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Quantum tunneling of flux lines in single-crystal $Bi_2Sr_2CaCu_2O_8$ with columnar defects

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We have studied the relaxation rate of the magnetization of a single crystal of $Bi_2Sr_2CaCu_2O_8$ before and after heavy-ion irradiation as a function of temperature for fields less than 1 T. We find a significant decrease in the thermally activated regime, and a moderate shift to lower values in the quantum regime after irradiation [respectively $S(0)\approx 0.022$ and $S(0)\approx 0.15$].

I. INTRODUCTION

The observation of anomalously high decay rates of the magnetization in a granular high-T_c superconductor was made by Mota *et al.*[1,2] and interpreted, on the basis of phase slippage in Josephson junctions, as the occurrence of quantum tunneling.[3] It was shown later that this phenomenon is not restricted to granular materials and also occurs in high quality $YBa_2Cu_3O_7$,[4] organic superconductors[5] and $Tl_2CaBa_2Cu_2O_8$ single crystals.[6] Recent calculations, extending the classical collective weak-pinning theory to the quantum regime, provide expressions for the relaxation rate at $T=0$ (Refs. 7 and 8) which have been compared favorably with experimental data.[5,7] We have performed in this study relaxation measurements at temperatures down to 150 mK on $Bi_2Sr_2CaCu_2O_8$ single crystals where well-defined columnar pinning defects were deliberately introduced. As demonstrated already, these defects dramatically enhance the pinning properties of the material in the thermally activated regime.[9] We show in the following that the quantum regime is only weakly affected by the presence of these defects.

II. EXPERIMENTAL PROCEDURES

We have carried out magnetic relaxation measurements on single-crystal $Bi_2Sr_2CaCu_2O_8$ ($T_c=86.5$ K) from 150 mK to 15 K. The sample had dimensions $625\times750\times15$ μm^3, with the c axis along the short dimension; it was prepared by the floating zone method.[10] After performing the relaxation experiments on this sample, the crystal was irradiated at room temperature using a 5.3-GeV Pb-ion beam along the c axis. As demonstrated by Hardy co-workers,[11,9] these irradiated crystals are pierced by continuous amorphous ion tracks along the c direction which can act as strong pinning centers in the flux creep regime, for fields along the irradiation direction. The total fluence was $F=10^{11}$ ions/cm^2, corresponding to an average distance between tracks ≈ 300 Å. The experiments were repeated on this irradiated sample.

The magnetization measurements were carried out using a Hall probe technique similar to the one in Ref. 4. A miniature InSb Hall probe of approximate dimensions 0.2×0.5 mm^2, with wire contacts on the ends and sides for the current and multiple voltage leads, respectively, was sandwiched between two micaplates with the sample mounted with vacuum grease on top, centered above the pairs of Hall contacts with the c axis parallel to the applied field. The Hall voltage that measures the field sensed by the probe, H_H, was given by an ac bridge using currents in the range 10 μA. We first calibrated the probe in the range of temperature used for our experiments by measuring the Hall voltage for external fields with no sample mounted on the probe. The sample and measuring system, with several thermometers (ruthenium oxide, carbon glass), were mounted on the end tip of a dilution refrigerator; for the low-temperature range, the temperature could be stabilized between 0.1 and 4 K; alternatively, the system could be operated as a standard ^{4}He cryostat with exchange gas and heater for temperatures above 4.2 K. We assumed that the field measured by the probe can be expressed as $H_H=H_{\rm exp}+\alpha M$, where $H_{\rm exp}$ is the external field and M is the total magnetization of the sample. This is only valid provided that each portion of a uniformly magnetized sample contributes equally to the local field on the probe. This could be reasonably assumed here, as the size of the Hall contacts were comparable to the sample dimensions. The calibration factor α was determined from a separable magnetization measurement in a superconducting-quantum-interference-device susceptometer. We recorded magnetic hysteresis curves $M_{\rm irr}(H)$ at each temperature in order to measure the irreversible component of the magnetization, using the equation $M_{\rm irr}=\alpha^{-1}(H_H^i-H_H^d)/2$ where H_H^i (respectively H_H^d) is the magnetic field sensed by the probe for increasing (decreasing) field. This is necessary in order to obtain the normalized relaxation rate

$S=(1/M_{irr})dM/d\ln t$ which is discussed in the following sections. The procedure used for the relaxation measurements was the following: at each temperature, the field was first cycled up to a maximum field H_H and then down to the measuring field H_{exp}. In order to achieve a full critical state in the sample, the cycling field was set larger than $2H^*+H_{exp}$, where H^* is Bean's critical field for full flux line penetration.[12] (H^* was obtained from the magnetic hysteresis loops as the point where the virgin magnetization curve merges with the cycled curve.) The cycling field $H_M=3$ T was found to satisfy this condition for the whole temperature range. With these conditions, more than one track is available on average for each flux line. After the field was ramped down to H_{exp}, we measured the relaxation of the magnetization over typically 2000 sec. The first point was taken 3 sec after reaching the measuring field. We have checked in a first run with no sample mounted that the setup did not bring any measurable relaxation of the field sensed by the probe.

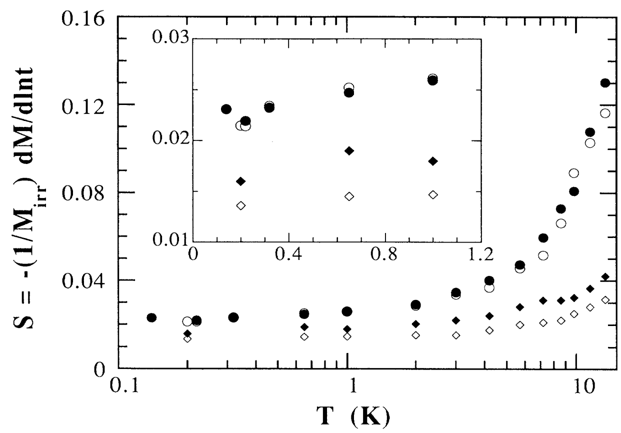

FIG. 2. Normalized relaxation rate of the irreversible magnetization. Circles: as-grown crystal; squares: irradiated crystal; empty: $H=0.2$ T; filled: $H=0.5$ T. Inset: enlargement of the low-temperature data.

III. RESULTS

Figure 1 shows magnetization-versus-time data at different temperatures. The dependence in time is roughly logarithmic and we obtain values of $dM/d\ln t$ by taking the tangent. However, the slope depends slightly on time and we observed a weak decrease of the relaxation rate as the relaxation proceeds. A similar behavior was reported by Van der Beek *et al.* and was attributed to the increase of the effective pinning barrier with the decrease of the screening current intensity.[13] For that reason we always took the slope at the same arbitrary time $t_i=180$ sec. Figure 2 shows the normalized relaxation rate as a function of the temperature. The screening current $J(T,H)$ can be obtained from $M_{irr}(T,H)$, using the extension of Bean's model to disk samples[14] giving (assuming thickness $\ll$ radius) $M_{irr}\approx Jt/3$, where t is the sample thickness. However, the applied magnetic field used for our relaxation experiments was always less than H^* and the above relation—which underestimates the critical current—should be corrected to account for demagnetizing effects.[14] Above a temperature close to 4 K, we find an exponential increase of J with decreasing temperature, characteristic of the flux creep regime.[15] However, below 4 K, the increase shows a saturation, as would be expected in the presence of a nonactivated relaxation process (Fig. 3). This low-temperature behavior is even clearer on the $S(T)$ curve (Fig. 2). $S(T)$ increases almost linearly with temperature between 2 and 15 K; but below 2 K, we measure a nearly constant value $S(0)\approx 2.2\times 10^{-2}$, comparable to the value found in Ref. 4. The field dependencies of $J(T)$ and $S(T)$ were tentatively investigated by performing the experiments for $H_{exp}=0.2$ T and $H_{exp}=0.5$ T. As shown in Figs. 2 and 3, the critical current decreased with field, while the normalized relaxation rate remained roughly constant. An additional check of this last point was done by measuring the normalized rate for the virgin sample at $T=400$ mK, for

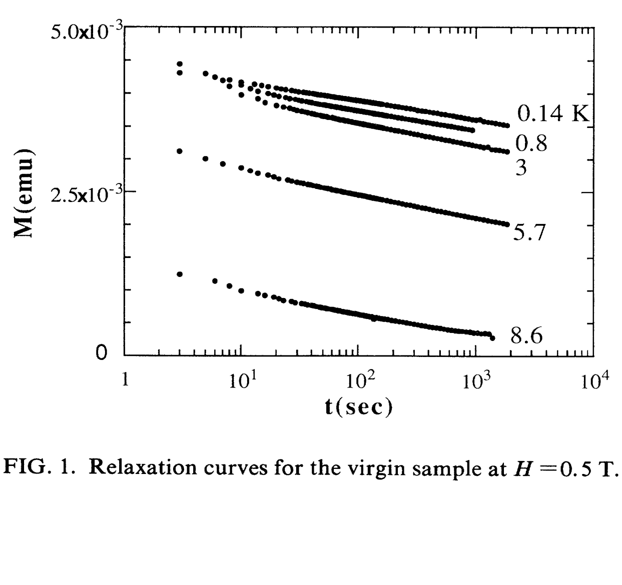

FIG. 1. Relaxation curves for the virgin sample at $H=0.5$ T.

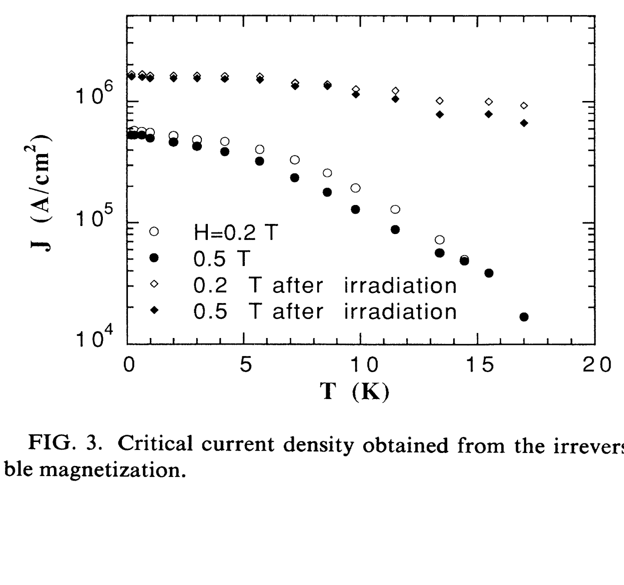

FIG. 3. Critical current density obtained from the irreversible magnetization.

fields up to 3 T: we could not detect any appreciable change of the normalized rate above 0.5 T. Thus it is unlikely that flux line curvature effects, as expected for flat samples in low fields,[14] could have been playing a role in our experiments.

After irradiation, the results were changed dramatically. In the high-temperature regime, the temperature dependence of the critical current density as well as the normalized relaxation rate were drastically reduced (Figs. 2 and 3). The field dependence for $S(T)$, however, significantly increased. In the low-temperature regime, we find a decrease of the limiting relaxation rate down to $S(0)\approx 1.5\times 10^{-2}$ and a slight increase of the field dependence of this value (Fig. 2).

IV. DISCUSSION

The observation that the magnetic relaxation rate does not vanish as $T\rightarrow 0$, as would be predicted from the classical flux creep model, is now widely attributed to the occurrence of vortex motion by quantum tunneling.[16,17,5] Blatter, Geshkenbein, and Vinokur computed the zero-temperature relaxation rate in bulk superconductors, using the macroscopic quantum tunneling theory, including dissipation, in the framework of the collective weak-pinning theory.[7] The case of the layered superconductors and the finite-temperature effects were addressed later by Blatter and Geshkenbein.[8] As found by these authors, the results of the three-dimensional (3D) anisotropic theory should be valid for the layered materials when $\theta > \epsilon$, where θ is the angle between the induction and the plane direction and $\epsilon^2 = m/M$ is the ratio of the effective masses. In the limit of single vortex pinning and strong dissipation, for pinning potential with minimal characteristic length $r_p \approx \xi$, the coherence length, and pinning wells d_p apart, the zero-temperature relaxation rate is given by[7]

$$S(0)\approx \hbar/(L_c \eta d_p^2) , \tag{1}$$

where L_c is the length of the correlated hopping vortex segment along the c axis and $\eta = \phi_0 H_{c2}/c^2\rho_n$ the flux flow resistivity, with H_{c2} the upper critical field in that direction and ρ_n the normal-state resistivity. The expression for η given above is valid assuming that the carrier mean free path is larger than ξ. In the dirty limit, the viscosity should be smaller than this value.[18] In the framework of the collective pinning theory, for screening currents close to the critical current J_c (i.e., in the limit of vanishing effective pinning energies), in the single vortex regime ($H \ll H_{c2}J_c/J_0$ where J_0 is the depairing current) and for large damping, we have[7]

$$L_c \approx \epsilon d_p (J_0/J_c)^{1/2} , \tag{2}$$

$$S(0)\approx (e^2\rho_n \xi^2/\epsilon\hbar d_p^3)(J_c/J_0)^{1/2} . \tag{3}$$

This theory has been found to agree with experimental data for the quasi-two-dimensional (BEDT-TTF) based radical salt and for $YBa_2Cu_4O_8$.[5] It is fair to notice at this point that, due to the uncertainty on the various parameters entering Eq. (3), the agreement between experimental data and theory is more qualitative than strictly quantitative. Taking $\rho_n \approx 100$ $\mu\Omega$ cm, $\xi\approx 27$ Å, $\epsilon\approx \frac{1}{50}$, $J_c\approx 5\times 10^5$ A/cm^2 (Fig. 3), $J_0 \approx 10^8$ A/cm^2, and $d_p \approx \xi$ for the virgin sample gives $S(0)\approx 0.2$. This value is one order of magnitude larger than our experimental data for the virgin sample. A naive application of (3) to the case of pinning by tracks shows that the quantum rate should be affected through the change in d_p and J_c. To make sense, d_p could not be given by the tracks separation (300 Å), which would lead to unmeasurable rates. It seems more likely that dense, secondary pinning wells are necessary to explain the substantial low-temperature rate measured for the irradiated sample. In a more general way, it is surprising that materials as different as $YBa_2Cu_3O_7$, $YBa_2Cu_4O_8$, $Bi_2Sr_2CaCu_2O_8$—irradiated or not—and BEDT-TTF based salts exhibit comparable low-temperature rates. Considering the very much different anisotropies (ϵ) for the Bi-based and Y-based materials, for instance, one would expect a strongly enhanced tunneling rate for the latter, which is indeed not observed. A more general expression for the quantum relaxation rate is given by Eq. (1). Taking $d_p \approx \xi$, the in-plane coherence length, gives

$$S(0)\approx \hbar\rho_n/(L_c\phi_0^2) . \tag{4}$$

Taking $L_c \approx s \approx 15$ Å, the plane spacing, and $\rho_n \approx 100$ Ω cm gives $S(0)\approx 10^{-2}$ as observed in our case. A qualitative agreement is found also in the case of $YBa_2Cu_3O_8$ (Ref. 4) and for the organic compound (BEDT-TTF)$_2$Cu(SCN)$_2$.[5] Although expression (4) with $L_c \approx s$ cannot account for the small variation of the relaxation rate after irradiation, it has the merit of being independent of the anisotropy, while Eq. (3) would predict a strong dependence on this parameter.

In order to check further the validity of the collective weak-pinning theory, we have analyzed our data for the virgin sample also in the thermally activated regime. A good approximation for the creep activation barrier in the single-vortex regime is[19]

$$U(J)=U_0 \ln(J_c/J) \tag{5}$$

giving, in the case of full flux penetration,

$$S(T)\approx kT/U_0 . \tag{6}$$

As demonstrated by Maley *et al.*,[20] the effective pinning potential $U(J)$ can be obtained from relaxation curves taken at different temperatures, using

$$U(J)=AT - T\ln(dM/dt) , \tag{7}$$

where A is a constant. We repeated their procedure with our data, between 5 and 15 K, for the as-grown crystal: the result is shown in Fig. 4(a) where we have used $A\approx 15$ ($A\approx 14$ was used in Ref. 20). As shown in Fig. 4(b), our data can be reasonably described by Eq. (5). In addition, we have for the highest value of J, $J\approx J_c/3$. This is consistent with the assumption of a single-vortex regime, in the limit of vanishing effective pinning energies. We now focus on the high-J regime. In the collective weak-pinning limit, U_0 is given by the elastic energy

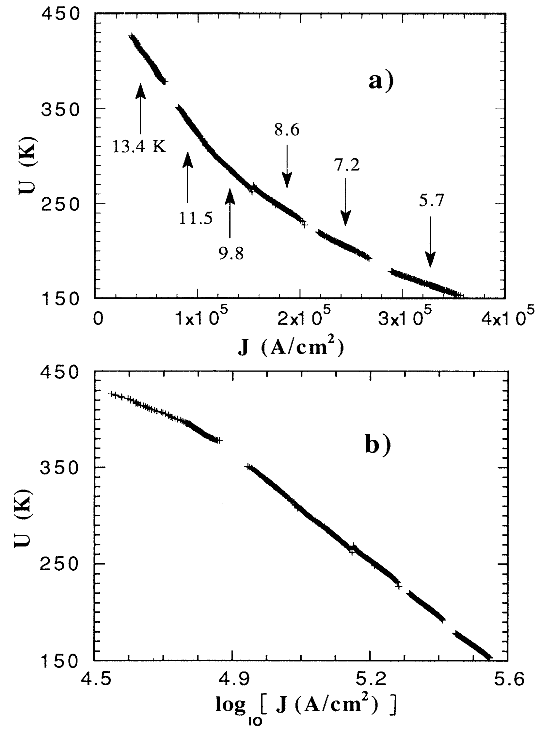

FIG. 4. Effective pinning energy given by (6) and $A=15$ for the as-grown crystal.

of the segment hopping from one site to the other:[7]

$$U_0 \approx \epsilon^2 \epsilon_l d_p^2 / L_c \quad (8)$$

where $\epsilon_l = (\phi_0/4\pi\lambda)^2$ is the line tension in the isotropic case. Using (2), we have in the limit of vanishing energies

$$d_p \approx U_0 (J_0/J_c)^{1/2} / \epsilon \epsilon_l \ ,$$
$$L_c \approx U_0 (J_0/J_c) / \epsilon_l \ . \quad (9)$$

From Fig. 4(b), we have $U_0 \approx 150$ K. Taking $\lambda \approx 2500$ Å, we obtain $d_p \approx 3000$ Å and $L_c \approx 700$ Å from (8). The incompatibility between the generally assessed value $d_p \approx \xi$ (Refs. 7 and 8) and the one found here arises from the large value for U_0. As noticed already by Vinokur, Kes, and Koshelev,[21] taking $d_p \approx \xi$ leads, in the framework of the weak 3D collective pinning theory, to $U_0 \approx 1$ K, which cannot be identified with the energy barrier for flux creep measured for the high-T_c superconductors. Then, the relevance of the 3D collective pinning theory in the case of $Bi_2Sr_2CaCu_2O_8$ with point defects (and *a fortiori* in the one of the irradiated sample) is questionable. A 2D collective approach may, in this case, be more adequate.

V. CONCLUSION

We have shown that the quantum relaxation rate in the 2:2:1:2 compound is not reduced after the introduction of columnar defects, which nevertheless substantially increase the critical current density. Nor is this rate affected by the large anisotropy of this material, as compared to previous experiments on the weakly anisotropic 1:2:3 compound.

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[1]A. C. Mota, A. Pollini, P. Visani, and K. A. Müller, and J. G. Bednorz, Phys. Scr. **37**, 823 (1988).
[2]A. C. Mota, A. Pollini, P. Visani, K. A. Müller, and J. G. Bednorz, Phys. Rev. B **36**, 4011 (1987).
[3]E. Simanek, Phys. Rev. B **39**, 11 384 (1989).
[4]L. Fruchter, A. P. Malozemoff, I. A. Campbell, and J. Sanchez, Phys. Rev. B **43**, 8709 (1991).
[5]A. C. Mota, G. Juri, P. Visani, A. Pollini, T. Teruzzi, K. Aupke, and B. Hilti, Physica C **185-189**, 343 (1991).
[6]F. Zuo, A. C. Shi, A. J. Berlinsky, H. Duan, and A. M. Hermann (unpublished).
[7]G. Blatter, V. B. Geshkenbein, and V. M. Vinokur, Phys. Rev. Lett. **66**, 3297 (1991).
[8]G. Blatter and V. Geshkenbein (unpublished).
[9]V. Hardy, J. Provost, M. Hervieu, B. Raveau, S. Durcok, E. Pollert, J. C. Frison, J. P. Chaminade, and M. Pouchard, Physica C **191**, 85 (1992).
[10]N. Motohira, T. Hasegawa, K. Kishio, and K. Kitazawa, J. Ceram. Soc. Jpn. **97**, 1009 (1989).
[11]V. Hardy, J. Provost, D. Groult, B. Raveau, and S. Bouffard, Nucl. Instrum. Methods B **54**, 472 (1991).
[12]C. P. Bean, Rev. Mod. Phys. **36**, 31 (1964).
[13]C. J. van der Beek, P. H. Kes, M. P. Maley, M. J. V. Menken, and A. A. Menovsky, Physica C **185-189**, 2507 (1991).
[14]L. W. Conner and A. P. Malozemoff, Phys. Rev. B **43**, 402 (1991).
[15]V. N. Zavaritsky and N. V. Zavaritsky, Physica C **185-189**, 2141 (1991).
[16]A. V. Mitin, Zh. Eksp. Teor. Fiz. **93**, 590 (1987) [Sov. Phys. JETP **66**, 335 (1987)].
[17]A. Hamzic, L. Fruchter, and I. A. Campbell, Nature **345**, 515 (1990).
[18]J. Bardeen and M. J. Stephen, Phys. Rev. **140**, A1197 (1965).
[19]V. V. Vinokur, G. Blatter, M. V. Feigel'man, V. B. Geshkenbein, and A. I. Larkin, Physica C **185-189**, 276 (1991).
[20]M. P. Maley, J. O. Willis, H. Lessure, and M. E. McHenry, Phys. Rev. B **42**, 2639 (1990).
[21]V. M. Vinokur, P. H. Kes, and A. E. Koshelev, Physica C **168**, 29 (1990).