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Zero sound in normal and superconducting molybdenum

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Abstract. We discuss the physical nature of electron sound signals excited in molybdenum by an acoustic wave and propagating at the Fermi velocity. The experimental temperature dependences of the amplitude and the phase velocity of these signals have been studied in the normal and superconducting state. We interpret this effect observed earlier in Ga by Burma *et al* as the excitation of a weakly damped zero-sound wave caused by the Fermi-liquid interaction between charge carriers. A dominating role of the electron–electron collisions in the zero-sound damping in Mo was established, and the corresponding relaxation time was estimated. Theoretical calculations of the expected zero-sound behaviour in a superconductor are in good agreement with the experimental data and enable us to determine the intensity of the Fermi-liquid interaction.

1. Introduction

Recently the first observation of electron sound (i.e. signals having a linear dispersion law and propagating at velocities of the Fermi range order), which has been excited and recorded for ultrapure gallium via piezoelectric transducers responding to the elastic component of a wave, was reported [1]. Qualitative considerations about its Fermi-liquid (zero-sound) nature were expressed in [1] on the basis of some experimental results revealing an analogy with the zero-sound behaviour in superfluid ^{3}He [2]. The problem of acoustic excitation and propagation of a zero-sound wave in metals, especially in the superconducting state, has not been studied theoretically up to now and represents a very interesting topic within the Fermi-liquid approach for studying the electron structure of metals. In this paper we present the results of a detailed experimental investigation of this effect in molybdenum and develop the theory describing zero-sound behaviour in metals including superconductors. The comparison between theoretical and experimental results reveals good agreement and permits us to determine the isotropic part of the Landau correlation function that has almost never been found by any methods known earlier.

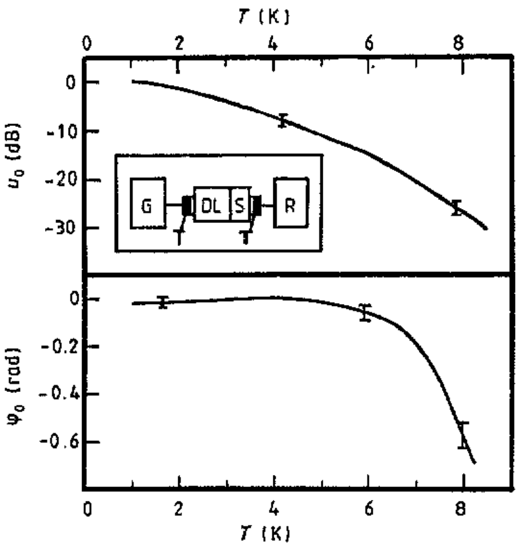

Figure 1. Experimental temperature dependences of the amplitude u_0 and the phase φ_0 of the electron sound in the normal state of Mo ($\omega/2\pi$ = 52.3 MHz; $q \| [111]$; $L = 0.284$ cm). The inset shows the experimental arrangement: G, RF generator; DL, delay line; S, sample; R, receiver; T, piezoelectric transducers.

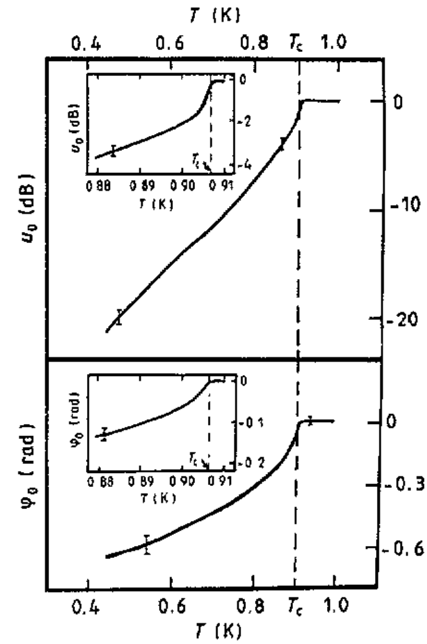

Figure 2. Experimental temperature dependences of the amplitude change δu_0 and the velocity change δv_0^{-1} of the electron sound in Mo below T_c ($\omega/2\pi = 52.3$ MHz; $q \| [111]$; $L =$ 0.284 cm). In the insets their behaviour near T_c is presented.

2. Experimental results

The experimental layout is shown in the inset of figure 1. The piezoelectric transducers with longitudinal vibrations were excited by short (about 5×10^{-7} s) RF pulses with a carrier frequency $\omega/2\pi$ of 50 and 150 MHz. The delay line was used to separate the rapid electron sound signal from the percolating exciting pulse. The Mo sample of [111] orientation and length $L = 0.284$ cm was prepared from a material with a residual resistance rate $R_{300K}/R_{4.2K} \simeq 10^5$ and a momentum relaxation time $\tau_p \simeq 3 \times 10^{-9}$ s, both determined by the sound velocity measurements at liquid-helium temperature in the magnetic field $\boldsymbol{H} \perp \boldsymbol{q}$, where $H = 15$ kOe [3].

As in [1], the receiving system recorded, as well as the sound signal, a rapid weak signal passing through the sample almost without any delay. Contrary to [1], we have not made any direct estimation of its absolute velocity v_0 because of the weakness of this signal in Mo. Nevertheless, the full analogy of its behaviour with that observed earlier in Ga enables us to assume that $v_0 \propto v_F$.

We studied the temperature dependences of the rapid signal amplitude u_0 and phase φ_0 both above and below T_c. The transformation efficiency K, i.e. the relation between u_0 and the exciting signal amplitude u_{ex} at $T = 1.6$ K was found, allowing for the attenuation Γ_{ac} of the output sound signal:

$$u_0/u_{ex} = -(95 \pm 0.5) \text{ dB at } 50 \text{ MHz}$$
$$u_0 u_{ex} = -(101 \pm 1) \text{ dB at } 150 \text{ MHz}. \quad (1)$$

The temperature dependences of u_0 and φ_0 in the normal state of Mo are shown in figure 1. We have succeeded in our measurements up to 8 K; above this temperature the signal cannot be distinguished from the noise. The variation in φ_0 describes the phase velocity decrease with increasing temperature.

In a superconducting state u_0 decreases like the sound attenuation, and the variation in φ_0 corresponds to decreasing phase velocity (figure 2). The recorded critical temperature of 0.907 K was somewhat lower than its accepted value of 0.92 K because the sample overheated as a result of powerful sound pulses.

The Fermi velocities in the [111] direction, which are essential for further discussion, were determined in our experiments by the tilt effect [4]:

$$v_{F1} = (7 \pm 0.3) \times 10^7 \text{ cm s}^{-1}$$
$$v_{F2} = (2.6 \pm 0.2) \times 10^7 \text{ cm s}^{-1}. \quad (2)$$

3. Discussion of the experimental results

To the best of our knowledge the possible mechanisms of electron sound waves are zero sound [5], acoustic plasmons [6] and quasiwaves [7]. We can note that the temperature dependence of signal amplitude can be described for all cases by the relaxation factor

$$u_0(T) \propto \exp[-L/v_0\tau(T)]. \quad (3)$$

Here $\tau(T)$ is the temperature-dependent relaxation time; $v_0 \equiv v_F$ for quasiwaves, and $v_0 \geqslant v_F$ for zero sound. The variations in $u_0(T)$ versus T^2 (and versus T^3 for comparison) are plotted in figure 3. In the whole temperature region the function $\tau^{-1}(T)$ is described well by the quadratic dependence that indicates the dominating role of electron–electron collisions. This is not surprising, because the same dependence in Mo was found in [8] even for the RF size effect, where the small-angle scattering is substantial. In our experiment, the small-angle scattering cannot give any contribution to wave damping because of the long electron sound wavelength, so that normal electron–phonon collisions are not essential at $T < 10$–12 K.

From the data shown in figure 3 we obtained an inverse relaxation length $[v_0\tau(T)]^{-1} = 0.17T^2$, yielding $\tau^{-1}(T) = 1.2 \times 10^7 T^2 \text{ s}^{-1}$ for $v_0 \simeq v_{F1}$. This value is in good agreement with the data for the RF size effect [8], although somewhat lower than in [8]. Note that, if we evaluate the temperature-dependent contribution to τ^{-1} employing the value $v_0 = (v_{F1}v_{F2}/3)^{1/2}$ [6] for the acoustic plasmon, the coefficient in the T^2-dependence of τ^{-1} will be too small.

The quasiwaves, being a single-electron effect, are caused by the ballistic transfer of an elastic perturbation via electrons with the extreme Fermi velocity. In this case it is difficult to expect noticeable phase velocity variations upon a change in τ. At the same time, for $v_0 = 7 \times 10^7 \text{ cm s}^{-1}$ at $T = 1$ K the whole phase change gives a wave velocity decrease of 50% at 8 K. This indicates the wave nature of the recorded signals, because the velocity of a real eigenmode of the electron system evidently depends on the relaxation time.

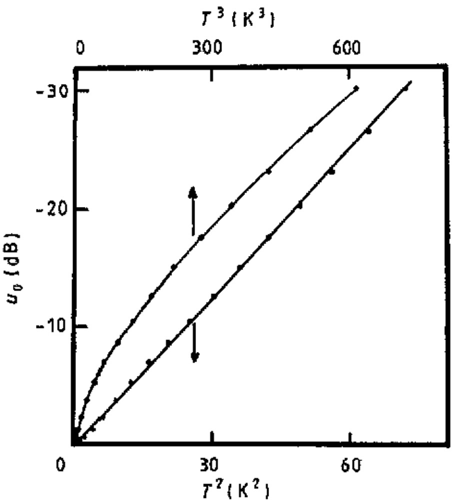

Figure 3. Experimental variations in the zero-sound amplitude u_0 versus T^2 (lower curve) and versus T^3 (upper curve).

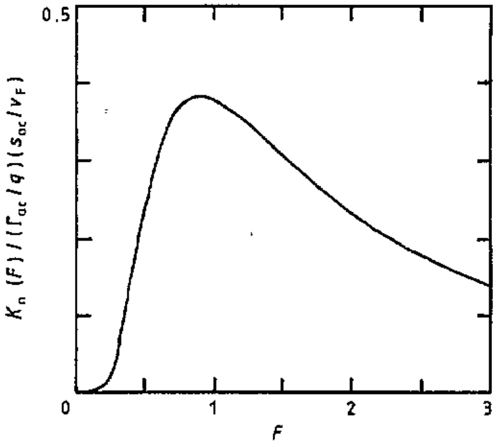

Figure 4. Numerically calculated dependence of the transformation coefficient K_n on the Fermi-liquid parameter F. The magnitude of K_n is normalized by the value $(\Gamma_{ac}/q)(s_{ac}/v_F)$ evaluated by using the same model.

The acoustic plasmon differs from the zero sound first of all in the presence of Landau damping Γ_L:

$$\Gamma_L/\omega \simeq v_{F1}^{-1} = 1.4 \times 10^{-8}\ \mathrm{cm\ s^{-1}}. \tag{4}$$

Assuming that u_0 is given by

$$u_0 = u_{ex} K \exp(-L/v_0\tau) \exp(-\Gamma_L L) \tag{5}$$

we approximate the measured values of u_{ex}/u_0 for different frequencies by a linear function and estimate the possible value $\Gamma_L/\omega \simeq 0.4 \times 10^{-8}\ \mathrm{cm^{-1}\ s}$. The difference between this magnitude and the theoretical value (4) is, however, not sufficiently significant to abandon finally the acoustic plasmon concept.

A more definite conclusion could be drawn from measurements on samples with lengths larger by a factor of 2–3, but unfortunately we could not realize this idea because of signal weakness. However, we found that for a Ga sample with a similar value of L, where the signal was more intensive, the frequency-dependent part of the damping increased much more slowly than L when the latter increased [9]. Most probably the observed frequency dependence of u_0 reflects the diffraction losses, which are essential here because all the typical sizes, such as wavelength, sample length and radiator diameter, are of the same order (Fresnel diffraction region). Thus we can say that the Landau damping is actually absent in the rapid signals; hence they have a zero-sound nature.

4. Theory and its comparison with experiment

To treat our experimental data theoretically within the framework of the Fermi-liquid approach we have used the simple model of the Fermi liquid with two groups of charge carriers possessing spherical Fermi surfaces. This model was employed for the description of zero-sound propagation in normal metals in [10, 11], where it was shown in particular that in a two-band metal with equal carrier masses the dispersion equation for zero sound has the following form:

$$1 + Fa(s) = 0 \qquad a(s) = 1 - (s/2)\ln[(s+1)/(s-1)] \qquad s = \omega/qv_{\mathrm{F}}. \tag{6}$$

Here $F > 0$ is the difference between the intraband and interband Fermi-liquid interaction parameters. Since the zero-sound velocity is unknown in Mo, we cannot determine F directly from the normal-state data by using (6). To extract additional information from the measurements in the superconducting state we must generalize the theory for this case.

The question of the existence of collective modes in a superconductor was formulated immediately after the BCS theory appeared. In a neutral superfluid Fermi system at $T = 0$ the sound-like mode associated with density fluctuations has the dispersion law $\omega = qv_{\mathrm{F}}/\sqrt{3}$ [12, 13]. This result was generalized in [14–16] and applied to ^{3}He with allowance for the Fermi-liquid interaction. In the charged system, where the density fluctuations are suppressed by the Coulomb forces, an analogous mode can exist only in a metal with at least two groups of carriers. This mode, being predicted in [17] for $T = 0$, represents oscillations of the difference $\chi_1 - \chi_2$ between the phases of superfluid condensates, accompanied by fluctuations in their partial densities upon full charge neutrality. It will be shown in this section that in the presence of the Fermi-liquid interaction these oscillations exist under some conditions at all temperatures below T_{c} as the coupled zero-sound vibrations of the quasiparticle gas and superfluid condensate, which can be excited by the acoustic wave.

The superconducting state of a metal with some carrier groups is described by the set of complex ordering parameters $\Delta_i \exp(\mathrm{i}\chi_i)$. Within the BCS model, the small oscillations of χ_i are independent of the fluctuations of Δ_i. Thus the initial set of equations includes the charge neutrality condition

$$\delta n = \sum_{i=1,2} \delta n_i = 0 \tag{7}$$

and the continuity equations, describing the oscillations of charge δn_i and current densities $\boldsymbol{j}_i$ of each group,

$$\delta \dot{n}_i + \operatorname{div} \boldsymbol{j}_i = 0. \tag{8}$$

To calculate the response $\delta n_i, \boldsymbol{j}_i$ of a clean ($\omega\tau \to \infty$) superconductor to slowly varying fields ($\omega \ll \Delta(T)$, $q \ll \Delta(T)/v_{\mathrm{F}}$), we use the quasiparticle representation of the density matrix equations [18, 19], introducing the distribution function $f(\boldsymbol{p}, \boldsymbol{r}, t)$ of quasiparticles with the local energy spectrum $\varepsilon = \varepsilon(\boldsymbol{p}, \boldsymbol{r}, t)$ [20]. Following [18], we allowed for the Fermi-liquid change $\delta\xi(\boldsymbol{p}, \boldsymbol{r}, t)$ of the electron energy $\varepsilon_0(\boldsymbol{p})$ in the spectrum $\varepsilon(\boldsymbol{p}, \boldsymbol{r}, t)$ derived in [21] for the superconductor with deformed lattice:

$$\varepsilon(\boldsymbol{p}, \boldsymbol{r}, t) = \varepsilon_p + \boldsymbol{v} \cdot \boldsymbol{p}_{\mathrm{s}} + (\boldsymbol{p} - m\boldsymbol{v}) \cdot \dot{\boldsymbol{u}} \qquad \varepsilon_p = \sqrt{\xi_p^2 + \Delta^2} \tag{9}$$

$$\xi_p = \varepsilon_0(\boldsymbol{p}) - \varepsilon_{\mathrm{F}} + \Phi + (\lambda_{\alpha\beta} + p_\alpha v_\beta)(\partial u_\alpha/\partial x_\beta) + \delta\xi(\boldsymbol{p}, \boldsymbol{r}, t) \tag{10}$$

$$\delta\xi(\boldsymbol{p}, \boldsymbol{r}, t) = \int \mathrm{d}\tau_{p'}\, f(\boldsymbol{p}, \boldsymbol{p}') \frac{\xi_{p'}}{\varepsilon_{p'}} [f(\boldsymbol{p}, \boldsymbol{r}, t) - n_{\mathrm{F}}(\xi_0(\boldsymbol{p}))]. \tag{11}$$

Here $\Phi = (1/2)(\partial\chi/\partial t) + e\varphi$ is the electrochemical potential, $\boldsymbol{p}_{\mathrm{s}} = (1/2)\nabla\chi - (e/c)\boldsymbol{A}$ is the momentum of the superfluid condensate, both describing oscillations of the phase χ and electromagnetic potentials φ, A; $\lambda_{\alpha\beta}$ is the deformation potential; $f(p, p')$ is the Landau correlation function. For briefness we omit the band indices in (9)–(12).

Considering only the isotropic part of the Landau function $F(p_i, p_k) \equiv \nu_{Fk} f(p_i, p_k) \rightarrow F_0\delta_{ik} + F_1(1 - \delta_{ik})$, where ν_{Fi} is the partial density of states, we linearize the kinetic equation for $f(p, r, t)$ [20] with respect to the small perturbation ψ of the local-equilibrium distribution function $n_F(\varepsilon(p, r, t))$:

$$\partial\psi/\partial t + v \cdot (\xi/\varepsilon)\nabla\psi + (\xi/\varepsilon)(\delta\dot{\xi} + \dot{\Phi} + \Lambda_{\alpha\beta}\dot{u}_{\alpha\beta}) + v \cdot \dot{p}_s = 0. \qquad (12)$$

The solution of (12) determines the system response given by

$$\delta n_i = -e\nu_{Fi} \sum_k (1 + \hat{F})^{-1}_{ik} \left(\Phi - \int d\xi \frac{\xi}{\varepsilon} \frac{\partial n_F}{\partial \varepsilon} \langle \psi \rangle \right)_k \qquad (13)$$

$$j_i = e\nu_{Fi} \left(\tfrac{1}{3}\rho_s v_F^2 p_s + \int d\xi \frac{\partial n_F}{\partial \varepsilon} \langle v\psi \rangle \right)_i \qquad (14)$$

and the Fermi-liquid addition to the excitation energy given by

$$\delta\xi_i = -\sum_k \left(\frac{\hat{F}}{1 + \hat{F}} \right)_{ik} \left(\Phi - \int d\xi \frac{\xi}{\varepsilon} \frac{\partial n_F}{\partial \varepsilon} \langle \psi \rangle \right)_k . \qquad (15)$$

Here $\Lambda_{\alpha\beta} = \lambda_{\alpha\beta} - \langle \lambda_{\alpha\beta} \rangle$, ρ_s is the condensate density, and the angular brackets denote averaging over the Fermi surface. The coupled longitudinal oscillations of the electron and ion subsystems are described by the lattice dynamics equation

$$\omega^2 u = q^2 s_{ac}^2 u + f_e \qquad (16)$$

where s_{ac} is the velocity of sound, f_e is the electron force given by

$$f_e = \frac{1}{\rho} \sum_{i,k} \nu_{Fi} \nabla_q \int d\xi \frac{\xi}{\varepsilon} \frac{\partial n_F}{\partial \varepsilon} \langle \Lambda_i (1 + \hat{F})^{-1}_{ik} \psi_k \rangle \qquad (17)$$

ρ is the metal density, $\Lambda \equiv \Lambda_{qq}$ is the longitudinal component of the deformation potential.

Let us consider first waves in the two-component superconducting electron system, neglecting the electron–lattice coupling and omitting the deformation term in (12). Assuming for simplicity that $\Delta_1 = \Delta_2 = \Delta(T)$, we derive the dispersion equation

$$A(s, F) \equiv 1 + F[a - (a - c)^2/(a + b - 2c)] = 0 \qquad (18)$$

with the linear response coefficients a, b and c defined in [21]:

$$a = 1 - \int d\xi \left(\frac{\xi}{\varepsilon} \right)^2 \frac{\partial n_F}{\partial \varepsilon} \langle \omega R \rangle \qquad (19)$$

$$b = -\frac{\rho_s}{3s^2} - \int d\xi \left(\frac{\varepsilon}{\xi} \right)^2 \frac{\partial n_F}{\partial \varepsilon} (1 + \langle \omega R \rangle) \qquad (20)$$

$$c = -\int d\xi \frac{\partial n_F}{\partial \varepsilon} (1 + \langle \omega R \rangle) \qquad (21)$$

$R = (q \cdot v\xi/\varepsilon - \omega)^{-1}$ is the resolvent of the kinetic equation (12); $F = F_0 - F_1$.

In the normal state we have $a = b = c = a(s)$; so equation (18) coincides with equation (6). In the superconducting state at $T = 0$, where $a = 1$, $c = 0$ and $b = -1/3s^2$, equation (18) describes the propagation of superfluid condensate oscillations at the velocity

$$s(0) = v_F[(1+F)/3]^{1/2}. \tag{22}$$

The condition $v_0(0) > v_F$ providing the weakness of zero-sound damping in the whole temperature region $0 < T < T_c$ has the form $F > 2$ following from (22). In the opposite case the intermediate-temperature range can exist, where Landau damping suppresses zero sound, although this effect could be quite weak because of a rapid decrease in the excitation number.

Expanding (18)–(21) over the small parameter Δ/T_c, we obtain the temperature-dependent correction $\delta s = s(T) - s(T_c)$ describing the zero-sound velocity in the lower vicinity of T_c:

$$\delta s(T)/s(T_c) = -FN_s[(s^2-1)/(1+F-s^2)]f_1(s) \qquad N_s = \pi\Delta/4T_c \tag{23}$$

$$f_1(s) = s/\sqrt{s^2-1} - 2s[s\sin^{-1}(1/s)]^2/[\sqrt{s^2-1} + s^2\sin^{-1}(1/s)] > 0 \tag{24}$$

where $s = s(T_c)$ obeys equation (6).

In studying the excitation of zero sound by the acoustic wave we use the effective-mass approximation for $\Lambda_{\alpha\beta}$:

$$\Lambda_{\alpha\beta} = m^*(\delta_{\alpha\beta}v_F^2/3 - v_\alpha v_\beta) \tag{25}$$

where $m^* = m_e$ for electrons and $m^* = -m_e$ for holes. Within the framework of this model, at $m_1^* = m_2^*$ the coupling between zero sound and elastic deformation is absent ($f_e = 0$), because no macroscopic fields accompany the zero-sound wave in this case. However, in a compensated metal ($m_1^* = -m_2^*$) the electron force (17) has a non-zero value:

$$f_e(q,\omega) = F_e(q,\omega)u(q,\omega) \qquad F_e = (6m_e/M)(qv_F)^2\,B(s,F)/A(s,F) \tag{26}$$

$$B(s,F) = -[s^2 - (1+F)/3]^2/F \tag{27}$$

(M is the ion mass); therefore zero sound can be excited by an acoustic wave on the sample surface. The calculation of the elastic component of zero sound was carried out by means of the well known method of boundary problem solution in the half-space $0 < x < \infty$. Executing the odd continuation of the longitudinal electric field $E(x,\omega)$ and the deformation $u(x,\omega)$ to $x < 0$, we can apply the Fourier transformation to (7)–(17) on the whole x axis, allowing for the jump $2u(0)$ of $u(x,\omega)$ at $x = 0$. Solving the set of inhomogeneous equations arising from this and carrying out the inverse Fourier transformation, we have

$$u(x,\omega) = u(0)\int_{-\infty}^{+\infty}\frac{\mathrm{i}\,\mathrm{d}q}{\pi q}\left(\frac{\omega^2}{\omega^2 - q^2 s_{ac}^2 + F_e(q,\omega)} - 1\right)\exp(\mathrm{i}qx). \tag{28}$$

Here $u(0)$ is the exciting signal amplitude.

The expression under the integral sign in (28) has two poles, one of them being associated with the usual acoustic wave renormalized slightly by the electron force, and the other with the excited zero-sound wave. This reflects the surface acoustic excitation of two eigenmodes of the coupled electron and ion subsystems of a metal which propagate independently with sufficiently different velocities. The second pole contribution to (28) yields the general expression for K:

$$K = 12(m_e/M)(B/s^3)(\partial A/\partial s)^{-1} \quad (29)$$

having the following form in the normal state:

$$K_n = 12(m_e/M)[B(s, F)/(Fs)^2][1 + F - s^2)/(s^2 - 1)]. \quad (30)$$

Equation (30) contains a complex implicit dependence on F, plotted in figure 4 by numerical calculation. It has the maximum at $F \sim 1$ and tends to zero at $F \to 0$ and $F \to \infty$.

In the superconducting state, according to (22), (26) and (27), the function $K(s, F)$ goes to zero at $T \to 0$. Physically this means that the zero-sound non-equilibrium state of the electron system could be excited in the vibrating lattice only in the presence of excitations, disappearing at $T = 0$. In the vicinity of T_c the change δK is proportional to $\Delta(T)$:

$$\delta K(T)/K_n = \delta s(T)[(\partial f_1(s)/\partial s)/f_1(s) + (\partial K_n(s)/\partial s)/K_n(s)]. \quad (31)$$

Passing to the experimental data treatment, on the basis of the theoretical results, we draw attention first to the good qualitative agreement of the measured magnitudes of the transformation coefficient (1) and their behaviour below T_c (figure 1) with the theoretical predictions (30) and (31). In particular, $u_0(T)$ varies near T_c in a similar way to the sound attenuation, being linear with $\Delta(T)$ in accordance with (31).

As regards the temperature dependence of the zero-sound velocity, we emphasize that all these data were obtained by measuring the phase $\varphi(T) = qL = \omega L/v_0(T)$ of the signal which had passed through the sample, so the variation $\delta v_0^{-1}(T)$ is really the measured quantity possessing a sufficiently high accuracy. Thus we can employ in principle the measured phase difference between the normal state and the deep superconducting state:

$$\delta\varphi(0) = \varphi(T \to 0) - \varphi(T_c) = \delta(qL) = (\omega L/v_F)[s^{-1}(0) - s^{-1}(T_c)] \quad (32)$$

to determine F by using the theoretical dependences of $s(0)$ (22) and $s(T_c)$ (6) on F. Unfortunately, in this method there arises the problem of choosing the v_F-value in (32), which is moreover known with a lower accuracy than 5–10% and certainly does not coincide with the model value of v_F.

To avoid this problem, let us use equation (23), which shows that the normalized phase variation $\delta\varphi(T)/\delta\varphi(0) \equiv \delta v_0^{-1}(T)/\delta v_0^{-1}(0)$ near T_c is a linear function of $\Delta(T)/\Delta(0)$ with the slope depending on the only model parameter F. Therefore, by using the numerical calculation of this dependence (figure 5) we can determine F directly by employing the experimentally measured slope. The linear experimental dependence of $\delta v_0^{-1}(T)/\delta v_0^{-1}(0)$ on $\Delta(T)/\Delta(0)$ near T_c, presented in figure 6, confirms our theoretical result (23) and yields the following values of the Fermi-liquid parameter F and the normalized zero-sound velocity $s(T) \equiv v_0(T)/v_F$ obtained from (6) and (22):

$$F = 1.06 \qquad s(T_c) = 1.05 \qquad s(0) = 0.83 \qquad (\boldsymbol{q} \| [111]). \quad (33)$$

Finally, from equations (32) and (33) we found a quite acceptable model value of $v_F = 3.7 \times 10^7$ cm s^{-1} which proved to be of the order of the measured magnitudes (2); this shows the internal self-consistency of the suggested theoretical model.

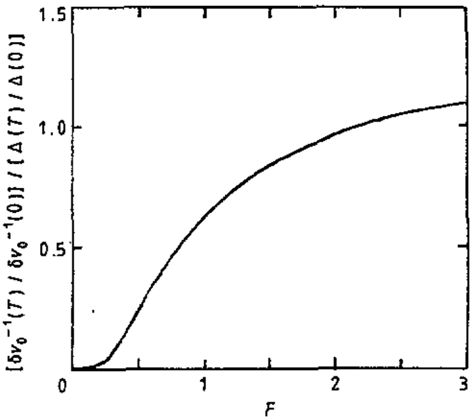

Figure 5. Numerically calculated value of the linear dependence slope of the normalized zero-sound velocity variation $\delta v_0^{-1}(T)/\delta v_{\parallel}^{-1}(0)$ on $\Delta(T)/\Delta(0)$ plotted as a function of F.

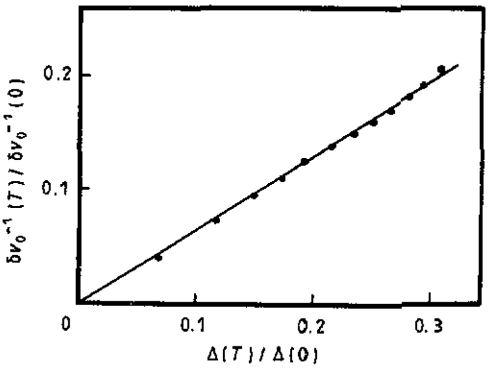

Figure 6. Experimental dependence of $\delta v_0^{-1}(T)/\delta v_0^{-1}(0)$ on $\Delta(T)/\Delta(0)$ (BCS) near T_c.

5. Conclusion

The discovery of zero sound in Mo extends the set of materials in which sound can propagate in normal and superconducting state. We have observed such signals also in magnetic fields up to 15 kOe [22] at any orientation of $\boldsymbol{H}$. The theoretical analysis [22] confirms the Fermi-liquid nature of this wave which has a velocity $v_H < v_0$ independent of H. The remarkable feature of the experiments in [22] was the possibility of the electromagnetic excitation of the zero-sound signal, which becomes comparable with the corresponding acoustic signal. Moreover, our preliminary investigation of Al shows the possibility of zero-sound excitation also in uncompensated metals, which could be explained beyond the simplest approximation of $\Lambda_{\alpha\beta}$ (25). So we consider zero sound both as an observable phenomenon in metal plasmas and as the direct and powerful instrument for Fermi-liquid interaction diagnostics in metals.

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