Figure 2 shows together with curves also the experimental data obtained in [5].

The experimental data were normalized to the point $\eta = 0$.

As seen from the figure, the agreement between the experiment and the curve corresponding to 1500 GeV is very good.

Thus, our analysis shows the following: a) in the intergy interval from 20 to 1500 GeV the inclusive experiment is completely described by the laws governing a phase space whose uniform filling is distorted only by the smallness of the transverse momentum; b) a convenient and sufficiently accurate approximation of these laws is the two-temperature distribution (1) with parameters (2) and (4).

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ANISOTROPY OF SURFACE SUPERCONDUCTIVITY OF LEAD SINGLE CRYSTALS

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Observation of surface superconductivity in type I superconductors is possible only if $\kappa \geq 0.418$ [1]. The only pure metal satisfying this requirement is lead. The occurrence of a superconducting surface layer in lead was confirmed, for example, in [2 - 4]. However, all the experiments performed to date were on polycrystalline samples with large impurity contents, or on films. Mention of measurements on a single crystal is contained in [4], but there are no data whatever on the anisotropy of the effect, and even the crystallographic orientation of the sample is not indicated.

We report here measurements of the surface resistance of lead single crystals whose superconductivity has been destroyed by a magnetic field. The electron mean free path ℓ, estimated from the cyclotron-resonance line width [5], amounts to 0.1 - 0.3 mm at T = 1.5°K and 0.05 - 0.1 mm at T = 4.2°K. Thus, the case of an extremely pure type-I superconductor ($\ell \gg \xi_0$) is realized in these experiments, and a large influence of non-local effects is to be expected.

The samples, in the form of disks 17.8 mm in diameter and 1 and 0.2 mm thick, were grown in a dismountable polished quartz mold [6] and placed in a strip resonator tuned to 9.2 - 9.6 GHz. The sample surface was not subjected to any additional treatment.

The change of the surface resistance of the sample in the magnetic field (Fig. 1) was revealed by the change in a klystron signal passing through a

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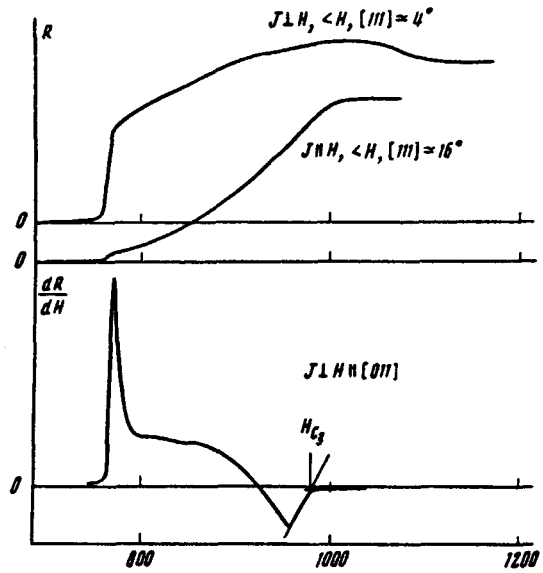

Fig. 1

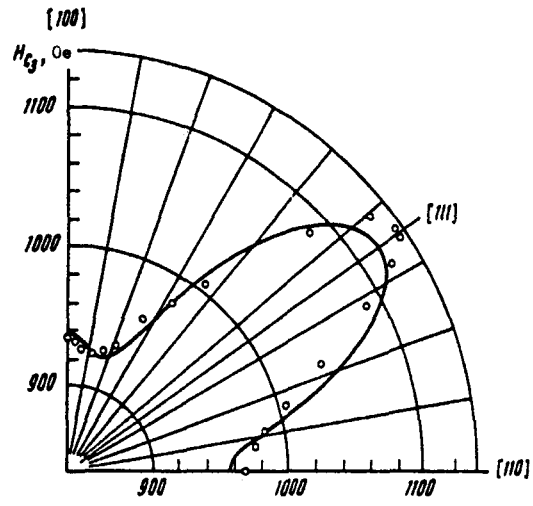

Fig. 2

Fig. 1. Change of the surface impedance of single-crystal lead in a magnetic field for three directions in the (011) plane and two polarizations of the microwave current, $\vec{J} \perp \vec{H}$ and $\vec{J} \parallel \vec{H}$. Temperature T = 1.5°K, sample thickness 0.2 mm.

Fig. 2. Anisotropy of critical field of the surface conductivity H_{c_3} in the (011) plane at T = 1.5°K. Points - experimental values, curve - approximation by the series $H_{c_3}(\alpha, \beta, \gamma) = 975[1 + 0.23K_4(\alpha, \beta, \gamma) + 6.49K_6(\alpha, \beta, \gamma) + 1.78K_8(\alpha, \beta, \gamma)]$, where K_{2n} are cubic harmonics [7], $\alpha = \cos\phi$, $\beta = \gamma = (\sqrt{2}/2)\sin\phi$, and ϕ is the angle between $\vec{H}$ and [100].

resonator weakly coupled to the external lines. The use of large-amplitude frequency modulation of the klystron in conjunction with peak detection of the signal has made it possible to register only the change in the coefficient of transmission to the resonator due to the change in Q, but independent of the resonator frequency shift.

The magnetic field was produced with an electromagnet and oriented parallel to the sample surface with accuracy 5 - 10'. A magnetic-field measurement accuracy of ∿0.1% was ensured with a NMR running-water magnetometer. In the measurement of the surface-conductivity critical field H_{c_3}, the magnetic field was modulated at a frequency 12 Hz and amplitude ∿2 - 6 Oe, so that the registered signal (Fig. 1, lower curve) was proportional to dR/dH. Figure 1 shows the graphic procedure used to determine H_{c_3}. The reproducibility of this quantity for different samples with slightly-oxidized surfaces at the same value of T and identical field orientation relative to the crystallographic axes was 0.3 - 0.5%.

The anisotropy of the field H_{c_3} was investigated in detail in the (011) plane, and the results are shown in Fig. 2. Since $H_{c_3} = C\kappa\sqrt{2}H_c = CH_{c_2}$, it is natural to relate the anisotropy of H_{c_3} to the anisotropy of H_{c_2}, which in turn is shown in [7] to be connected, for metals of cubic symmetry, with the nonlocality of the interaction that determines the superconducting-ordering parameter, and with the anisotropy of the Fermi surface. By retaining three angle-dependent terms in the series expansion of cubic harmonics [7] we obtain the

angular dependence shown by the continuous line in Fig. 2 (the points are experimental). The coefficients of the first and second cubic harmonics in the expansion depend on the relative temperature $t = T/T_c$ like $1 - t$ and $(1 - t)^2$, respectively, in qualitative agreement with the theory of [7]. It should be noted that the anisotropy of H_{c_3} in lead, observed in the present study, is more than double the anisotropy of H_{c_2} in pure niobium [7]. For a complete description of the observed anisotropy, as seen from Fig. 2, it is necessary to take into account terms of higher order in the expansion in cubic harmonics.

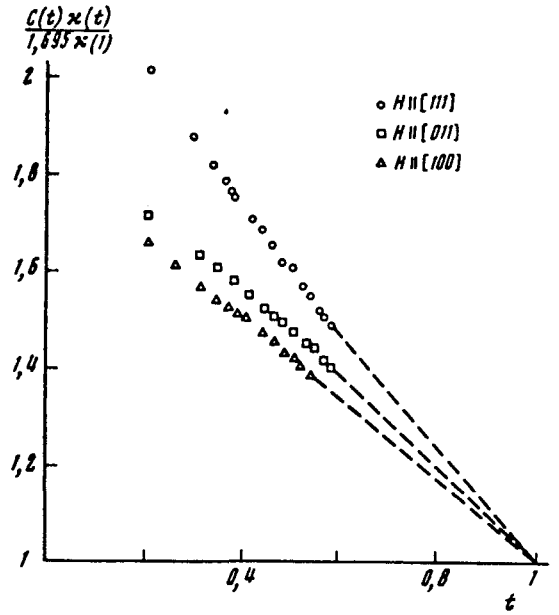

Fig. 3. Temperature dependence of the parameter $C(t)\kappa(t)/C(1)\kappa(1)$ for three rational directions in the (011) plane. Points - experimental values, dashed lines - extrapolation in accordance with the $1 - t$ law. ($C(1) = 1.695$, $\kappa(1) \approx 0.306 \pm 0.01$).

Figure 3 shows the temperature dependence of $H_{c_3}(t)/\sqrt{2}H_c(t) = C(t)\kappa(t)$, normalized to its values at $t = 1$, for three rational directions in the (011) plane. The value of $C(1)\kappa(1)$ is assumed to be 0.52, corresponding to $\kappa(1) \approx 0.306 \pm 0.01$ at $C = 1.695$ [1]. Fisher [3] obtained for lead films $\kappa(1) = 0.328$ and his value of $H_{c_3}(t)$ also exceeds by 7 - 10% our values of the critical field $\overline{H_{c_3}}(t)$ averaged over all the directions in the (011) plane ($\overline{H_{c_3}}(t)$ is the zeroth term of the expansion of $H_{c_3}(t)$ in cubic harmonics).

A characteristic feature, demonstrated in Fig. 1, is the essentially different behavior of R(H) in the dependence on the polarization of the microwave currents. When $\vec{J} \perp \vec{H}$, a maximum is observed on the R(H) plot, with $(R_{max} - R_n)/R_n \approx 0.1 - 0.3$, where R_n is the surface resistance in the normal state. According to the theory [8], for contaminated type-II superconductors, the increase of the surface resistance at $\vec{J} \perp \vec{H}$ in comparison with $\vec{J} \parallel \vec{H}$ is connected with the collective fluctuations of the ordering parameter in the superconducting layer, which in the case of the vortical state are none other than the vortex oscillations. It is shown in [9] that even at a small angle of inclination of the magnetic field to the surface of the metal, a homogeneous superconducting surface layer can go over into the vortical state. It is clear that when microwave currents flow transversely to the vortex structure, the vortex oscillations lead to additional absorption in comparison with the case $\vec{J} \parallel \vec{H}$. The question of the calculation of this absorption for pure lead remains open, however. Fisher and Klein [10] also observed an R(H) dependence with a maximum, but judging from the figure in their paper, the maximum was not repeated when the temperature was lowered. It must be emphasized that in our experiments the absorption maximum was observed for all seven investigated single crystals at $1.5°K \leq T \leq 4.2°K$ provided only $H_{c_3}(T) > H_c(T)$ and at a microwave power level variation of 40 dB (maximum power ~1 mW).

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SPIN-ECHO ENHANCEMENT BY DOUBLE-FREQUENCY PUMPING

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We propose the following explanation for the experimentally observed enhancement of the echo in ferrites [1]. Owing to the internal nonlinear interactions, the oscillations of the magnetic moment, which are coherent with the first weak pulse (signal), become parametrically amplified by oscillations produced by the second powerful pulse (pump). As a result, when the time interval between the exciting pulses (τ) is increased, the echo signal first increases to a certain value, and then decreases as a result of the relaxation processes.

The propcsed mechanism makes it possible to predict qualitatively a new phenomenon in spin echo, namely, enhancement when pumped at double frequency.

For a quantitative description of the enhanced echo we write down the equations of motion of the oscillation amplitudes of the ν-th mode (a_ν), which interact parametrically with the oscillations produced by the second pulse

$$\dot{a}_\nu - i\omega_\nu a_\nu + i\alpha_\nu \exp[2i\omega_\nu(t-\tau)]\,a_\nu^* = 0. \qquad (1)$$

Here ω_ν is the complex natural frequency of the oscillations of the ν-th mode ($\omega_\nu = \omega_\nu' + i\omega_\nu''$), α_ν is a quantity that depends on the magnetic moment produced by the second pulse and the system parameters.

Solving Eq. (1) with the initial condition $a_\nu = a_\nu^I \exp(i\omega_\nu\tau)$ at $t = \tau$, where a_ν^I is the amplitude of the oscillations produced by the first pulse, we can show that an enhanced echo signal becomes phased-in at $t = 2\tau$. The expression for the gain is

$$K = \frac{1}{4}\left\{\exp\left[\frac{|\alpha|}{2\omega''}(1-\exp[-2\omega''\tau])\right] - \exp\left[\frac{|\alpha|}{2\omega''}(\exp[-2\omega''\tau]-1)\right]\right\}\exp(-4\omega''\tau). \qquad (2)$$