Peak Effect in Twinned Superconductors

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A sharp maximum in the critical current J_c as a function of temperature just below the melting point of the Abrikosov flux lattice has recently been observed in both low- and high-temperature superconductors. This peak effect is strongest in twinned crystals for fields aligned with the twin planes. We propose that this peak signals the breakdown of the collective pinning regime and the crossover to strong pinning of single vortices on the twin boundaries. This crossover is very sharp and can account for the steep drop of the differential resistivity observed in experiments.

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The discovery of the high-temperature copper-oxide superconductors has renewed the experimental and theoretical interests in the properties of the mixed state of type-II superconductors in a magnetic field [1]. Various experimental techniques, including standard current versus voltage curves, are used to measure the critical current density J_c needed to depin the flux-line array and to investigate its temperature and field dependence. Naively J_c is expected to decrease monotonically as the temperature or the applied field are raised towards the mean field $H_{c2}(T)$. It has, however, been known for some time that an abrupt increase in J_c as a function of field or temperature can occur in conventional low-temperature superconductors near H_{c2} [2]. A qualitative explanation of this phenomenon, referred to as "peak effect," was proposed a long time ago by Pippard [3], who argued that the increase in J_c is associated with the softening of the shear modulus c_{66}. A more quantitative explanation of the peak effect as arising from the softening of all the elastic moduli of the flux lattice near H_{c2} was presented by Larkin and Ovchinnikov [4].

Recently a sharp maximum in J_c as a function of temperature has been observed in both untwinned [5–8] and twinned [9] Y-Ba-Cu-O (YBCO) crystals, as well as in some low-temperature superconductors [10,11]. The new feature is that in this case the peak occurs below H_{c2}, at the temperature T_m, where the flux lattice melts into a flux-line liquid. In view of the old suggestion by Pippard [3] and the work by Larkin and Ovchinnikov [4], it is natural to associate it with the softening of the shear modulus c_{66} at the melting point. In twinned YBCO crystals the peak depends strongly on the orientation of the applied field relative to the twin planes: It is largest for flux motion along the twin planes and external fields aligned with the c axis, and it weakens as the field is tilted out of the plane of the twins [9]. In untwinned YBCO single crystals the peak is much smaller: It shifts towards T_c and becomes less pronounced as the sample purity is increased [6].

In this paper we propose two distinct mechanisms for the peak effect. First by examining the temperature dependence of the critical current from collective pinning by point defects [1,12,13], we show that J_c can exhibit a sharp rise near T_m for a narrow range of magnetic fields due to the abrupt decrease of the shear modulus. This may provide a mechanism for the small peak effect observed in untwinned single crystals. New results on anisotropic collective pinning in samples with a family of parallel twin planes are also presented and show that the same mechanism is, in principle, operative when vortices are pinned collectively by an array of twin planes. On the other hand, collective pinning is very weak in this case and cannot account for the large increase in J_c observed in twinned samples. In the second part of the paper we discuss a second mechanism for a peak in J_c that is operative in twinned samples: the strong pinning of individual vortices on the twins. This mechanism can account for the sharp drop in the resistivity observed in twinned materials.

The critical current density of an elastic medium pinned by weak disorder can be calculated using the collective pinning theory [4]. Weak disorder destroys the translational order of the flux lattice and results in the coherent pinning of vortex bundles of extent L_c and R_c in the directions parallel and perpendicular to the applied field $\mathbf{H}$. The pinning lengths R_c and L_c, defined as the distances at which the lattice distortion due to disorder is of the order of the range ξ of the pinning potential, are determined in terms of the elastic constants of the lattice by balancing the elastic deformation energy against the pinning energy. The critical current J_c is the current where the Lorentz force balances the pinning force, or $BJ_c/c \approx \sqrt{W/V_c}$, where $W = n_p\langle f^2\rangle$ is the mean square pinning force, with n_p the volume density of pins and f the elementary pinning force, and $V_c = R_c^2 L_c$.

We consider a three-dimensional flux-line array in a sample with an external magnetic field aligned with the c axis, which is chosen as the z direction. Disorder

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is described as a quenched random potential per unit length $V(\mathbf{r})$ with zero mean and Gaussian correlations, $\overline{V(\mathbf{r})V(\mathbf{r}')} = \Gamma(\mathbf{r}, \mathbf{r}')$. The overbar denotes the disorder average, and the correlator $\Gamma(\mathbf{r}, \mathbf{r}')$ is determined by the strength and geometry of the disorder. The static elastic deformation of the lattice due to disorder can be evaluated by a perturbation theory in the pinning potential [1,12]. To lowest order in perturbation theory the components of the mean square displacement $U_{ij}(\mathbf{r}) = \overline{\langle \Delta u_i(\mathbf{r}) \Delta u_j(\mathbf{r}) \rangle}$ induced by the random potential [$\Delta u_i(\mathbf{r}) = u_i(\mathbf{r}) - u_i(0)$ and the brackets denote a thermal average] are given by [1,12]

$$U_{ij}(\mathbf{r}) = \int \frac{d\mathbf{q}}{(2\pi)^3} \int \frac{d\mathbf{q}'}{(2\pi)^3} (1 - e^{i\mathbf{q}\cdot\mathbf{r}})(1 - e^{i\mathbf{q}'\cdot\mathbf{r}'}) \times G_{ik}(\mathbf{q}, \omega = 0) G_{jl}(\mathbf{q}', \omega = 0) \tilde{W}_{kl}(\mathbf{q} + \mathbf{q}'), \quad (1)$$

with $G_{ij}(\mathbf{r}, t)$ the elastic Green function of the lattice, $\mathbf{f}_p(\mathbf{r}) = -\hat{n}(\mathbf{r}, t)\vec{\nabla} V(\mathbf{r})$ the pinning force per unit volume, and $\tilde{W}_{kl}(\mathbf{q}, \mathbf{q}') = \overline{f_{p,i}(\mathbf{q}) f_{p,j}(\mathbf{q}')}$ the pinning force correlator. Here $\hat{n}(\mathbf{r}, t) = \sum_n \delta^{(2)}(\mathbf{r}_\perp - \mathbf{r}_n(z, t))$, with $\mathbf{r} = (\mathbf{r}_\perp, z)$, is the coarse-grained microscopic vortex density field [the flux lines are parametrized by their trajectories $\{\mathbf{r}_n(z, t)\}$]. The main contribution to Eq. (1) for the case of interest below comes from the transverse part of the elastic Green function, given by

$$G_{ij}^T(\mathbf{q}, \omega) = \frac{\mathcal{P}_{ij}^T}{-i\omega\zeta + c_{66} q_\perp^2 + c_{44}(q_\perp, q_z) q_z^2}, \quad (2)$$

where $\mathcal{P}_{ij}^T = \delta_{ij} - \hat{q}_{\perp i} \hat{q}_{\perp j}$, with $\hat{q}_{\perp i} = q_{\perp i}/q_\perp$, and c_{66} and $c_{44}(q_\perp, q_z)$ are the shear and tilt moduli of the vortex lattice, respectively. Thermal fluctuations can be incorporated in the perturbation theory by separating out in the vortex positions $\mathbf{r}_n(z, t)$ the deviation from equilibrium due to pinning from that due to thermal effects, as described in [13]. The main effect of thermal fluctuation is the replacement of the upper cutoff $q_0 \approx \xi^{-1}$ in the wave-vector integral by a thermal cutoff $q_T \approx (\xi^2 + \langle u^2 \rangle_{\text{th}})^{-1/2} = q_0(1 + T/T_{\text{dp}})^{-1/2}$. The depinning temperature T_{dp} is defined by $\langle u^2(T_{\text{dp}}) \rangle_{\text{th}} \approx \xi^2$, where $\langle u^2(T) \rangle_{\text{th}}$ is the mean square thermal excursion of the vortices about their equilibrium positions. When the vortex array is described as an elastic continuum, $T_{\text{dp}} = 2a_0^2 \xi^2 \sqrt{c_{66} \hat{c}_{44}}/\lambda$, with $a_0 = \sqrt{\phi_0/B}$ the mean intervortex separation, λ the penetration length in the a-b plane, and $\hat{c}_{44} = c_{44}(q_\perp = 0, q_z = 0)$ [14].

When the flux array is pinned by isotropic point disorder, the random potential is short ranged in all directions and $\Gamma(\mathbf{r}, \mathbf{r}') = \gamma \delta^{(3)}(\mathbf{r} - \mathbf{r}')$ and $\gamma \approx (U_0 \xi^3)^2 n_p [1 + \mathcal{O}(n_p \xi^3)]$, with U_0 the depth of an individual pinning potential per unit length. The pinning force correlator is isotropic, $W_{ij}(\mathbf{q}) = W \delta_{ij} (2\pi)^3 \delta^{(3)}(\mathbf{q})$, with $W = \gamma/\xi^4 a_0^2$. The pinning lengths and the critical current for this disorder geometry have been calculated elsewhere [12], but it is instructive to display the dependence of J_c on the elastic constants. For low defect densities ($R_c > \lambda$, this is the so-called large-bundle regime), the dispersion of the tilt modulus can be neglected and

$$J_c \approx j_0 \left(\frac{j_{\text{sv}}}{j_0} \frac{H_c^3}{4\pi} \right)^3 \frac{3\sqrt{3}B}{16 H_{c2}} \frac{1}{c_{66}^2 \hat{c}_{44}} (1 + T/T_{\text{dp}})^{-11/2}.$$

Here $j_0 = cH_c/3\sqrt{6}\pi\lambda$ is the depairing current and $j_{\text{sv}} = j_0 (W a_0^2 \xi / \epsilon_0^2)^{2/3}$ is the single-vortex critical current density. The critical current contains both an explicit temperature dependence from the thermal smoothing of the pinning potential and an implicit T dependence through the superconductor's parameters that determine the elastic constants. To extract the strong temperature dependence of c_{66} near melting, we write $c_{66} = c_{66}^0 r(T/T_m)$, where c_{66}^0 depends weakly on T and the function r drops sharply from unity to zero at T_m. This is defined by $\langle u^2(T_m) \rangle_{\text{th}} \approx c_L^2 a_0^2$, with $c_L \approx 0.1-0.3$ the Lindemann parameter. The critical current is then $J_c \sim 1/[c_{66}^0 r(T)]^2 [1 + T/T_{\text{dp}}^0 r^{1/2}(T)]^{-11/2}$, where $T_{\text{dp}}^0 = 2a_0^2 \xi^2 \sqrt{c_{66}^0 \hat{c}_{44}}/\lambda$. The temperature dependence of J_c near T_m is controlled by the parameter $\alpha = T_m/T_{\text{dp}}^0 = (c_L a_0/\xi)^2$. For $H = 6$ T and temperatures near T_c, we use $\xi \approx 100$ Å and $c_L \approx 0.2$ to obtain $T_m/T_{\text{dp}}^0 \approx 0.1$. At low temperatures ($T \ll T_{\text{dp}}, T_m$) J_c decreases very slowly with T. At higher temperatures, but still well below T_m, the elastic constants are only very weakly temperature dependent and the temperature dependence of J_c is controlled by thermal fluctuations, yielding a decrease of J_c as T grows. As T_m is approached from below, c_{66} softens and the flux lattice can better adjust to the pinning centers, raising J_c. Finally, at T_m the function $r(T)$ drops sharply in a narrow temperature range giving rise to a sharp maximum in J_c, provided $T_m/T_{\text{dp}}^0 < 1$. This mechanism can therefore yield a peak in J_c only at large fields, where this condition can be satisfied.

For larger defect densities the dispersion of c_{44} is important (this is referred to as the small-bundle regime). In this case

$$J_c \approx j_0 \frac{6\pi\sqrt{3}}{H_c^2} c_{66} (1 + T/T_{\text{dp}})^{1/2} \times \exp\left[-\left(\frac{j_0}{j_{\text{sv}}} \frac{4\pi}{H_c^2} \right)^{3/2} \frac{16\sqrt{2\pi}}{B} \hat{c}_{44}^{1/2} c_{66}^{3/2} (1 + T/T_{\text{dp}})^3 \right].$$

Again, the critical current can be written as $J_c \sim \exp\{-a r^{3/2} [1 + T/T_{\text{dp}}^0 r^{1/2}]^3\}$, where a is practically independent of temperature near T_m. In this case we find $J_c \sim e^{-a}$ below T_m and $J_c \sim e^{-\alpha^3 a}$ as $T \to T_m^-$, yielding a sharp rise of J_c in a very narrow temperature range. This mechanism can be responsible for the small peak effect observed in untwinned single crystals. It can, however, only account for the *rise* in J_c at T_m. The subsequent drop of J_c to zero predicted by the collective pinning theory occurs above T_m, in a region where the theory does not apply. Furthermore, the first order transition is in practice smeared out by sample inhomogeneities, and one can never access the region where the function r is so small that $T_m/T_{\rm dp}^0\sqrt{r}$, or $r \leq 0.01$ [15]. For these reasons we associate the observed *drop* in J_c with the onset of plastic motion of vortices discussed at the end of the paper.

We now consider collective pinning in a sample with a single family of twin boundaries of mean separation d spanning the z-y plane. The field is in the z direction and a current J is applied normal to the twins. This is the experimental geometry where the peak in J_c is strongest [9]. We define two different pinning lengths, $R_{c\parallel}$ and $R_{c\perp}$, corresponding to the size of the vortex bundle in the directions parallel and transverse to the twin planes, respectively. If $R_{c\perp} \gg d$, pinning occurs via the collective action of many twin planes. Each twin is described as a sheet with a large concentration of point defects. The correlator of the random potential is given by $\Gamma(\mathbf{r},\mathbf{r}') = \gamma_1 g(|x - x'|)\delta(y - y')\delta(z - z')$, where $\gamma_1 \approx (U_0\xi^3)^2 n_p^{(2)}[1 + \mathcal{O}(\xi/d)]$ is proportional to the *areal* density $n_p^{(2)}$ of pins on each twin plane, and $g(x)$ describes correlations in the distribution of twin planes. On distances large compared to the twin spacing d the twins are essentially uncorrelated $[g(x) \approx (1/d)\delta(x)]$ and the pinning force correlator is $\tilde{W}_{ij}(\mathbf{q} + \mathbf{q}') = W_T\delta_{ij}(2\pi)^3\delta^{(3)}(\mathbf{q} + \mathbf{q}')$, with $W_T = \gamma_1/\xi^4 d a_0^2$. Collective pinning in this regime is very similar to collective pinning by point defects in bulk. The dependence of J_c on the elastic constants and temperature is identical to that obtained for isotropic point disorder, with the replacement $W \to W_T$. A peak effect in densely twinned samples may then in principle arise from the same mechanism discussed above for untwinned crystals. On the other hand, the mean squared pinning force W_T is still determined by the effective *volume* density of pins, which is now given by $n_p^{(2)}a_0/\xi d$. As a result, the anisotropy due to the twin planes increases the pinning volume of a factor $(d/d_p)^3$, with $d_p \sim (n_p)^{-1/3}$, correspondingly decreasing the critical current. For this reason collective pinning in twinned samples is very weak, especially if $d \gg a_0$, and cannot account for the observed critical currents.

The dominant pinning mechanism in twinned crystals, particularly in sparsely twinned samples, is the strong pinning of individual vortex lines on the twin boundaries. As T_m is approached from below intervortex, interactions weaken and the vortices on the twins become more strongly pinned than those in the channels between twins. The main contribution to the critical current arises then from pinning of single vortices on the twins, and particularly from those vortex segments that are strongly pinned in rare regions with an excess of impurities. As described below, it is the rise in the fraction of such strongly pinned vortex segments on the twins with T that gives a peak in J_c in twinned samples.

To evaluate the critical current due to strong pinning in regions with excess impurities, we consider a representative vortex line trapped near a twin plane by the large concentration of point defects on the twin and interacting with its neighbors at an average distance a_0 in the lattice. The remainder of the lattice, even though not directly pinned by the twin, is held in place by interactions. The magnitude of the elastic force associated with displacing a length L of the representative fluxon a transverse distance u from its equilibrium position is

$$F_{\rm el}(u,L) \sim \tilde{\epsilon}_1\frac{u}{L} + c_{66}uL\,. \tag{3}$$

The first term is the force associated with tilting the representative vortex, with $\tilde{\epsilon}_1$ the tilt energy per unit length. The second term arises from the interaction with the neighbors. The typical pinning force exerted on a vortex segment of length L is $[\overline{F_p^2(L)}]^{1/2} \sim (W_1\xi^2 L)^{1/2}$, with $W_1 = \gamma_1/\xi^5 a_0^2$ the mean squared pinning force per unit volume due to a single twin of thickness ξ. The most effective pinning arises from rare regions with an anomalously large impurity concentration that pin strongly the vortex segment. The pinning forces F_p in these regions exceed the typical pinning force $[\overline{F_p^2(L)}]^{1/2}$ and give the dominant contribution to J_c. The problem of strong pinning of vortex lines is analogous to that of incommensurate charge density waves and can be rigorously discussed following Ref. [16]. Here we prefer, however, to follow the more phenomenological, but physically intuitive, discussion given by Coppersmith [17]. The condition for the strong pinning is $F_p(\xi) > F_{\rm el}(\xi)$ [4,18]. The critical current J_c is proportional to the density n of the strongly pinned vortex segments where this condition is satisfied. To find it we note that the pinning force F_p scales as the impurity excess in the region, and it can be shown to be Gaussian distributed with variance $(\overline{F_p^2})^{1/2}$ [17]. The density n of vortex segments that are strongly pinned in these excess-impurity regions is then

$$n \sim \int_0^\infty dL \int_{F_{\rm el}(L)}^\infty dF_p e^{-(F_p)^2/2[\overline{F_p^2(L)}]}, \tag{4}$$

where we have used $u \sim \xi$. Using Eq. (3), one can show that the integral over L is dominated by the length scale $L^* \sim \sqrt{\tilde{\epsilon}_1/c_{66}}$, where the single-vortex tilting force balances the force of interaction of the pinned vortex with the rest of the lattice. The critical current J_c is then

$$J_c \sim n \sim \exp[-(c_{66}/c_{66}^*)^{3/2}], \tag{5}$$

where $c_{66}^* = (2W_1/\sqrt{\tilde{\epsilon}_1})^{2/3}$. As T approaches T_m, the shear modulus softens and drops to zero. Correspondingly, J_c grows according to Eq. (5). The result given in Eq. (5) applies for $c_{66} \geq c_{66}^*$. The mechanism of strong pinning just described is also operative in untwinned samples where vortices are pinned by isotropic point disorder. We note that strong pinning of isolated vortices yields the same functional dependence of J_c on c_{66} as collective pinning of small bundles. In both cases $J_c \sim \exp(-ac_{66}^{3/2})$.

In untwinned crystals a is essentially the same for these two pinning mechanisms. In twinned crystals collective pinning of vortex bundles is very weak ($a \sim 1/W$) and strong single-vortex pinning in regions with excess impurities ($a \sim 1/W_1$) controls the critical current even when the condition for single-vortex pinning is satisfied only locally. In this case the main contribution to J_c arises from pinning energy barriers which are large compared to the typical barrier $E_p(L^*) \sim (W_1\xi^2/c_{66}^*)(c_{66}^*/c_{66})^{1/4}$, but small compared to the scale of the elastic energy of interaction of the vortex segment with the rest of the lattice $c_{66}\xi L^* \sim \sqrt{\tilde{\epsilon}_1 c_{66}}\,\xi \sim (W_1\xi^2/c_{66}^*)(c_{66}/c_{66}^*)^{1/2}$.

Most experiments do not measure the critical current as defined theoretically, but rather the nonlinear resistivity. For comparison with experiments it is important to discuss the small finite resistivity due to the creep of strongly pinned vortex segments at low measuring currents, $J < J_c$. To find the energy barriers that determine the creep rate, we consider the distribution of barriers separating different metastable states in the regime of strong pinning. As discussed above for pinning forces, it has been shown that the impurity energies δE of strongly pinned vortex segments of length L are described by a Gaussian distribution $\sim \exp[-(\delta E)^2/2W_1L\xi^4]$ [17]. A gain δE in pinning energy is associated with a cost $\delta\mathcal{F}_{\rm el} = \tilde{\epsilon}_1\xi^2/L + c_{66}\xi^2L$ in elastic energy. The resulting activation barrier is $E_p = \delta E - \delta\mathcal{F}_{\rm el}$. The distribution function of creep barriers can be evaluated in the following: [17],

$$P(E_p) \sim \int_0^\infty dL e^{-(E_p + \delta\mathcal{F}_{\rm el})^2/2W_1L\xi^4} \sim e^{-2E_pc_{66}/W_1\xi^2}. \tag{6}$$

The typical creep barrier at $J \sim J_c$ is $E_p \sim W_1\xi^2/c_{66}$. At arbitrary currents the creep barrier becomes $E_p(J) \sim W_1\xi^2/c_{66}f(J/J_c)$. The function $f(J/J_c)$ decreases as J/J_c grows, but its explicit form cannot be obtained by our dimensional analysis. The resulting thermally activated resistivity is

$$\rho \sim \frac{1}{J_c}\exp\left(-\frac{W_1\xi^2 f(J/J_c)}{Tc_{66}}\right), \tag{7}$$

with J_c given by Eq. (5). The temperature dependence of the resistivity at a fixed current is governed by the temperature dependence of J_c. When c_{66} decreases near melting, the creep resistivity drops exponentially. This result applies for $J \leq J_c$ and temperatures near, but below, melting, where c_{66} is not too small, $c_{66} > \sqrt{W_1}\,\xi$.

In the regime just described dominated by strong pinning of individual vortex segments on twin planes, interactions are still strong enough to hold the lattice together so that the remainder of the flux array can be described as an elastic continuum. On the other hand, in the presence of an applied current the competition between strong single-vortex pinning at the twin boundaries and the elastic deformations of the portions of flux lattice between twin planes eventually leads to the development of large strains in the regions next to the twins. This results in the breakup of the lattice and the onset of plastic flow. An approximate condition for the onset of plastic flow can be obtained by equating the total Lorentz force per unit area on the flux lattice in the channel between two twins $\sim f_L d$, with $f_L = BJ/c$ the Lorentz force per unit volume parallel to the twins, to the force due to elastic stresses $\sim c_{66}(u/R^2)R$ in a region of linear size $R \sim a_0$, for $u \sim a_0$. This defines a current scale $J_p \sim cc_{66}/Bd$. The onset of plastic flow at $J \sim J_p$ corresponds to a sharp rise in the differential resistivity and a drop in the critical current.

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[1] G. Blatter *et al.,* Rev. Mod. Phys. **66**, 1125 (1994).
[2] H.R. Kerchner, J. Low Temp. Phys. **50**, 337 (1983).
[3] A.B. Pippard, Philos. Mag. **19**, 217 (1969).
[4] A.I. Larkin and Yu.N. Ovchinnikov, J. Low Temp. Phys. **34**, 409 (1979).
[5] X.S. Ling and J. Budnick, in *Magnetic Susceptibility of Superconductors and Other Spin Systems,* edited by R.A. Hein, T.L. Francavilla, and D.H. Liebenberg (Plenum, New York, 1991), pp. 377-388.
[6] G. D'Anna *et al.,* Europhys. Lett. **25**, 225 (1994).
[7] W.K. Kwok (private communication).
[8] N.-C. Yeh (private communication).
[9] W.K. Kwok *et al.,* Phys. Rev. Lett. **73**, 2614 (1994).
[10] S. Bhattacharya and M.J. Higgins, Phys. Rev. Lett. **70**, 2617 (1993).
[11] W.R. White, A. Kapitulnik, and M.R. Beasley, Phys. Rev. B **50**, 6303 (1994).
[12] M.V. Feigel'man *et al.,* Phys. Rev. Lett. **63**, 2303 (1989).
[13] M.V. Feigel'man and V.M. Vinokur, Phys. Rev. B **41**, 8986 (1990).
[14] Even though only the long wavelength tilt modulus enters in this expression, $T_{\rm dp}$ is mainly determined by $q_\perp \simeq a_0^{-1}$.
[15] For the same reason we neglect the thermal softening of the pinning potential in our subsequent discussion of strong pinning by twin planes.
[16] V.M. Vinokur, M.B. Mineev, and M.V. Feigel'man, Sov. Phys. JETP **54**, 1138 (1981).
[17] S.N. Coppersmith, Phys. Rev. B **44**, 2887 (1991).
[18] R. Labusch, Crystal Latt. Def. **1**, 1 (1969).