delta/\omega_s$ and $\omega\tau_s$. Although the complete expression for $\Delta q^\pm$ contains some other small terms in this case—terms which do not come from the expansion of $\Phi_\pm$—this point is irrelevant for order-of-magnitude estimates. It turns out that the polarization parameters calculated under this assumption are smaller by at least an order of magnitude than those observed experimentally. Furthermore, it is not possible to explain the large phase shift of the oscillations in $\Delta\Gamma^+$ and $\Delta\Gamma^-$. We thus conclude that in this sense the quantities ω/ω_s, $\omega\Delta/\omega_s$, and $\omega\tau_s$ are not small under these experimental conditions and that the strong-field asymptotic behavior has not been attained.

Is it worthwhile to consider all three of the parameters in the functions $\Phi_\pm$? Since the polarization dependence is important, it is clear that we do need to consider $\omega\Delta/\omega_s$ or ω/ω_s. Otherwise, the theoretical H dependence of the amplitudes of the polarization parameters will be monotonic, while the observed dependence is definitely not monotonic. The values of ε and ϕ will be small quantities of the same order as in the strong-field approximation. The ratio ω/ω_s thus plays an important role in the functions $\Phi_\pm$, reflecting the interaction of the sound with the helical wave. If the condition $\omega\tau_s \ll 1$ held, the wave would have been slightly damped, and the interaction would have been resonant. The interaction would have occurred in a narrow field interval near the resonant field and would have resulted in a sharp increase in the polarization parameters. Since these effects were not observed, the helical wave was not slightly damped under our conditions; i.e., $\omega\tau_s \sim \omega/\omega_s \sim 1$. We find that the observed values of the amplitudes of the oscillations in ε and ϕ can be explained only if we take account of the quantity $\omega\Delta/\omega_s$ with $\Delta \sim 10$.

To compare the theoretical H dependence of $\Delta q^\pm$ with the experimental dependence in more detail, we single out the dependence of ω_s and τ_s on the magnetic field: $\omega_s = C_1 H^3$, $\tau_s = C_2/H^2$. We use the constants C_1, C_2, and Δ as adjustable parameters, bearing in mind that they have actually already been determined in order of magnitude. It can be seen from Figs. 1 and 2 that the ε and ϕ oscillations reach their maximum amplitudes in the interval 60–61 kOe; as the field is increased, these amplitudes decrease more slowly than when the field is lowered. These results suggest that the resonant field $H_{\rm res}$ has a value of about 65 kOe. From the conditions $H_{\rm res} = (\omega/C_1)^{1/3}$ and $\omega C_2/H^2_{\rm res} = 1$ we find C_1 and C_2, respectively. Varying C_1, C_2, and Δ around the approximate values given above, we find the optimum values: $C_1 = 3.6\cdot 10^{-6}\,{\rm s}^{-1}\cdot{\rm Oe}^{-3}$, $C_2 = 2.7\,{\rm s}\cdot{\rm Oe}^2$, and $\Delta = 12.7$. The oscillatory parts of Re $\Delta q^\pm$ and Im $\Delta q^\pm$ were adjusted simultaneously in three field intervals; the monotonic contributions found from the experimental results were taken into account in each interval.

After the values found for the constants were substituted into (12) and (13), it was found that the discrepancy between the theoretical $\Delta q^\pm(H)$ dependence and the experimental results generally does not exceed 20%. One particular reason for this discrepancy is that the parameters of the ellipsoids which were used in the calculations were determined to within the same error. Another source of discrepancy is the uncertainty in the orientation of the magnetic field along the sound wave vector. A possible uncontrollable field inclination of 2–3° would lead to an error in the determination of the field strength H, which in turn appears in the oscillating functions. The agreement would thus be improved if the scale for the field strength were varied over the range of this error during the plotting of the theoretical curves (Fig. 5).

We might add that the shape of the oscillations observed in the absorption differs from the theoretical shape to

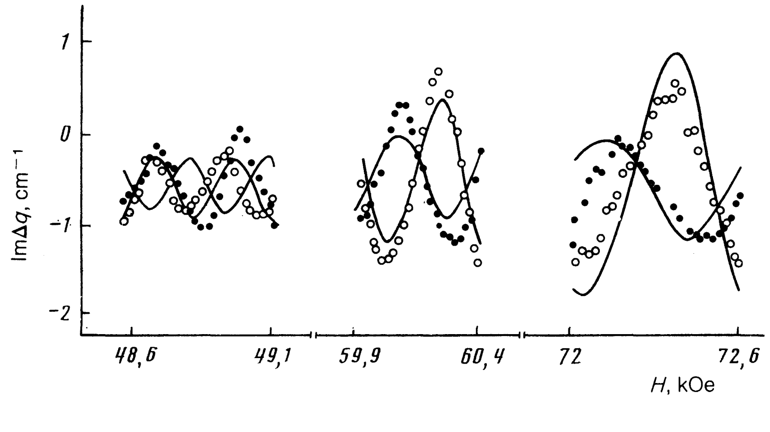

FIG. 5. Magnetic-field dependence of the imaginary part of the wave vector of the ultrasound. ●, ○—experimental data for the − polarization and the + polarization, respectively; solid lines—theoretical predictions from expression (13).

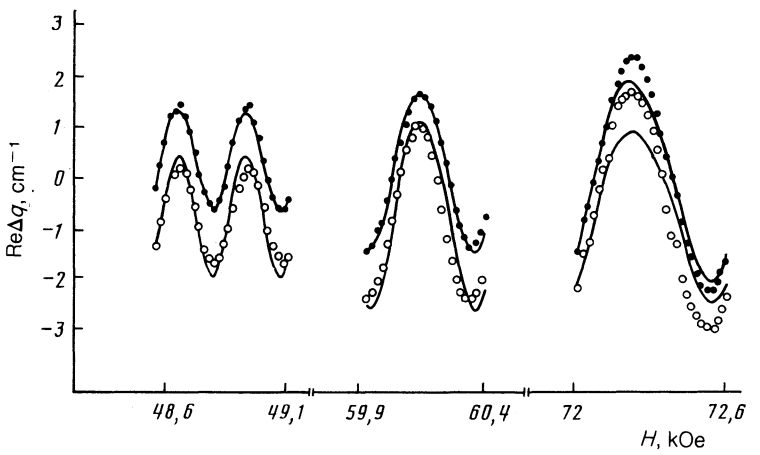

FIG. 6. Magnetic-field dependence of the real part of the wave vector of ultrasound. ●, ○—Experimental data for the − polarization and the + polarization, respectively; solid lines—theoretical predictions from (12).

an extent which is slightly greater than that in the case of the velocities. The probable reason for this result is the approximate nature of the theoretical description of the collisionless absorption. We know[11] that the shape of the giant quantum oscillations at high frequencies may depend strongly on the nature of the electron scattering as well as other factors. In the interval 40–50 kOe, where the discrepancies are greatest, the doppleron-phonon resonance may also be having some effect.

These results, shown in Figs. 5 and 6, demonstrate that expressions (12) and (13), with the optimum parameter values, give a competent description of the quantum oscillations observed over the field range 50–80 kOe. We thus have quantitative support for the existence of an interaction of ultrasound with a helical wave in tungsten and the related amplification of the oscillations in polarization parameters. As a result we were also able to find the frequency of the spiral wave, $\omega_s = C_1 H^3$, the quantity $\tau_s = C_2/H^2$, which is a measure of the damping of this wave at $\omega = 10^9\ s^{-1}$, and the constant Δ, which is proportional to the average value of the strain energy of tungsten. The coefficient C_1 can serve as a characteristic of the Fermi surface, and C_2 as a characteristic of the average collision rate. Expressions for these constants are given in Ref. 3. Manifestations of helical waves of the type studied here may also be seen in other metals, particularly in molybdenum, which is an electronic analog of tungsten.[22]

CONCLUSION

1. Quantum oscillations have been observed experimentally in the polarization ellipticity of ultrasound and in the angle through which the polarization plane is rotated in tungsten over a broad range of magnetic fields (30–80 kOe) beyond the cyclotron-absorption edge. Basic features of the oscillations are large amplitudes and a nonmonotonic dependence on the magnetic field.

2. It has been found that the oscillations in the ellipticity stem from a difference between the absorption amplitudes and phases, while the oscillations in the rotation angle stem from a difference in only the amplitudes of the oscillations in the velocities of the circularly polarized ultrasonic waves.

3. Good quantitative agreement has been found between experimental data and a theory incorporating the action of the ultrasound with a damped helical electromagnetic wave. This agreement supports the existence of such a wave in tungsten and allows us to determine its frequency, its damping, and also the average strain energy.

4. The following conditions must be satisfied in order to observe quantum oscillations in the polarization parameters of ultrasound beyond the cyclotron-absorption edge: The collisionless absorption of electromagnetic waves must be low. The value of the off-diagonal component of the strain-energy tensor, $\Lambda^{\pm} = \Lambda_{xz} \pm i\Lambda_{yz}$, averaged over the cyclotron period, must be nonzero. There must be fluctuations in the longitudinal component of the carrier velocity on the effective orbit.

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