Scaling in the angular dependence of the critical current and temperature-dependent anisotropy ratio in $Bi_2Sr_2CaCu_2O_8$

R. Fastampa, S. Sarti, and E. Silva
Dipartimento di Fisica, Università "La Sapienza," Piazzale Aldo Moro 2, 00185 Roma, Italy

E. Milani
Dipartimento di Ingegneria Meccanica, Università di Roma "Tor Vergata," V.O. Raimondo 8, 00173, Roma, Italy

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We present measurements of the critical current in an epitaxial $Bi_2Sr_2CaCu_2O_8$ film as a function of the magnetic field, the temperature, and of the angle ϑ between the field and the a,b planes. The orientational study reveals an increasing anisotropy in J_c with increasing temperature. The isothermal curves $J_c(H)$ taken at various angles $0° \le \vartheta \le 90°$ are found to collapse to a single curve when the magnetic field is normalized to a scaling function $f(\vartheta)$: $J_c(H,\vartheta)=J_c[H/f(\vartheta)]$. The best fit for $f(\vartheta)$ is obtained using an expression originally derived for $H_{c2}(T,\vartheta)$ by Tinkham. The increase in $J_c(0°)/J_c(90°)$ with increasing temperature is reflected by the increase in the anisotropy ratio $f(0°)/f(90°)$, consistent with the quasi-two-dimensional Tinkham expression. Consistently, the same model quantitatively describes the increasing anisotropy with increasing temperature.

I. INTRODUCTION

One of the most striking phenomena in the transport properties on highly anisotropic high-temperature superconductors in a magnetic field is the behavior of the observed in-plane resistivity ρ as a function of the angle ϑ, between the applied magnetic field **H** and the CuO planes, and the angle α between **H** and the current density **J**. Some of these phenomena are particularly evident in $Bi_2Sr_2Ca_2CuO_8$, and they undermine the traditional picture of fluxons moving under the opposite influence of a Lorentz-like driving force and of some pinning mechanism due to structural imperfections: First, there is no dependence of the observed resistivity[1,2] and of the critical current[3] on the angle α, contrary to what one could expect if the driving force were Lorentz-like. Further, the resistivity does not vanish if the field is parallel to the current. Second, columnar defects produced by heavy ion irradiation, that should give rise to an enhancement of the pinning force when the field is aligned parallel to the columns, seem to be ineffective: the ϑ behavior of the critical current, in fact, shows no additional peak after irradiation.[4] Finally, in a small angular range near $\vartheta=0°$ ($\Delta\vartheta=1°$), an unexpected increase of the dissipation has been observed, giving rise to a sharp peak in the resistivity when measured at high current densities[5] and in the microwave surface resistance.[6]

All these puzzling features suggested the need of a more subtle fluxon statics and dynamics, possibly connected to the high anisotropy of these materials. In particular, it has been proposed that the flux lines could penetrate the sample in a stepwise manner: fluxons would then be made up of pancake vortices lying in the planes[7,8] and Josephson fluxons parallel to the planes themselves, the last being locked to the interplane region and thus ineffective for the dissipation (intrinsic pinning model).[9,10] The predictions of this model agree with the J_c experiments only at low temperatures [namely, 4.2 K (Refs. 3 and 11)], while at higher temperatures the experimental framework is still controversial.[3,11]

A different approach is given by considering the dimensionality of the superconductor itself, instead of analyzing the detailed structure of the flux lines. If we think of the $Bi_2Sr_2Ca_2CuO_8$ sample as composed of a stack of superconducting and nonsuperconducting layers, with thicknesses d and s, respectively (this scenario seems to suit well several known properties of $Bi_2Sr_2Ca_2CuO_8$), due to the very short out-of-plane coherence length $\xi_\perp$ one can have a dimensional crossover in the superconducting material, with the following meaning: If $d<\xi_\perp(0)<s/\sqrt{2}$, each superconducting layer behaves like a quasi-two-dimensional (2D) superconductor;[12] approaching T_c, $\xi_\perp(t)$ can overcome $s/\sqrt{2}$; the layers become interacting, and the superconductor is better described by a coupled-layers model, resulting in a 3D, anisotropic behavior.[13] Such a crossover, experimentally observed in $Bi_2Sr_2Ca_2CuO_8$ close to T_c at very low current densities as depending on the temperature[14,15] and on the magnetic field,[16,17] is a peculiar feature of layered superconductors, and it could be responsible for several physical properties of $Bi_2Sr_2Ca_2CuO_8$.

In this paper we present a set of measurements of the critical current density J_c in an epitaxial $Bi_2Sr_2Ca_2CuO_8$ film, as a function of the external magnetic field H, applied at variable angles ϑ, in the range 40 K $<T<$ 80 K. The experimental curves, $J_c(H,\vartheta)$, are found to exhibit scaling features: It turns out that $J_c(H,\vartheta)$

$=J_c[H/f(\vartheta)]$, where $f(\vartheta)$ is a (temperature-dependent) point-defined scaling function. In other words, each curve collapses onto the orthogonal-field curve, $J_c(H,90°)$, if it is plotted as a function of $H/f(\vartheta)$. The results clearly show strong deviations from the $\sin\vartheta$ scaling, thus ruling out the intrinsic pinning description in this range of temperature. Moreover, the obtained scaling function is fully described by a quasi-2D model in which the magnetic field has to be scaled with ϑ according to the Tinkham law, written originally for the upper critical field H_{c2}:[18]

$$\frac{|H_{c2}(\vartheta)\sin(\vartheta)|}{H_{c2\perp}}+\left[\frac{H_{c2}(\vartheta)\cos(\vartheta)}{H_{c2\parallel}}\right]^2=1 . \quad (1)$$

In addition, the resulting anisotropy ratio $f(0°)/f(90°)$ is found to grow with the temperature; it is shown that this behavior is fully consistent with the Tinkham-like scaling rule, giving internal consistency to the explanation proposed. To explain the results only general arguments on the dimensionality are needed, and no specific model based on specific forces acting on fluxons has to be invoked.

II. EXPERIMENTAL SECTION

The sample is a Bi-2:2:1:2 film, epitaxially grown on a (001) oriented $NdGaO_3$ substrate by liquid phase epitaxy. The mosaic spread is less than 0.15°.[19] The film thickness is about 1 μm. The film was patterned by scratching in a 2-mm-long, 100-μm-wide strip. A standard four-probe configuration was obtained by gluing test leads with silver paint. Contacts were annealed for 1 h in air atmosphere at a temperature of 500 °C. The resistive transition measured at low current density (≈ 1 A/cm^2) showed metallic behavior followed by a sharp transition ($\Delta T_c=3$ K, 10–90 % criterion) with zero resistance (within the experimental resolution) at 82.4 K. The resistivity at room temperature is 210 $\mu\Omega$ cm, and the extrapolated residual resistivity at zero temperature is about one tenth of that at room temperature.

The critical current measurements were performed by a stepped current source and a nanovoltmeter. A voltage criterion of 2 μV/mm was chosen (no significant variation on J_c was found at 41 and 79 K by halving the voltage criterion). The sample holder was aligned with the a,b plane of the film lying on the vertical plane. The magnetic field (up to 2T) was supplied by a traditional electromagnet, that could rotate in a horizontal plane. The rotation angle ϑ between the magnetic field and the basal plane was accurate to 0.1°, that is comparable to the width of the mosaic spread, so that no smearing of $J_c(\vartheta)$ curves should be due to the misorientation of the film. The current was directed along the a,b plane, and made the same angle ϑ with the field. The insensitivity of the J_c measurements with respect to the relative orientation of the current and the field was checked, according with the literature.[1,2] Special care was taken in the alignment for the parallel orientation. The temperature was measured by a platinum sensor, and it ranged between 40 and 80 K. Stabilization was within ±0.03 K during each measurement. At every temperature and angle, measurements of the critical current were taken at several fields.

III. EXPERIMENTAL RESULTS

In Fig. 1 the temperature behavior of J_c for $H=0$ and for $H=2$ T, applied parallel ($\vartheta=0°$), nearly parallel ($\vartheta=1°$), and perpendicular ($\vartheta=90°$) to the basal planes, is shown. The value of the critical current density at 77 K is about 1.5×10^3 A/cm^2, while at low temperatures it tends to the value of 3×10^4 A/cm^2. A perpendicular applied field is shown to induce a strong depression in the critical current density. Therefore, we extended the high-temperature measurements down to 1 mT to obtain a reasonably detailed field dependence of J_c. From the figure, two main qualitative features appear. First of all, the critical current at $\vartheta=0°$ is depressed with respect to that at $\vartheta=1°$, from low temperature up to 65 K: this is equivalent to the anomalous enhancement of the dissipation in parallel field orientation found with different techniques.[5,6] Second, the anisotropy in the critical current *increases* with increasing temperature.

The anomalous behavior of the dissipation at $\vartheta=0°$ is reflected in the J_c measurements at low angles, presented in Fig. 2. As can be seen, the unexpected depression of the critical current density is more evident at low temperatures and disappears in the measurement at 71 K. At the present stage, a quantitative explanation of this anomalous dissipation peak at $\vartheta=0°$ is not available, and the effect seems to be sample dependent (for example, no such peak was found in J_c in Refs. 3 and 20); therefore, we will limit ourselves to the discussion of the data in the angular range that is not affected by the presence of the peak (the anomalous dissipation exists in a very narrow angular range—$\Delta\vartheta\approx0.5°$—so that for our measurements, only the data points at $\vartheta=0°$ and possibly at $\vartheta=0.5°$ should be neglected).

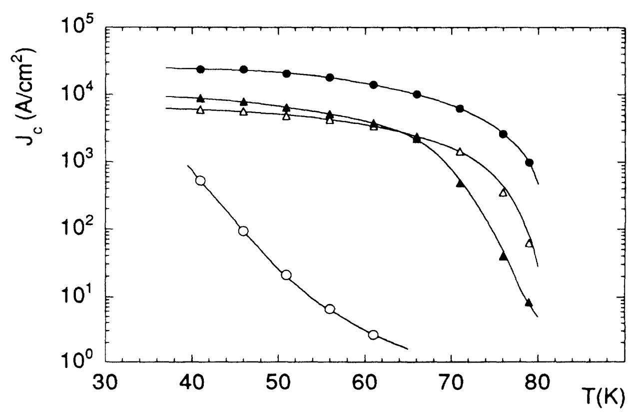

FIG. 1. Temperature dependence of J_c in zero field (full circles) and at 2 T for different orientations: orthogonal field ($\vartheta=90°$), open circles; parallel field ($\vartheta=0°$), empty triangles; slightly misaligned ($\vartheta=1°$), full triangles. Continuous lines are guides for the eye. Anomalously, the critical current at $\vartheta=0°$ is *lower* than at $\vartheta=1°$ up to ≈65 K. The orthogonal field is shown to induce an abrupt fall in J_c.

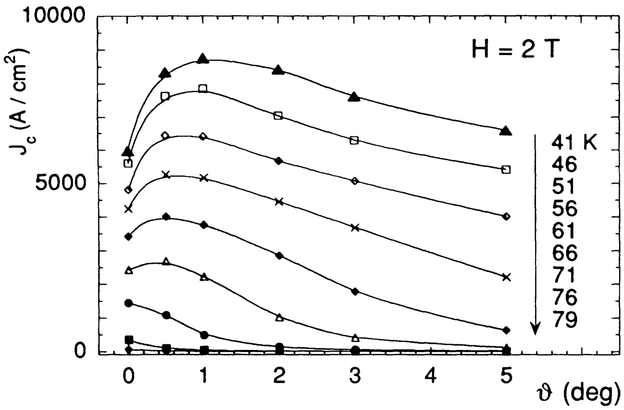

FIG. 2. Low-angle measurements of J_c at 2 T, for several temperatures. The anomalous dip (increased dissipation) visible at $\vartheta=0°$ is found to disappear by raising the temperature. Continuous lines are guides for the eye.

To show the increasing anisotropy with increasing temperature, in Fig. 3 some of the angular measurements of the critical current density normalized to the value in orthogonal field, at several temperatures, are reported. We use here a logarithmic scale for the angle, to enhance the behavior close to the parallel orientation (this choice does not allow us to display the data points at $\vartheta=0°$). The measurements were found to be symmetrical with respect to $\vartheta=0°$, so that no information is lost with displaying only the data for $\vartheta>0°$. The data clearly show that the anisotropy of J_c in a fixed applied field increases with increasing temperature and increasing field, ranging from ~ 2 at 41 K and 70 mT, up to ~ 2000 at 71 K and 0.7 T. We stress that the opposite behavior is present in $YBa_2Cu_3O_7$: the anisotropy in the critical current density *decreases* with increasing temperature.[21]

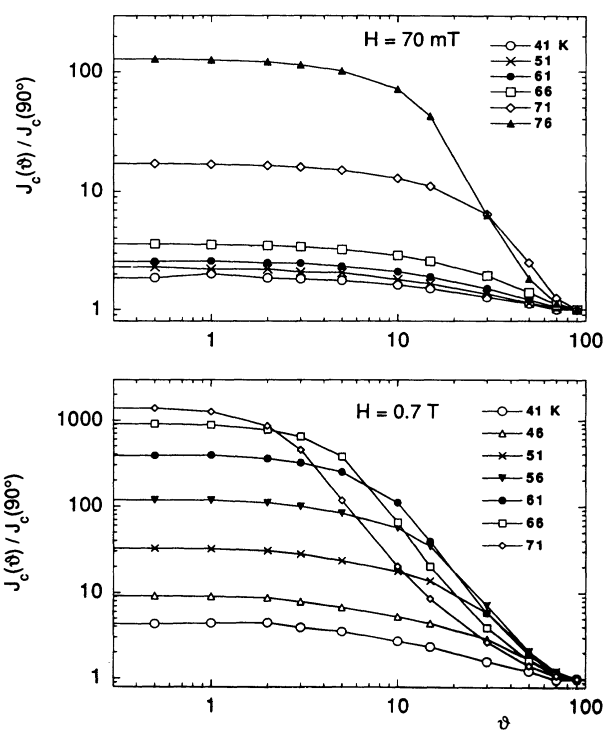

FIG. 3. Angular measurements for J_c normalized to the $\vartheta=90°$ values, at different fixed fields and different temperatures. The anisotropy in J_c increases with increasing temperature at low as well as high fields. The ϑ scale is logarithmic to enhance the low-angles behavior. Continuous lines are guides for the eye.

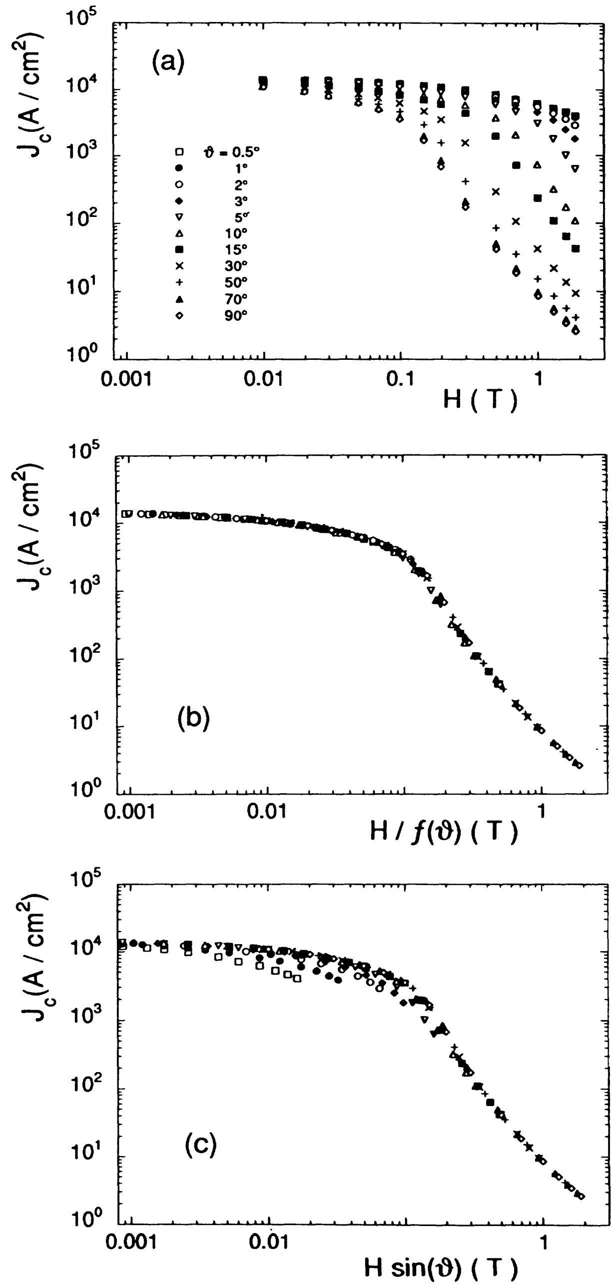

FIG. 4. Typical scaling procedure. (a) Field dependence of J_c for different angles, raw data; (b) scaled data; each curve $J_c(H)$ scales over the $J_c(H_\perp)$ curve when plotted as J_c vs $H/f(\vartheta)$ with an appropriate parameter $f(\vartheta)$; (c) $\sin\vartheta$ scaling: curves are plotted as J_c vs $H\sin\vartheta$. This scaling breaks down at relatively large angles ($\vartheta \approx 5°$). Data are taken at $T=61$ K.

In Fig. 4(a) we report a typical measurement of J_c at fixed temperature, as a function of H for several values of ϑ. All the curves at a given temperature appear to be translated with respect to each other along the x axis, when plotted on a logarithmic scale. This strongly suggests the definition of a parameter $f(\vartheta)$ such that $J_c(H,\vartheta)=J_c[H/f(\vartheta)]$. This is the basis of the scaling procedure described in the following section.

IV. SCALING PROCEDURE AND DISCUSSION

Before describing in detail the scaling procedure and the informations obtained from it, some general considerations are needed. Several attempts have been made to describe the field dependence of the critical current density in $Bi_2Sr_2Ca_2CuO_8$, most of these being based on the analysis of the so-called volume pinning force.[3,22] The agreement with the well-known result by Kramer[23] $F_p \propto (H/H_{c2})^{1/2}(1-H/H_{c2})^2$ is quantitative, if a different characteristic field H^* replaces H_{c2}. However, since the Kramer's expression for the pinning force is obtained from the balance between the pinning force and the Lorentz force acting on a flux line, and the dissipation in $Bi_2Sr_2Ca_2CuO_8$ has been shown to be independent on the latter,[1-3] we think that the experimental frame is not fully consistent with this kind of approach. For these reasons, in the discussion of the data, we will not focus on the volume pinning force, but instead we will work directly on the $J_c(H)$ curves, taken at different angles ϑ, showing that they scale on the $J_c(H_\perp)$ curve with a temperature-dependent scaling function, and we will identify such function with the one corresponding to Eq. (1). While a scaling procedure has been employed in the past for J_c measurements,[3,24] up to now no detailed determination of the scaling function and of its temperature dependence has been given.

In Fig. 4(b) we plot the $J_c(H)$ curves at various ϑ versus $H/f(\vartheta)$, where $f(\vartheta)$ is the scaling parameter necessary to make the curves collapse on the curve $J_c(H,\vartheta=90°)$. For comparison, in (c) the same curves are plotted as a function of $H\sin\vartheta$, that is assuming that only the transverse component of the field is effective. It is immediately apparent that the simple $\sin\vartheta$ scaling undergoes a breakdown for $\vartheta \approx 5°$ [this result is in agreement with microwave measurements on samples from the same batch (Refs. 6 and 25)]. We stress that the breakdown of the $\sin\vartheta$ scaling appears at angles much larger than the mosaic spread, and also much larger than the width of the dip at $\vartheta=0°$, so that we can attribute it to intrinsic properties. The breakdown in the $\sin\vartheta$ scaling involves that also the parallel component of the applied magnetic field causes a dissipation. In Fig. 5 collapsed curves at several temperatures are shown. The data at $\vartheta=0°$ were not scaled due to the presence of the anomalous dip in J_c, but this lack of data does not prevent us from obtaining interesting results.

From the scaling procedure, we then obtain at each temperature a point-defined function $f(\vartheta)$, to be compared with theoretical predictions. In Fig. 6 we report the scaling functions for several temperatures (a log axis

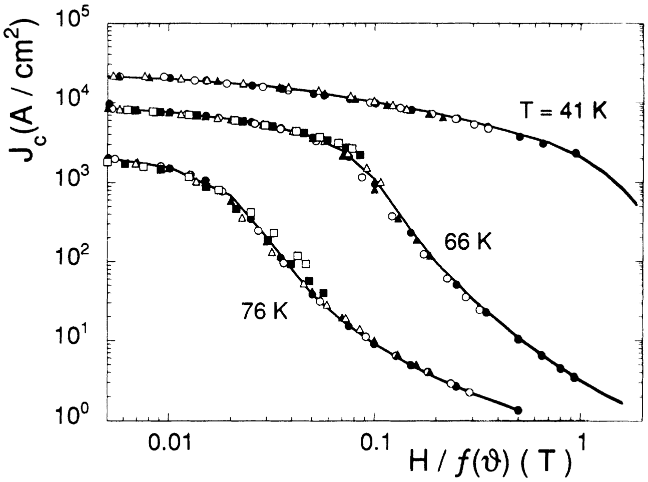

FIG. 5. Collapsed J_c vs $H/f(\vartheta)$ at several temperatures. The presented data are at $\vartheta=30°$ (full circles), $\vartheta=10°$ (open circles), $\vartheta=5°$ (full triangles), $\vartheta=2°$ (open triangles), $\vartheta=1°$ (full squares), and $\vartheta=0.5°$ (open squares). The scaling curves, $J_c(H,90°)$, are plotted with continuous lines.

for the angle has been chosen to make the low-angles behavior clear). By construction, the point at $\vartheta=90°$ is always 1; from this plot, the fact that the anisotropy increases with increasing temperature assumes a quantitative aspect.

Due to the failure of the $\sin\vartheta$ scaling, a different theoretical curve is needed for $f(\vartheta)$. Alternative scaling functions are suggested by the work of Hao and Clem:[26] they demonstrated that the Gibbs free-energy difference between the superconducting and normal state depends on H only through the ratio $H/H_{c2}(\vartheta)$; moreover, from the fact that the anisotropic properties of $Bi_2Sr_2CaCu_2O_8$ do not depend on the relative direction between the field and a transport current, but only on the direction between the field and the crystallographic axes, they inferred that also nonequilibrium quantities such as the resistivity and the critical current should depend on H/H_{c2} only. This result allows us to use the anisotropic behavior of the critical field as an angular scaling function, as will be done below.

The best known models that led to explicit analytical functions for $H_{c2}(\vartheta)$ are the Tinkham model for a single slab of finite thickness,[18] Eq. (1), and the Lawrence-Doniach (LD) model for Josephson-coupled layers.[13] The latter model gives the result

$$\left[\frac{H_{c2}(\vartheta)\sin\vartheta}{H_{c2\perp}}\right]^2+\left[\frac{H_{c2}(\vartheta)\cos\vartheta}{H_{c2\parallel}}\right]^2=1 \tag{2}$$

with the same meaning of the symbols as in Eq. (1). The scaling functions are obtained from Eqs. (1) and (2), written in the form

$$H_{c2}(\vartheta)=H_{c2\perp}f(\vartheta) .$$

The resulting scaling functions are, respectively,

$$f_T(\vartheta)=\frac{1}{2(H_{c2\perp}/H_{c2\parallel})^2\cos^2\vartheta}\times[\sqrt{\sin^2\vartheta+4(H_{c2\perp}/H_{c2\parallel})^2\cos^2\vartheta}-|\sin\vartheta|]=\frac{1}{2\varepsilon^{-2}\cos^2\vartheta}[\sqrt{\sin^2\vartheta+4\varepsilon^{-2}\cos^2\vartheta}-|\sin\vartheta|] \quad (3)$$

and

$$f_{LD}(\vartheta)=\frac{1}{\sqrt{\sin^2\vartheta+(H_{c2\perp}/H_{c2\parallel})^2\cos^2\vartheta}}=\frac{1}{\sqrt{\sin^2\vartheta+\varepsilon^{-2}\cos^2\vartheta}}\ , \quad (4)$$

where ε is the anisotropy ratio.

A striking difference between the two models lies in the temperature dependence of the inverse anisotropy ratio, defined as $1/\varepsilon=H_{c2\perp}/H_{c2\parallel}$: in the LD model both $H_{c2\parallel}$ and $H_{c2\perp}$ have the same temperature dependence, so that ε does not depend on T, while in the quasi-2D scenario a temperature dependence appears. In this model, one has $H_{2c\perp}=\Phi_0/2\pi\xi_\parallel\xi_\perp\sim(1-T/T_c)$, but the size of the fluxon core in parallel field is limited by the thickness d of the layer,[18] and $H_{c2\parallel}=\Phi_0/2\pi d\xi_\parallel\sim(1-T/T_c)^{1/2}$. As a consequence, the resulting anisotropy ratio depends on T:

$$1/\varepsilon=H_{c2\perp}/H_{c2\parallel}=d/\xi_\parallel(T)=[d/\xi_\parallel(0)](1-T/T_c)^{1/2}\ . \quad (5)$$

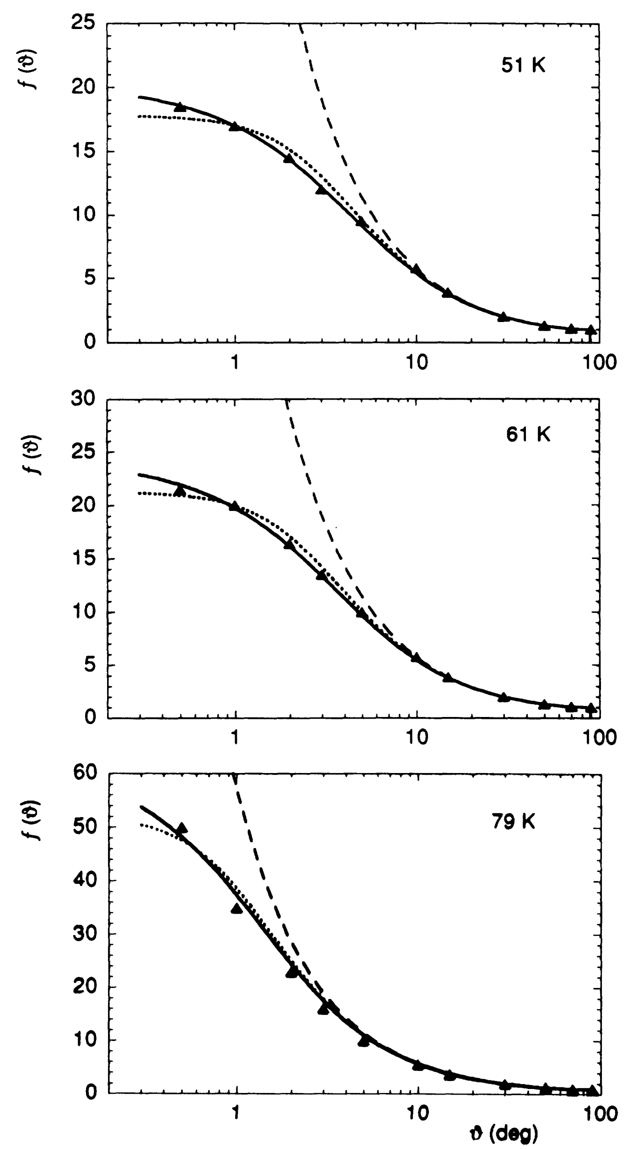

FIG. 6. Scaling function $f(\vartheta)$ at different temperatures. Full triangles are obtained from the scaled J_c curves; solid and dotted lines are least-squares fits by Eqs. (3) and (4), respectively. Equation (3) fits slightly better than Eq. (4) at all the temperatures. The breakdown of the simple $|\sin\vartheta|$ scaling (dashed lines) appears at angles as large as 5°.

In Fig. 6 we compare our $f(\vartheta)$ with Eqs. (1) and (2), taken as scaling functions; the parallel field is left as the only fitting parameter. Even if the Tinkham formula (continuous lines) gives slightly better fits than the LD one, this result alone could not allow one to decide against the latter model (similar uncertainty was reported in the scaling of the resistivity in Ref. 27, where only two temperatures were investigated). However, the analysis of the temperature dependence of the anisotropy ratio ε obtained from the fits gives a clear result: as is shown in Fig. 7 (where we report $1/\varepsilon^2$ vs T), a temperature dependence is evident, and, as a consequence, the LD model is ruled out. Moreover, the resulting $1/\varepsilon^2$ is found to lie on a straight line, that is $1/\varepsilon$ follows the square-root law given by Eq. (5) for a stack of isolated layers of finite thickness. The Tinkham formula describes our scaling results, including the temperature dependence of the scaling function.

The slope extracted from the $1/\varepsilon^2$ vs T plot by means of a linear fit, allows for some speculation: in principle, knowing the thickness d, one could extract $\xi_\parallel(0)$. As a crude estimate, we can take $d\approx3$ Å, that is the thickness of the double-CuO layers. Our data, for this choice of d, give $\xi_\parallel(0)\approx37$ Å, close enough to the commonly accepted value $\xi_\parallel(0)\approx25$ Å. The critical temperature from the same fit is $T_c=82.25$ K, extremely close to the measured zero-resistance temperature, $T_{R=0}=82.4$ K. For a precise determination of the true Ginzburg-Landau coherence length $\xi_\parallel(0)$ one should perform upper critical field measurements.

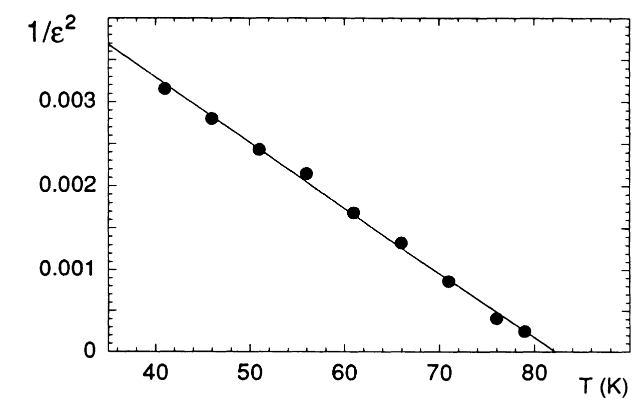

FIG. 7. $1/\varepsilon^2$ as a function of T (full points). The temperature dependence of the anisotropy ratio is very well described by the quasi-2D model [from Eq. (5), solid line].

Finally, we note that the dimensional crossover from a quasi-2D to a 3D situation obtained by further increasing the temperature[14,15,17] could not be observed by the study of the critical currents: the strong depression of J_c in perpendicular field does not allow a detailed study at temperatures close to T_c, and different techniques are needed.

V. CONCLUSIONS

We presented measurements of the critical current density in an epitaxial $Bi_2Sr_2CaCu_2O_8$ film, as a function of the external magnetic field for several field orientations and temperatures. The magnetic-field curves of J_c at different angles are found to collapse on the orthogonal-field curve; the scaling law is found to be the quasi-2D Tinkham formula, Eq. (3); the extracted anisotropy ratio is temperature dependent, and it increases with increasing temperature. This finding is consistent with the Tinkham's predictions; in particular the square-root law (5) is found to describe well the data. A coupled-layers scaling function (Lawrence and Doniach model) could describe the J_c scaling reasonably well, but is unable to explain the temperature-dependent anisotropy ratio.

[1] Y. Iye, S. Nakamura, and T. Tamegai, Physica C **159**, 433 (1989).

[2] T. Fukami, T. Kamura, T. Yamamoto, and S. Mase, Physica C **160**, 391 (1989).

[3] S. Labdi, H. Raffy, O. Laborde, and P. Monceau, Physica C **197**, 274 (1992).

[4] J. R. Thompson, Y. R. Sun, H. R. Kerchner, D. K. Christen, B. C. Sales, B. C. Chakoumakos, A. D. Marwick, L. Civale, and J. O. Thomson, Appl. Phys. Lett. **60**, 2306 (1992).

[5] Y. Iye, T. Tamegai, and S. Nakamura, Physica C **174**, 227 (1991).

[6] R. Fastampa, M. Giura, R. Marcon, and E. Silva, Europhys. Lett. **18**, 75 (1992).

[7] P. H. Kes, J. Aarts, V. M. Vinokur, and C. J. van der Beek, Phys. Rev. Lett. **64**, 1063 (1990).

[8] J. R. Clem and M. W. Coffey, Phys. Rev. B **42**, 6209 (1990); J. R. Clem, *ibid.* **43**, 7837 (1991).

[9] M. Tachiki and S. Takahashi, Solid State Commun. **70**, 291 (1989).

[10] M. Tachiki and S. Takahashi, Solid State Commun. **72**, 1083 (1989).

[11] T. Fukami, K. Miyoshi, T. Nishizaki, Y. Horie, F. Ichikawa, and T. Aomine, Physica C **202**, 167 (1992).

[12] L. N. Bulaevskii, V. L. Ginzburg, and A.A. Sobyanin, Zh. Eksp. Teor. Fiz. **94**, 355 (1988) [Sov. Phys. JETP **68**, 1499 (1988)]; L. N. Bulaevskii, Zh. Eksp. Teor. Fiz. **64**, 2241 (1973) [Sov. Phys. JETP **37**, 1133 (1973)].

[13] W. E. Lawrence and S. Doniach, in *Proceedings of the 12th International Conference on Low Temperature Physics, Kyoto, Japan, 1970,* edited by E. Kanda (Kiegaku, Tokyo, 1971), p. 361.

[14] R. Fastampa, M. Giura, R. Marcon, and E. Silva, Phys. Rev. Lett. **67**, 1795 (1991).

[15] R. Marcon, E. Silva, R. Fastampa, and M. Giura, Phys. Rev. B **46**, 3612 (1992).

[16] E. Silva, R. Marcon, R. Fastampa, M. Giura, and S. Sarti, Physica C **214**, 175 (1993).

[17] S. Sarti, E. Silva, R. Fastampa, M. Giura, and R. Marcon, Phys. Rev. B **49**, 556 (1992).

[18] M. Tinkham, Phys. Rev. **129**, 2413 (1963).

[19] G. Balestrino, V. Foglietti, M. Marinelli, E. Milani, A. Paoletti, and G. Paroli, Appl. Phys. Lett. **57**, 2359 (1990).

[20] P. Schmitt, P. Kummeth, L. Schultz, and G. Saemann-Ischenko, Phys. Rev. Lett. **67**, 267 (1991).

[21] L. Schultz, B. Roas, P. Schmitt, P. Kummeth, and G. Seamann-Ischenko, IEEE Trans. Magn. **27**, 990 (1991).

[22] R. A. Rose, S. B. Ota, P. A. J. de Groot, and B. Jayaram, Physica C **170**, 51 (1990); H. Yamasaki, K. Endo, S. Kosaka, M. Umeda, S. Yoshida, and K. Kajimura, Phys. Rev. Lett. **70**, 333 (1993).

[23] E. J. Kramer, J. Appl. Phys. **44**, 1360 (1973).

[24] G. Jakob, M. Schmitt, Th. Kluge, C. Tomé-Rosa, P. Wagner, Th. Hahn, and H. Adrian, Phys. Rev. B **47**, 12 099 (1993).

[25] E. Silva, M. Giura, R. Marcon, R. Fastampa, G. Balestrino, M. Marinelli, and E. Milani, Phys. Rev. B **45**, 12 566 (1992).

[26] Z. Hao and J. R. Clem, Phys. Rev. B **46**, 5853 (1992).

[27] H. Raffy, S. Labdi, O. Laborde, and P. Monceau, Phys. Rev. Lett. **66**, 2515 (1991).