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Lower Critical Field $H_{c1}$ and Barriers for Vortex Entry in Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ Crystals

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The experimental determination of $H_{c1}$ has been a challenging and controversial problem since the beginning of HTSC research. Not only the values reported for the first flux penetration field $H_p$ are scattered over a wide field range, but in strongly layered superconductors with $H$ perpendicular to the planes, striking features have been observed, such as a marked upturn of $H_p$ for temperatures $T < T_c/2$ [1–6]. The origin of the positive curvature of $H_p$ at low temperatures has been the matter of various speculations. Different mechanisms have been proposed: Bean-Livingston surface barriers [1,3,4,6], bulk pinning [2,6], a modification of the character of the field penetration in layered structures [5], a low-temperature enhancement of the superconducting order parameter in the normal layers [7], etc.

For type-II superconductors, there are at least two kinds of barriers which hinder the system from reaching a thermodynamic equilibrium state: (i) Surface and geometrical barriers [8–10] governing vortex penetration into the superconductor; (ii) bulk pinning barriers governing vortex motion in the superconductor. For the determination of $H_{c1}$, the relevant barriers are those which govern vortex entry into the sample. Of particular interest in this context is the phenomenon of vortex creep over surface and geometrical barriers. From the theoretical point of view this subject has been discussed in Refs. [10] and [11]; however, to our knowledge no systematic experimental study has been carried out so far.

In this Letter we investigate the dependence of $H_p$ on the magnetic field sweep rate $dH/dt$ of isothermal magnetization curves for the strongly layered Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ (Bi2212) material. This has been done for crystals with different cross sections along the $c$ axis. For specimens with ellipsoidal cross sections geometrical barriers can be neglected [10,12], so that the relevant barriers for vortex entry are expected to be surface barriers. Indeed, for our specimen with an ellipsoidal-like cross section, the $H_p$ vs $dH/dt$ dependence obtained at high sweep rates is well described by vortex creep over surface barriers [11]. However, for decreasing $dH/dt$ rates, these barriers are observed to collapse. The subsequent saturation of $H_p$ at lower rates indicates that the system is in equilibrium with respect to surface barriers. This is interpreted as strong evidence that we have reached the lower critical field $H_{c1}$ in our measurements. The obtained low-temperature dependence of $H_{c1}$ is in good agreement with recent microwave absorption measurements of the penetration depth $\lambda$ [13,14]. On the other hand, for specimens with rectangular cross sections, no significant vortex creep over the relevant barriers for vortex entry could be observed, in consistency with theoretical results for geometrical barriers [10].

The crystals under investigation have been grown with different techniques and have different shapes. The specimen with an ellipsoidal-like cross section along the $c$ axis has been grown in a ZrO$_2$ crucible [15] and is approximately $1 \times 1.3 \times 0.05$ mm$^3$ in size. The specimens with rectangular cross sections along the $c$ axis have been grown with the traveling solvent floating zone method.
[16] and are slightly larger in size. The experiment is performed in a noncommercial SQUID magnetometer where the sample is stationary in the pickup coil. The field $H$ is supplied by a superconducting coil working in a non-persistent mode. For the magnetization curve measurements, the sample is zero field cooled from above $T_c$ and stabilized at a fixed temperature (the residual field of the cryostat is $<10$ mOe). Further, the field $H$ is applied at a fixed rate $dH/dt$. For $H_p$, we select the field at which a deviation from Meissner shielding occurs (see the inset of Fig. 3, shown below).

Figure 1 shows the dependence of $H_p$ on the field sweep rate $dH/dt$ at three characteristic temperatures for the specimen with an ellipsoidal-like cross section. A sharp step in the $H_p$ vs $dH/dt$ curves is observed at all three temperatures. At high sweep rates, the curves display a finite slope which is most pronounced at $T = 11$ and 17 K and to a lesser degree at $T = 61$ K. These slopes are in good agreement with the theoretical results of Burlachkov et al. [11] for vortex creep over surface barriers (see the analysis below). Proceeding from the creep regime towards lower rates, the sharp drop signals a collapse of the surface barriers (or equivalently a heat pulse) which leads to a rapid flux entry, possibly in terms of flux jumps. Indeed, an avalanche type of flux penetration into a type-II superconductor has been reported by Durán et al. [17]. On decreasing the sweep rates even further, the flux penetrates into the sample through the occurrence of rare events when the applied magnetic field reaches $H_{c1}$. We thus identify the saturated value of $H_p$ with the lower critical field $H_{c1}$ and use it to obtain information about the nature of the superconducting state.

Analogous measurements have been done on thin specimens with rectangular cross sections. According to Zeldov et al. [10], for such specimens, geometrical barriers build up over a distance $s$ from the sample edges, where $s$ is the sample thickness. The energy required to overcome such a barrier is macroscopic $\sim \varepsilon_0 s$, so that vortex creep over geometrical barriers is expected to be very weak [here $\varepsilon_0 = (\Phi_0/4\pi \lambda)^2$ and $\Phi_0$ is the unit flux]. As a matter of fact, we were not able to detect significant vortex creep over geometrical barriers within the experimental time scale $10^{-5} \leq dH/dt \leq 10^{-1}$ Oe/s.

We proceed with a brief discussion of vortex creep over surface barriers. As shown by Burlachkov et al. [11], the surface barrier for pancake vortices is given by $U = \varepsilon_0 d \ln(0.76 H_c / H_p)$, where $H_c$ is the thermodynamic critical field and $d$ the interlayer distance. During the time $t$, thermal creep allows vortices to overcome barriers of size $U(t) = T \ln(t/t_0)$, where $t_0$ is a “microscopic” time scale [18]. Equating the two expressions one obtains the time dependence of the penetration field for the pancake-vortex regime,

$$H_p(t) = H_c (t/t_0)^{-T/\varepsilon_0 d}. \quad (1)$$

With the definition $T_0 = \varepsilon_0 d / \ln(t/t_0)$, Eq. (1) becomes $H_p = H_c \exp(-T/T_0)$. At higher temperatures creep proceeds via half-loop excitations of vortex lines [11]. The creation of a half-loop saddle configuration involves an energy $U = \pi \varepsilon \varepsilon_0^2 \ln(j_0/|j|)/2\Phi_0 j$, where $j_0$ is the depairing current density and $\varepsilon = (m/M)^{1/2} < 1$ is the anisotropy parameter. With a similar analysis as above, the time dependence of $H_p$ for the vortex-line regime takes the form

$$H_p(t) = H_c \frac{\pi}{2\sqrt{2}} \frac{\varepsilon \varepsilon_0 \xi}{T \ln t/t_0} \ln^2 \left( \frac{H_c}{H_p} \right), \quad (2)$$

where $\xi$ is the coherence length. According to Ref. [11], this expression gives a temperature dependence $H_p \propto (T_c - T)^{3/2}/T$. The crossover between the two regimes occurs when the size of the half-loop excitation is of the order of the interlayer distance $d$. This is expected to occur at a temperature $T^* = T_0(T^*) \ln(d/\varepsilon \xi)$.

According to (1), creep of pancake vortices over surface barriers results in a linear dependence of $\ln(H_c/H_p)$ on $\ln(t/t_0)$ with a slope $T/\varepsilon_0 d$. This dependence is given in Fig. 2(a) using the data at high cycling rates in Fig. 1. For $T = 11$ and 17 K, the values of $T/\varepsilon_0 d$ obtained from the fits are in good agreement with the values estimated from the penetration depth $\lambda_{ab}(T)$ (see Table I). On the other hand, for $T = 61$ K the above analysis for pancake vortices is not satisfactory since the $T/\varepsilon_0 d$ value obtained from the fit is 1 order of magnitude smaller than the calculated value. According to (2), for creep of vortex lines over surface barriers $(H_c/H_p) \ln^2(H_c/H_p)$ vs $\ln(t/t_0)$ is linear with a slope $2\sqrt{2}/T \pi \varepsilon \varepsilon_0 \xi$. For the $T = 61$ K data at high cycling rates in Fig. 1, such a representation is given in Fig. 2(b). The slope $2\sqrt{2}/T \pi \varepsilon \varepsilon_0 \xi$ obtained from the fit is $15 \pm 3$, in good agreement with the value 16, estimated with the help of the Ginzburg-Landau formula for the coherence length $\xi(T)$, with
the time origin $t = 0$ at $H = H_{c1}$. The critical field $H_{c1}(T) = H_c(0)(1 - T/T_\ast)$ is calculated with $\lambda_{ab}(0)$ as obtained from the $H_{c1}$ data in Fig. 3 and the coherence length $\xi(0) = 25\,\text{Å}$. We assumed $\xi_0 = 10^{-6}\,\text{s}$. (b) $T = 61\,\text{K}$ data at high sweep rates in Fig. 1 in a convenient representation. The lines in (a) and (b) are fits according to Eqs. (1) and (2), respectively.

In Fig. 3 we make use of the curves in Fig. 1 to determine $H_p(T)$ with three criteria: (i) $H_p$ data ($\bigcirc$) are taken at low sweep rates so as to guarantee that they lie in the regime where the system is in equilibrium with respect to surface barriers [these data represent $H_{c1}(T)$], (ii) $H_p$ data ($\ast$) are taken at high rates so that they always lie in the creep regime, and (iii) the field $H$ is swept at the constant rate $dH/dt = 8 \times 10^{-3}\,\text{Oe/s}$ for all temperatures. In this case, the values of $H_p$ ($\bigtriangledown$) lie on the $H_{c1}$ curve for $T \geqslant 60\,\text{K}$, they go through the steplike crossover for temperatures $30 \leqslant T \leqslant 60\,\text{K}$, and, finally, for $T \leqslant 30\,\text{K}$ they lie in the creep regime.

![Figure 2](image2.png)

**FIG. 2.** (a) $T = 11$ and $17\,\text{K}$ data of the upper plateau in Fig. 1 in a different representation. Here we choose the time origin $t = 0$ at $H = H_{c1}$. The critical field $H_{c1}(T) = H_c(0)(1 - T/T_\ast)$ is calculated with $\lambda_{ab}(0)$ as obtained from the $H_{c1}$ data in Fig. 3 and the coherence length $\xi(0) = 25\,\text{Å}$. We assumed $\xi_0 = 10^{-6}\,\text{s}$. (b) $T = 61\,\text{K}$ data at high sweep rates in Fig. 1 in a convenient representation. The lines in (a) and (b) are fits according to Eqs. (1) and (2), respectively.

![Figure 3](image3.png)

**FIG. 3.** Temperature dependence of the first flux penetration field $H_p$ of the specimen with an ellipsoidal-like cross section for different field sweep rates: (C) $dH/dt = 1 \times 10^{-4}\,\text{Oe/s}$, (a) $dH/dt = 5 \times 10^{-2}\,\text{Oe/s}$, and (>) $dH/dt = 8 \times 10^{-3}\,\text{Oe/s}$. The inset shows the deviation from the Meissner slope for $T = 11$ and $61\,\text{K}$. The upper curves (■) are measured at a rate $dH/dt = 5 \times 10^{-2}\,\text{Oe/s}$ and the lower ones (□) at $dH/dt \leqslant 1 \times 10^{-4}\,\text{Oe/s}$. The data are scaled with the demagnetization factor $N = 0.96(\pm 15\%)$.
fluctuations of the order parameter around its mean-field renormalization of the vortex-line free energy due to can be understood on the basis of an entropic downward

planes has been estimated to be smaller than $2\pi$ from measurements at magnetic fields in the range $H_{\parallel} \ll H < H_{\perp}$ (see Ref. [21]; here $H_{\parallel}$ and $H_{\perp}$ are the parallel and perpendicular penetration fields). For this configuration geometrical barriers are irrelevant. As shown in Fig. 4, for temperatures $T$ between 10 and 70 K, $H_{c1}(T)$ has an approximately linear behavior. According to Ref. [22], for temperatures not so close to $T_c$, the correct description of $H_{c1}$ in strongly layered materials is $H_{c1} = \Phi_0/(4\pi \lambda_{ab} \lambda_c) \left[ \ln(\lambda_{ab}/d) + 1.12 \right]$. With this formula we calculated the penetration depth in $c$-direction $\lambda_c$ using the previously determined $\lambda_{ab}$ data. For $T \leq 40$ K, $\lambda_c$ is approximately linear in $T$ with a slope of 0.1 $\mu$m/K. The extrapolated $T = 0$ value is $\lambda_c(0) = 15$ $\mu$m. Figure 4 further shows $H_{p}$ data at high sweep rates ($\ast$). Contrary to the case for $H \perp ab$, no upturn of $H_{p}$ is observed here for temperatures $T \leq T_c/2$. We attribute this difference to the absence of the strong pancake-vortex creep regime at low temperatures for $H \parallel ab$. From the $H_{p}$ vs $dH/dt$ dependence in the inset of Fig. 4, we obtain an activation barrier for vortex entry $U = 200$ K at $T = 12$ K and $U = 1300$ K at $T = 81$ K.

Close to $T_c$ a downward bending of $H_{c1}$ is observed for $H \perp ab$ as well as for $H \parallel ab$ (see Figs. 3 and 4). This can be understood on the basis of an entropic downhill renormalization of the vortex-line free energy due to fluctuations of the order parameter around its mean-field form [23]. The decrease in the free energy $f_1 = \epsilon_1 - Ts_1$ then leads to a drop in $H_{c1}$ as $T$ approaches $T_c$ (here, $\epsilon_1$ is the line energy of the vortex excitation and $s_1$ is the line entropy).

To conclude, isothermal magnetization curves were measured on Bi2212 crystals for configurations of negligible geometrical barriers. At very low sweep rates we found saturated values of the penetration field $H_{p}$ which we interpret as the true lower critical field $H_{c1}$. The low-temperature dependence of $H_{c1}$ is consistent with recent magnetic penetration depth measurements [13,14] suggesting a superconducting state with nodes in the gap function. Based on the results obtained for Bi2212, we argue that the frequently reported low-temperature upturns of $H_{p}$ in strongly layered superconductors with $H$ perpendicular to the planes [1–6] can be explained in terms of measurements where the system is out of equilibrium with respect to barriers for vortex penetration.

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