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Giant flux creep through surface barriers and the irreversibility line in high-temperature superconductors

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Magnetic and transport phenomena in high-temperature superconductors due to magnetic flux relaxation and transport over the surface barrier are investigated. Vortex dynamics controlled by the penetration both of pancake vortices and vortex lines is discussed. The penetration field H_p for pancakes decays exponentially with temperature. The size of the magnetization loop is determined by the decay of H_p during the process of relaxation, but its shape remains unchanged. The irreversibility line associated with the pancake penetration is given by $H_{\rm irr} \propto \exp(-2T/T_0)$ and may lie both above and below the melting line.

It has been recognized[1–4] that the Bean-Livingston surface barriers play an important role in vortex dynamics in high-temperature superconductors (HTSC) and control the penetration field $H_p > H_{c1}$,[1,2,4] where H_{c1} is the first critical field. For an absolutely perfect surface $H_p \cong H_c \cong \kappa H_{c1}/\ln\kappa$, where H_c is the thermodynamic field, $\kappa = \lambda/\xi$, λ is the penetration depth, and ξ is the coherence length. Because of large $\kappa \simeq 100$ one gets $H_c/H_{c1} \simeq 20$ in HTSC's, but due to surface imperfections H_p never reaches H_c and lies somewhere in between: $H_{c1} < H_p < H_c$.[2] For the case of weak bulk pinning the surface barriers determine the magnetization[2] and the position of the irreversibility line.[5] In clean $YBa_2Cu_3O_7$ (YBCO) crystals such a domination of surface irreversibility over the bulk one was proved at $T > 85$ K (with $T_c = 92$ K),[2] where bulk pinning becomes exponentially small with respect to surface barrier due to thermal fluctuations of vortices.[6]

The field distribution and magnetization curves at $H > H_p$ due to the surface barrier were calculated in Refs. 7, 8 and 3 for the whole field range where the London approximation holds, i.e., up to fields of order H_{c2}. However in HTSC the field H_p appeared to show a strong (exponential) temperature dependence and the irreversibility line $H_{\rm irr}(T)$ was found to lie much below H_{c2}.[1,2,4,5] These phenomena can be attributed to the important role of thermally activated processes in these materials. Suppression of the penetration field due to creep of pancakes through the barrier in strongly layered materials was considered in Ref. 1, and the barriers for flux entry and exit determined by the motion of vortex lines were calculated in Ref. 3. In this paper we find the temperature dependence of the whole magnetization loop including that of the irreversibility field for both two-dimensional (2D) and 3D cases. We also calculate the field, temperature, and current dependence of the thermally activated resistivity of superconductor due to the flux motion over the surface barrier.

Consider a strongly layered superconductor in external magnetic field $\mathbf{H} \| c$ and calculate the energy profile for a pancake surmounting the surface barrier following.[1] The energy of interaction of the pancake at the distance $r \ll \lambda$ with its mirror image and the Meissner current is given by

$$U(r) = \epsilon_0 d \ln\frac{2.94r}{\xi} - \frac{\Phi_0}{c} djr, \tag{1}$$

where $\epsilon_0 = (\Phi_0/4\pi\lambda)^2$, Φ_0 is the unit flux, and d is the period of the layered structure. The maximum of this energy corresponds to the distance $r_0 = c\epsilon_0/\Phi_0 j = (3\sqrt{3}/4)\xi j_0/j = \xi H_c/\sqrt{2}H$, where $j_0 = cH_c/3\sqrt{6}\pi\lambda$ is the depairing current and we used $j = cH/4\pi\lambda$ to relate the Meissner current with H. For the sake of simplicity we consider a geometry where the superconductor occupies half space $x > 0$. The modification of the results due to demagnetization effects is discussed at the end of the paper. The activation energy required to surmount the barrier (1) is

$$U_0 = \epsilon_0 d \ln(1.1 r_0/\xi) = \epsilon_0 d \ln(0.76 H_c/H).$$

The pancakes, which penetrate after the first one in adjacent layers, require less energy than U_0. Therefore, the process of the vortex line (stack of pancakes) formation is of avalanche-type,[9] and U_0 is the activation energy for penetration of the whole stack.

Since the probability of surmounting is proportional to $\exp(-U_0/T)$, the penetration field H_p should be determined[10] from the equation $U(H_p) = T\ln(t/t_0)$, where t is the time scale of the experiment and t_0 is the microscopic time scale which is determined by the preexponential factor in the penetration rate $t_0 \sim 10^{-10} - 10^{-12}$ s.[1] Then one obtains[1]

$$H_p \simeq H_c \exp(-T/T_0), \tag{2}$$

with $T_0 = \epsilon_0 d/\ln(t/t_0)$. The penetration field decays according to the power law as

$$H_p \simeq H_c (t/t_0)^{-T/\epsilon_0 d}.$$

This result holds as long as the cost of the Josephson energy ϵ_J to create a single pancake at distance r_0 at the surface, $\epsilon_J = (\epsilon^2 \epsilon_0 r_0^2/2d)\ln(d/\epsilon r_0)$ is small as compared to U_0 or, equivalently, $r_0 < \Lambda_J = d/\epsilon$,[10] where $\epsilon = \sqrt{m/M} < 1$ is the parameter of anisotropy. This condition corresponds to the field range $H > H_c \xi \epsilon/d \simeq \epsilon H_c$ in HTSC materials. Therefore in layered (Bi- and Tl-based) compounds with $\epsilon \simeq 10^{-2}$ the region, where the flux penetration is controlled by the single pancake creep over the surface barrier, is very wide. For realistic values of the superconducting parameters for $Bi_2Sr_2CaCu_2O_8$ (BSCCO) ($\lambda \simeq 2000$ Å,[11] $d = 15$ Å) one obtains $\epsilon_0 d \simeq 700$ K. Taking the typical value of $\ln(t/t_0) \simeq 30$ one obtains $T_0 \simeq 23$ K in reasonable agreement with the experiment.[1–4] Note that surface imperfections reduce the prefactor in Eq. (2), but cannot change the exponent.

Now we turn to the calculation of the magnetization loop. The structure of the vortex lattice near the surface of the sample was calculated by Ternovskii and Shekhata,[7] but in what follows we use a simplified consideration given by Clem[8] (see also Ref. 3). The Clem model predicts the existence of a vortex-free region of the width x_f near the surface of the sample and the constant vortex density for $x > x_f$. The distribution of the field in the vortex-free region ($x < x_f$) is given by $h(x) = B\cosh[(x - x_f)/\lambda]$ with the current density at the surface $j = (c/4\pi)B\sinh(x_f/\lambda)$, whereas at $x > x_f$ one gets $h(x) = B = \text{const}$. The current relaxation due to vortex creep over the barrier is determined by the condition $U(j) = \epsilon_0 d \ln(3\sqrt{6} j_0/4j) = T\ln(t/t_0)$. Therefore, $B\sinh(x_f/\lambda) = H_p(T)$ with $H_p(T)$ from Eq. (2). Making use of the boundary conditions $H = B\cosh(x_f/\lambda)$ one obtains straightforwardly $B = \sqrt{H^2 - H_p^2}$ for $H > H_p$, and the magnetization is

$$-4\pi M = H - \sqrt{H^2 - H_p^2} \approx H_p^2/2H, \quad \text{for} \quad H \gg H_p. \tag{3}$$

The magnetization loops preserve their shape during the relaxation process, and creep results only in the thermal renormalization of the penetration field $H_p(T,t)$. Since for $H \gg H_p$ one gets $-4\pi M \approx H_p^2/2H$, the width of the magnetization loop at high fields drops with the temperature as $\exp(-2T/T_0)$, i.e., much faster than H_p itself.

The above consideration holds as long as the hopping distance $r_0 < x_f$. As temperature increases, the hopping distance increases accordingly and when r_0 becomes of order of x_f, the barrier for pancake saturates and the pancake penetration becomes reversible. This occurs at $r_0 \sim \xi H_c/H_p \sim x_f$, which corresponds to $-4\pi M \simeq H_p^2/2H \simeq -4\pi M_{eq} \simeq (\Phi_0/8\pi\lambda^2)\ln(\eta H_{c2}/H)$, where η is the constant of order of unity. The irreversibility line is then given by

$$H_{irr} \simeq \frac{H_p^2}{H_{c1}} \frac{\ln\kappa}{\ln(\eta H_{c2}/H_{irr})} = \frac{H_{c2}}{\ln(\eta H_{c2}/H_{irr})} \exp(-2T/T_0). \tag{4}$$

At high-enough temperatures $T > T_0$ the expression (4) reduces to $H_{irr} \approx H_{c2}(T_0/2T)\exp(-2T/T_0)$. This dependence agrees favorably with numerous experimental data and provides a reasonable explanation for the observed behavior of the irreversibility line[12–14] in highly anisotropic HTSC's. Note that the irreversibility temperature $T_{irr} \simeq \epsilon_0 d \ln(H_{c2}/H)/2\ln(t/t_0)$ may lie both above and below the 2D melting temperature $T_m \simeq \epsilon_0 d/8\sqrt{3}\pi$ depending on the magnetic field range involved.

Above the irreversibility line the barrier for the pancake motion is smaller than $T\ln(t/t_0)$ and flux penetration is reversible. However, this barrier can give rise to a linear resistivity. The height of this barrier can be estimated from Eq. (1) replacing j by the equilibrium surface current $j_{eq} = c\sqrt{HM_{eq}}/(\sqrt{2}\pi\lambda)$:

$$U \approx \frac{\epsilon_0 d}{2} \ln\left(\frac{0.6 H_{c2}}{H\ln(H_{c2}/H)}\right). \tag{5}$$

A more elaborate calculation of U at equilibrium which accounts for the lattice deformations was performed in Ref. 15. The resulting expression for resistivity acquires the form $\rho = \rho_0 \exp(-U/T)$, in good agreement with experimental results.[16,17] If the transport current density near the surface becomes larger than $j_0 \xi (B/\Phi_0)^{1/2}$, the barriers decrease and the current-voltage characteristic becomes nonlinear being characterized by the activation barrier

$$U(j) \approx \epsilon_0 d \ln(j_0/j). \tag{6}$$

Note that the logarithmic dependence (6) of the activation energy on current gives rise to a power-law behavior of the current-voltage characteristic $V \propto j^{\epsilon_0 d/T}$. This behavior has been observed in several experiments.[18,19] Another point one has to be aware of when analyzing the experimental data is that the surface current flows through a thin layer ($\sim x_f$) near the surface, therefore, the current density should be determined as $I/x_f \times$(relevant length) rather than $I/$(relevant area).

The activation barriers depending logarithmically on the applied field were observed in experiments on multilayered compounds[16] and interpreted as the result of the disorder stimulated dissociation of dislocation pairs in 2D vortex lattice.[20] Here we propose a simpler explanation of this dependence. However, our estimations for the activation energy from Eq. (5) give $U=350$ K $\ln(100T/H)$, with $\lambda=2000$ Å for BSCCO, and $U=550$ K $\ln(200T/H)$, with $\lambda=1400$ Å for YBCO superlattice. The prelogarithmic factor is approximately 2–3 times larger than that observed experimentally. Since our theory expresses the barrier through d and λ which are well measured quantities with no other fitting parameters, this discrepancy is the relevant one. The possible explanation of this discrepancy is that λ near the surface can be smaller than in the bulk because of the suppression of superconductivity due to the variation of oxygen contents near the surface.

For less anisotropic HTSC materials (such as YBCO) or for Bi- and Tl-based compound at high temperatures where $H_p(T)$ drops below ϵH_c, the Josephson interaction of the pancakes becomes important and the continuous 3D anisotropic model is more appropriate. Then the thermal activation over the surface barrier occurs via a creation and further expansion of the half-loop excitation.[3,21,22] The energy as a function of the size of the loop along layers r is then given by

$$U(r)=\pi\epsilon\epsilon_0 r\ln(r/\sqrt{2}\xi)-\frac{\pi}{2}\frac{\Phi_0}{c}\epsilon j r^2,$$

which leads to the radius of the critical loop $r_0\simeq c\epsilon_0\ln(j_0/j)/\Phi_0 j$ and the activation energy

$$U(j)=\frac{\pi\epsilon\epsilon_0^2 c\ln^2(j_0/j)}{2\Phi_0 j}.$$

Proceeding further analogously to the 2D case, i.e., relating current to the magnetic field near the surface and taking into account that the characteristic barriers $U(j)=T\ln(t/t_0)$ one easily obtains the expression for the penetration field

$$H_p\simeq H_c\frac{\pi}{2\sqrt{2}}\frac{\epsilon\epsilon_0\xi}{T\ln(t/t_0)}\ln^2(H_c/H_p), \tag{7}$$

providing that creep is strong enough and H_p is smaller than H_c.

The magnetization loops for $H>H_p$ are given by the same formulas as in the 2D case (3) with H_p from Eq. (7). Since this equation gives $H_p\propto(T_c-T)^{3/2}/T$, the width of the hysteresis loop decrease with temperature and field as $M\propto(T_c-T)^3/TH$.

To estimate the location of the irreversibility line we equate the irreversible magnetization to $M_{\rm eq}$ and obtain

$$H_{\rm irr}^s\simeq\frac{\pi}{256}\frac{\epsilon\epsilon_0^2\Phi_0\ln^3(H_{c2}/H_{\rm irr}^s)}{T^2\ln^2(t/t_0)}. \tag{8}$$

The obtained result describes the irreversibility line controlled by the single vortex creep. However the difference between the 3D and 2D cases is that in the 3D case the single vortex penetration becomes impossible near the equilibrium state of vortex lattice. Indeed, the implantation of the newly penetrating vortex line into the ideal vortex configuration near the surface requires the infinite energy. This means that near the equilibrium state vortices penetrate through the surface via creation of critical nuclei consisting of several vortex lines. This collective penetration starts when the magnetization M of the sample approaches its equilibrium value $\Delta M=M_{\rm eq}-M\approx 0.3\Phi_0/(4\pi\lambda)^2$ (this corresponds to external field $H=H_{\rm eq}-0.3\Phi_0/4\pi\lambda^2$, where $H_{\rm eq}=B-4\pi M_{\rm eq}$).[21,15] Near, but not too close to the equilibrium $[M_{\rm eq}-M>\Phi_0/(4\pi\lambda)^2\ln^{3/2}(H_{c2}/B)]$ the magnitude of the activation barrier grows as

$$U(H,\Delta M)=U_1\frac{\Phi_0/(4\pi\lambda)^2}{\Delta M}. \tag{9}$$

Here

$$U_1=0.2\epsilon\epsilon_0\left(\frac{\Phi_0}{H}\right)^{1/2}\left(\ln\frac{H_{c2}}{H}\right)^{3/2}$$

is the barrier at the transition point from the single vortex to the collective penetration (see Ref. 21 for a detailed discussion). The irreversibility line is determined by the condition $U(H_{\rm irr},\Delta M_{\rm exp})=T\ln(t/t_0)$, where $\Delta M_{\rm exp}$ is experimentally resolvable irreversible magnetization (usually $0.1-1$ G). This gives

$$H_{\rm irr}^{\rm col}\simeq\left(\frac{\Phi_0/(4\pi\lambda)^2}{\Delta M_{\rm exp}}\right)^2 H_{\rm irr}^s>H_{\rm irr}^s \tag{10}$$

at $\Delta M_{\rm exp}<\Phi_0/(4\pi\lambda)^2$.

The irreversibility line is determined by the relation between $H_{\rm irr}^s$, $H_{\rm irr}^{\rm col}$, and B_m. If $H_{\rm irr}^s>B_m$, then the onset of the irreversibility is controlled by the single vortex penetration and the position of the irreversibility line is given by Eq. (8). If $H_{\rm irr}^{\rm col}>B_m>H_{\rm irr}^s$ the irreversibility line coincides with the melting line, and, at last, for $B_m>H_{\rm irr}^{\rm col}$ the irreversibility line is determined by Eq. (10). H_s differs from the conventional expression for the melting line $B_m\simeq 8\epsilon^2\epsilon_0^2c_L^4\Phi_0/T^2$ by the factor of $\ln^{3/2}(H_{c2}/H)/24\ln(t/t_0)$ instead of the Lindemann factor c_L^2. These factors are of the same order of magnitude which means that all three cases are possible.

On the contrary, in much more anisotropic BSCCO material the melting line at high temperatures may be close to H_{c1} and the irreversibility line determined by Eq. (8) is probably located *above* the melting line. Therefore in BSCCO compounds the irreversibility line at high temperatures may arise from the surface barrier rather than from the freezing of the vortex liquid into a vortex solid. Above the melting the activation energy over the surface barrier remains finite and its maximum magnitude at $M=M_{\rm eq}$ is

$$U_{\rm 3D}\simeq 0.2\epsilon\epsilon_0(\Phi_0/H)^{1/2}\ln^{3/2}(H_{c2}/H)\propto\frac{T_c-T}{\sqrt{H}}, \tag{11}$$

corresponding to the radius of the half-loop excitation $r_0\simeq x_f/2$. This barrier, being independent of current, gives rise to a linear resistivity $\rho\propto\exp(-U_{\rm 3D}/T)$ which can explain thermally activated resistivity in this materials.

In the presence of the bulk pinning the magnetization is just the sum of the bulk and surface contributions.[3] The fast drop of the surface component in the fields $H>H_p$ gives rise to a widely observed (see, for example, Ref. 23) sharp maximum of the magnetization at small $[H\simeq H_p(T)]$ fields, the height of the maximum being approximately equal to $4\pi H_p(T)$.

Finally, let us discuss briefly the effect of the sample shape.[24] Consider a rectangular sample with the thickness h, and the smallest transverse size w. The field near the surface of such a sample in the Meissner state is $\sqrt{w/d}$ times larger than the applied field, accordingly the penetration field is $\sqrt{d/w}H_p$ (Ref. 24) (note the distinction from the commonly used demagnetization factor of d/w for the sample of the ellipsoidal shape). The slope of the magnetization curve in the Meissner state below H_p is equal to $(H/4\pi)(w/d)$. In the field much greater than the penetration field $\sqrt{d/w}H_p$ the magnetization is equal to $M \simeq H_p^2/2H$ and does not depend on the demagnetization factor.

To conclude, we have shown that the creep over the surface barriers can explain the observed temperature dependence of magnetization curves and the position of the irreversibility line in different families of HTSC materials. It is important that the critical current due to the bulk pinning drops very fast at high magnetic fields [exponentially or as H^{-3} (Ref. 10)], whereas the surface critical current decays only as $1/H$.[6] As a result even in the samples with the strong bulk pinning controlling the width of the hysteresis loop well below the irreversibility line, the *position* of the irreversibility line itself can be still determined by surface barriers.[25]

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