Anomalously deep penetration of a magnetic field into metal-oxide superconductors

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The anomalous temperature dependence and the large magnitude of the London penetration depth in the new high-temperature superconductors are described satisfactorily by the model of a heavy charged Bose gas of bipolarons.

The published measurements of the depth to which a magnetic field penetrates into the high-temperature superconductors solidly support the assertion that the behavior of this penetration depth is anomalous in comparison with that of ordinary superconductors. The absolute values of λ_H at liquid-helium temperatures reach 2000–3000 Å for the lanthanum ceramics[1,2] and 6000 Å for yttrium superconductors.[3,4] Measurements of the temperature dependence $\lambda_H(T)$, including that at low temperatures,[4] also provide evidence of a qualitative difference from the prediction of the BCS theory.

In the present letter we offer an explanation for these results based on the theory of the bipolaron superconductivity of high-temperature superconductors which was developed in Refs. 5 and 6.

As was shown in Ref. 7, a bipolaron superconductivity, which occurs in the case of a strong electron-phonon coupling in the crystal (the electron-phonon coupling constant must satisfy $\lambda \gtrsim 1$), is a superfluidity of a charged, nonideal Bose gas with a short-range interaction potential between particles. The critical temperature of the bipolaron superconductivity and λ_H are given by the obvious relations

$$T_c \approx 3.31 \frac{n^{2/3}}{m^{**}}, \qquad \lambda_H(0) = \sqrt{\frac{c^2 m^{**}}{16\pi e^2 n}}, \tag{1}$$

where n is the density of the Bose gas, and m^{**} is the mass of a bipolaron, which describes a two-particle tunneling in a narrow band. Here and below, we are setting $\hbar = k_B = 1$. Substituting the experimental values $\lambda_H(0) = 1500$–6000 Å into expressions (1), we find the following values for the mass and density of the particles in the case of an yttrium ceramic ($T_c \approx 90$ K):

$$n = (10^{24} - 2.7\times10^{20})\ \mathrm{cm}^{-3}, \qquad m^{**} = (3200 - 14)\, m_e.$$

Noting the Hall-effect measurements ($n \approx 10^{21}$–10^{22} cm^{-3}), we find the value $m^{**} = (50$–$200)m_e$ to be the most plausible; this figure corresponds to a depth $\lambda_H = 2500$–4000 Å.

The large mass is evidence of a pronounced polaron narrowing of the electron band. The temperature dependence can be calculated in the standard way with the help of temperature Green's functions.[8] In a first approximation in the gas parameter $\eta = an^{1/3}$ (a is the scattering length for the scattering of particles by each other), which we will use as a model description, we easily find $\lambda_H(T)$:

$$\left[\frac{\lambda_H(0)}{\lambda_H(T)}\right]^2 = 1 + \frac{1}{3nm^{**}} \int \frac{d^3\mathbf{p}}{(2\pi)^3} p^2 \frac{\partial f}{\partial \epsilon} , \tag{2}$$

where $f(\epsilon) = [\exp(\epsilon(\mathbf{p})/T) - 1]^{-1}$. In the same approximation, the quasiparticle spectrum in (2) is given by the known expression

$$\epsilon(\mathbf{p}) = \sqrt{\frac{4\pi a n_0}{m^{**2}} p^2 + \left(\frac{p^2}{2m^{**}}\right)^2} , \tag{3}$$

where $n_0 = n(1 - t^{3/2})$ is the particle density in the condensate, and $t = T/T_c$.

Substitution of (3) into (2) yields

$$\left[\frac{\lambda_H(0)}{\lambda_H(T)}\right]^2 = 1 - \frac{1}{6} \frac{t^{3/2}}{\Gamma(\frac{3}{2})\zeta(\frac{3}{2})} \int_0^\infty \frac{x^{3/2}\,dx}{\sinh^2(\frac{\xi}{2})} , \tag{4}$$

where

$$\xi = \sqrt{x^2 + \delta x} , \qquad \delta = 4[\zeta(3/2)]^{3/2} \eta \frac{1 - t^{3/2}}{t} .$$

Here $\Gamma(x)$ is the gamma function, and $\zeta(x)$ is the Riemann zeta function. How these results are affected by an anisotropy of the crystal deserves a few words. Introducing the inverse-mass tensor $(1/m^{**})_{ik}$, in which the masses of the particles in the (x,y) plane are assumed to be equal and are denoted by $m_{\parallel}^{**}$, while those along the z axis are denoted by $m_{\perp}^{**}$, we find results similar to (1) for the penetration depths $\lambda_H^{\parallel}$ and $\lambda_H^{\perp}$,

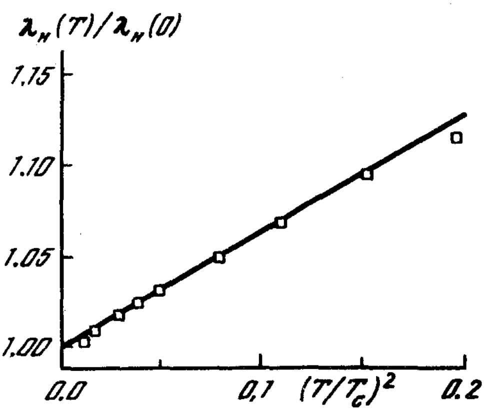

FIG. 1.

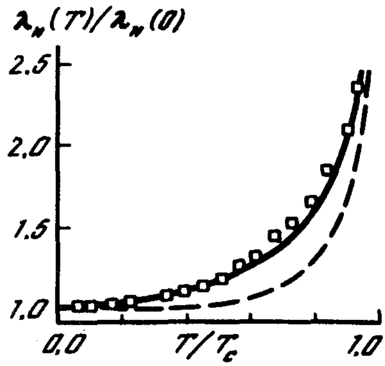

FIG. 2.

in which the masses m^{**} must be replaced by $m_{\parallel}^{**}$ and $m_{\perp}^{**}$, respectively. [Here $\lambda_H^{\parallel}$ is to be understood as the penetration depth in the (x,y) plane.] The temperature dependence of the quantity $\lambda_H(0)/\lambda_H(T)$ remains the same as before.

For a comparison with experiment we used the results of Ref. 4. From the data reported there we selected the results obtained for samples with the largest grain size and the strongest diamagnetic screening, as the most reliable results. Figure 1 compares the low-temperature data with a theoretical curve ($\eta = 0.02$); Fig. 2 makes the comparison for the entire temperature interval. The power-law behavior at low temperatures differs qualitatively from the law $(1 - (T/T_c)^4)^{-1/2}$, which interpolates the experimental data for ordinary BCS superconductors (the dashed line). This power-law behavior can be described satisfactorily in the model of a charged short-range Bose gas.

The long-range Coulomb potential in the new high-temperature superconductors can be suppressed by a screening by light carriers from broad bands; if the latter are absent, they can be suppressed by the giant static dielectric constant of the lattice, which can in principle reach values $\epsilon = 10^5$, as in BaBiPbO. The corresponding plasma frequency of the oscillations of a heavy charged Bose liquid is quite low:

$$\omega_p = \frac{c}{\sqrt{\epsilon}\,\lambda_H} \lesssim 20\,\mathrm{K}\,. \tag{5}$$

It is important to note that the model of a short-range charged Bose gas with a gapless excitation spectrum explains the experimental results (which seem at first glance to be contradictory), according to which a gap is observed in the IR and tunneling spectrum but not in the dependence $\lambda_H(T)$. As has been pointed out previously,[6] in the former case we are seeing a quasigap caused by a state density of the linear excitation spectrum, (3), at low energies which is lower than the single-particle state density in a normal Bose gas ($T \geqslant T_c$). The mean free path of the carriers in the new high-temperature superconductors with a resistivity $\rho \approx 100\ \mu\Omega\cdot$cm is at least an order of magnitude greater than the coherence length ($\zeta_0 \approx 10$ Å). Accordingly, the large value of $\lambda_H(0)$ at carrier densities $n > 10^{21}$ cm^{-3} seems to us to be unambiguous evidence for an anomalously large effective mass of these carriers. Recent measurements of λ_H for single crystals[9,10] generally confirm the existence of the anomalies in λ_H which we are discussing here and which have been seen previously in ceramic samples. For example, the absolute values found for $\lambda_H(0)$ in experiments with $EuBa_2Cu_3O_{7-\delta}$ single crystals are[9] $\lambda_H^{\parallel} = 1480$ Å and $\lambda_H^{\perp} = 5110$ Å. Measurements of the temperature dependence $\lambda_H(T)$ in $YBa_2Cu_3O_{7-\delta}$ (Ref. 10) very convincingly confirm the results of Ref. 4, which we selected for comparison with the theory. We believe that an interpretation of the experimental results on μSR presently available is hindered by the complex distribution of the magnetic field in ceramics.

We thus see that experiments carried out to measure λ_H can be described quite well by the model of a heavy charged Bose gas. In principle, the charged bosons themselves could be of various species. However, tunneling experiments; measurements of the heat capacity, the thermal conductivity, the thermal emf, and the isotopic effect; and certain other experiments provide evidence in favor of a bipolaron nature of the heavy carriers.[11]

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