Vortex-glass transition in the $(K,Ba)BiO_3$ cubic superconductor

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We present a detailed characterization of the vortex-glass transition in the cubic $(K,Ba)BiO_3$ superconductor (with critical temperature $T_c \sim 30$ K). Combining transport measurements down to 0.3 K and third-harmonic susceptibility close to T_c we show that the vortex-glass transition field can be well fitted by a power law: $H_g \sim (1-T/T_c)^n$ with $n \sim 1.5$ over three magnetic field decades. In striking contrast with cuprates, this simple behavior holds down to our lowest temperature. The E-J characteristics are consistent with the vortex-glass transition for both $J \| H$ and $J \perp H$ configurations assuming that, close to T_g, the vortex line is making an average angle $\sim 20°$ with the external field. The vortex-glass correlation length reaches several μm in the vicinity of the transition, showing that vortices are correlated over more than $200a_0$ (where a_0 is the intervortex spacing). The characteristic exponents $\nu = 1.0 \pm 0.2$ and $z = 5.0 \pm 0.6$ are field independent and do not depend on the relative orientation of the magnetic field and the current density. [S0163-1829(98)04942-X]

I. INTRODUCTION

It is now well established that both high T_c (Ref. 1) and low T_c (Ref. 2) superconductors may present a vortex solid to liquid phase transition. This transition line lies very close to the upper critical field in conventional low-T_c materials whereas the liquid phase can occupy a large part of the H-T phase diagram in highly anisotropic materials such as Bi-Sr-Ca-Cu-O. It has then become common to fit this transition line using a simple $(1-T/T_c)^n$ power law. It has been shown both theoretically[3] and experimentally[4] that an exponent $n=2$ can be used to described the first-order melting of a ''pure'' vortex lattice close to T_c. On the other hand, the irreversibility line of a large variety of cuprates presents a power law variation with $n \sim 3/2$ at high temperature followed by a more rapid variation at lower temperatures.[5–8] Various models have been invoked to explain this crossover: (i) a crossover from a weakly correlated to a strongly correlated rigid lattice,[9] (ii) a transition from three-dimensional (3D) vortex lines to 2D independent pancakes, [8–10] and (iii) a crossover from 3D XY to lowest-Landau-level (LLL) fluctuation regimes.[11]

In this context, the $(K,Ba)BiO_3$ superconductor ($T_c \sim 30$ K) appears to be particularly interesting since it presents a fully isotropic ($\gamma = 1$) cubic structure and a phase diagram which is experimentally accessible down to the lowest temperatures [the upper critical field $H_{c2}(0) < 30$ T]. We present here magnetotransport (down to $T = 0.3$ K and up to $H = 27$ T) and third-harmonic susceptibility measurements performed on $(K,Ba)BiO_3$ single crystals. The existence of a vortex-glass transition[12] in this system has been previously discussed (up to 20 T) in Ref. 13 and the present paper is thus organized as follows. The most characteristic features of the transition (E-J curves, magnetic field dependence of the resistance in the Ohmic regime) are briefly reviewed in Sec. II. In contrast to cuprates we will show in Sec. III that the vortex-glass transition field deduced from both third harmonic susceptibility and magnetotransport measurements can be well fitted using a simple $H_g \sim (1-T/T_c)^{1.5}$ law down to 0.3 K, i.e., over three magnetic field decades (a small deviation from this simple behavior can be observed in sample A—see below). Finally, taking advantage of the cubic structure of our system, we will compare in Sec. IV transport properties performed for different orientations of the current density J vs magnetic field: $J \perp H$, $\widehat{(J,H)} = 45°$ and $J \| H$. We will show that the transition temperature as well as the exponents does not depend on the relative orientation of the current versus magnetic field as expected for isotropic pinning by weak point disorder. We will show that the vortex-glass correlation length ξ_{vg} is very large in the vicinity of the transition line, reaching several μm. Finally Sec. V contains a summary of the main results of our work.

II. VORTEX-GLASS TRANSITION

The experiments have been performed on dark blue single crystals grown by electocrystallization.[15] The transport measurements have been done on sample A ($T_c \sim 24.2$ K), sample C ($T_c \sim 22.1$ K), and sample D ($T_c \sim 31.2$ K) and the susceptibility measurements on samples A and B ($T_c \sim 30.8$ K). Samples A and D presented very sharp transitions in zero-field transport measurements whereas samples B and C presented somehow larger transitions (1–2 K), in-

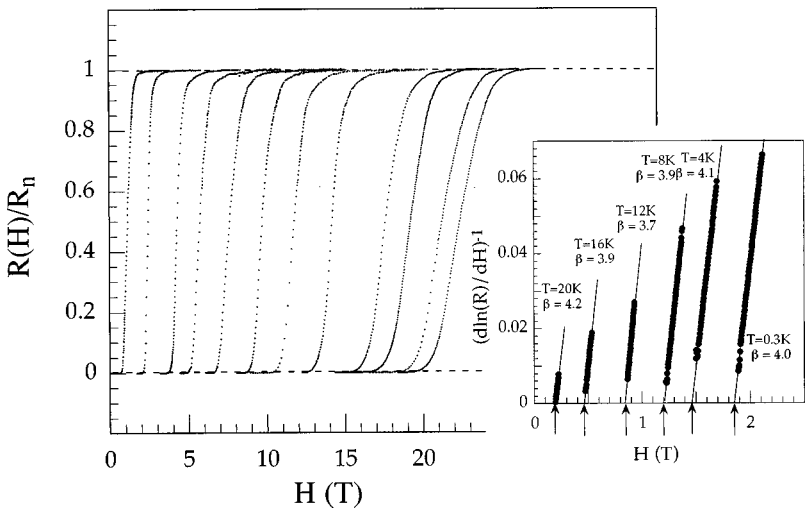

FIG. 1. Magnetoresistance measurements of a (K,Ba)BiO_3 single crystal (sample A) from the left to right T (K) = 22-20-18-16-14-12-10-8-4.2-3.0-1.5-0.3. In the inset $[d(\ln R)/dH]^{-1}$ vs magnetic field showing that $\beta = \nu(z-1)$ (i.e., the slope of the straight lines in the vortex-glass transition model) is field independent down to 0.3 K; H_g can be deduced from the zero intercept of the lines (as indicated by the arrows).

dicating the presence of larger sample inhomogeneities. As the magnetic field is increased, all samples present a smooth decrease of the resistivity without any sign of a first-order transition (i.e., without any sharp jump in the resistivity). Transport data for sample A up to 27 T are presented on Fig. 1 down to 0.3 K.

The vortex-glass transition scaling theory appeared to be particularly well adapted to describe the freezing of the vortex liquid in presence of quenched disorder.[14] As expected for this transition, for sufficiently low current densities (see below), the electrical field (E) versus current density (J) characteristics become Ohmic above the transition and the corresponding resistance is then expected to scale as[12]

$$R \sim [1 - T/T_g(H)]^{\nu(z-1)} \sim [1 - H/H_g(T)]^{\nu(z-1)}. \quad (1)$$

From Eq.(1), one gets $(d\ln R/dH)^{-1} = (H - H_g)/\nu(z-1)$; $\nu(z-1)$ and H_g can thus be obtained by plotting $(d\ln R/dR)^{-1}$ vs H (see the inset of Fig. 1, sample A). For all measured samples we obtained $\nu(z-1) = 3.9 \pm 0.3$ which does not depend on H down to the lowest temperatures.

As the current density is increased, the E-J characteristics are changing from an Ohmic behavior dominated by the thermal fluctuations to a critical (i.e., power law) behavior above a characteristic current $J^+(T)$ (see, for instance, Fig. 2 at $H = 1000$ G, $J \| H$). At $T = T_g$ the critical regime is observed for all J with $E \sim J^{(z+1)/2}$ and finally a change in the curvature is observed as the temperature is lowered below T_g [$E \sim \exp(1/J^{\mu})$ below T_g for small J]. z can thus be extracted from the power law behavior at T_g and one gets $z = 5.0 \pm 0.6$ and hence $\nu = 1.0 \pm 0.2$ up to 8 T (i.e., our highest field for the E-J characteristics). Those values are consistent with the one expected from the vortex-glass theory and are very similar to those obtained in Y-Ba-Cu-O samples.[14]

We have shown in a previous paper[16] that the upper critical field extracted from tunneling measurements lies very close to the top of the resistive transition curves. As shown on Fig. 1, the width of the transition, i.e., the liquid phase, is thus rather large in our system: $H_g - H_{c2} \sim 5$ T $\sim H_{c2}/3$ around $T_c/2$. Surprisingly, this transition remains very large down to the lowest temperatures (~ 6 T at 0.3 K). A similar behavior has been previously observed in superconducting films[17] and attributed to the presence of quantum fluctuations (see discussion below). Note, however, that torque magnetometry measurements are indicating that the irreversibility line may actually be rising up very fast below 1 K.[18]

The origin of this rather large liquid phase in the (K,Ba)BiO_3 system is, however, not straightforward. Indeed the strength of the thermal fluctuations is given by the Ginzburg number G_i,[1] which is pretty small in this 3D system: $G_i \sim 4 \times 10^{-5}$ ($G_i \sim 10^{-2}$–1 in anisotropic cuprates). Following Ref. 3, the transition line would be expected to scale as c_L^4/G_i (where c_L is the Lindemann number) and, taking $c_L \sim 0.15$, the liquid phase should only exist on a very narrow temperature range close to T_c. However, a vortex solid-liquid transition has also been observed in other isotropic systems with similar G_i values such as Rb_3C_{60} (Ref. 19) or YNi_2B_2C (Ref. 20). The origin of this melting line has then been attributed to a small value of the Lindemann criterion

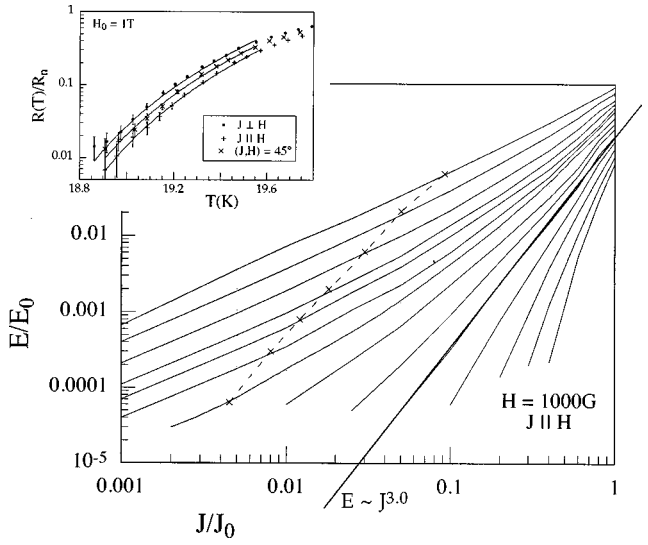

FIG. 2. E-J characteristic of a (K,Ba)BiO_3 single crystal (sample A) at $H = 1000$ G for $J \| H$ [from top to bottom T (K) = 22-21.5-21-20.9-20.8-20.7-20.6-20.5-20.4-20.3-20.2-20.1-20.0]. The solid line is an $E \sim J^{3.0}$ fit to the data at 20.4 K (i.e., $T = T_g$) as expected from the vortex-glass transition model. The dotted line marks the onset of nonlinearity in the characteristics. In the inset: enlargement of the foot of the resistive transition ($H_0 = 1$ T) for different orientation of the current density vs magnetic field. Solid lines are $(T/T_g - 1)^{\nu(z-1)}$ fits to the data with $T_g = 18.45 \pm 0.05$ K and $\nu(z-1) = 4.0 \pm 0.1$ for all J-H configurations.

as predicted by Brandt for 3D systems.[21] Alternatively, it has been shown by Blatter and Ivlev[23] that the width of the liquid phase can be significantly increased by the presence of quantum fluctuations or by the suppression of the order parameter close to H_{c2} (see also Ref. 24) and the melting field H_m is then expected to scale as

$$H_m/H_{c2} \sim \frac{4.\theta^2}{(1+\sqrt{1+4S\theta T_c/T})^2}, \qquad (2)$$

with $\theta = c_L^2(\beta_{th}/G_i)^{1/2}(T_c/T-1)$ $(\beta_{th}\sim 2)$ and $S = q + c_L^2(\beta_{th}/G_i)^{1/2}$ where q is quantifying the strength of the quantum fluctuations and $c_L^2(\beta_{th}/G_i)^{1/2}$ has been introduced to take into account the proximity of the H_{c2} line. Note, however, that Eq. (2) has been derived for the melting of a pure system without any static disorder. This expression thus cannot be directly used in our system for which no sign of any first-order melting has been observed. It has been shown, for instance, that weak point disorder can actually increase the vortex line wandering and shift the transition line downwards.[22] Nevertheless, in the absence of any quantitative expression for the position of the melting line in the presence of weak pinning centers, we will use Eq. (2) in order to get an estimation of the position of this line. The width of our liquid phase ($H_m/H_{c2}\sim 0.3$ at $T\sim T_c/2$) would then be consistent with Eq. (2) assuming that $S\sim 7$.

Taking $c_L\sim 0.15$ and $G_i\sim 4\times 10^{-5}$, we actually get directly $c_L^2(\beta_{th}/G_i)^{1/2}\sim 7$; however, quantum fluctuations (i.e., q) may also play a significant role in (K,Ba)BiO_3. Indeed, q is given by[23] $q\sim 0.15(Q/\sqrt{G_i})\Omega\tau$, where $Q = (e^2/\hbar)(\rho_n/\xi)$ (ρ_n is the normal state resistivity and ξ the coherence length, ~ 30 Å at low temperature), Ω is a cutoff frequency of the order of $c/\lambda\sim 2\times 10^{15}$ Hz in (K,Ba)BiO_3 (λ is the penetration depth, ~ 1500 Å at low temperature), and τ an effective relaxation time of the order of the elastic scattering time. Taking the carrier density[25] $n\sim 10^{22}$ cm^{-3} we get $q>1$ [$q\sim 1/n$ and can be pretty large in (K,Ba)BiO_3 due to the proximity of a metal-insulator transition, i.e., small n]. Even though this numerical estimation should be taken with care given the rather large error bars on each of the parameters, it shows that quantum fluctuations may actually be pretty large in our system. The origin of the liquid phase in the (K,Ba)BiO_3 system is thus rather complex and a complete description of this melting would probably have to take into account (i) the presence of large amount of pinning centers increasing the vortex wandering, (ii) the reduction of the c_L coefficient in a 3D system, (iii) the proximity of the H_{c2} line, and (iv) possible large quantum fluctuations.

III. TEMPERATURE DEPENDENCE OF THE MELTING FIELD

Let us now discuss the temperature dependence of the vortex-glass transition field. At high field and low temperature $H_g(T)$ has been deduced from our transport experiments as explained above ($R\sim[1-H/H_g(T)]^{\nu(z-1)}$). Close to T_c, T_g has been determined by measuring the onset of the amplitude of the third-harmonic susceptibility which marks the onset of nonlinearity[26] in the magnetic response. T_{onset} is expected to tend towards T_g as the frequency of the modulation field tends towards zero and we have thus used a low frequency $f=8.9$ Hz in our experiments. T_{onset} can be clearly identified since the onset is very sharp (see Fig. 3, sample A) and does not depend on the amplitude of the modulation field. The vortex-glass transition line deduced from both transport and third-harmonic susceptibility is presented on Fig. 4. As shown, $H_g(T)$ can be well described using a simple $(1-T/T_c)^{1.5}$ law on the entire field range (samples B, C, and D), i.e., on three decades, although small deviations can be observed for sample A. In this sample, the exponent is slightly varying ranging from ~ 1.6 close to T_c to ~ 1.3 at low temperature.

An exponent $n=1.5$ was obtained by Müller *et al.*[5] in Ba-La-Cu-O samples and the irreversibility line has thus been called the ''quasi Almeida-Thouless'' line by comparison with spin glasses. An explanation for this behavior was

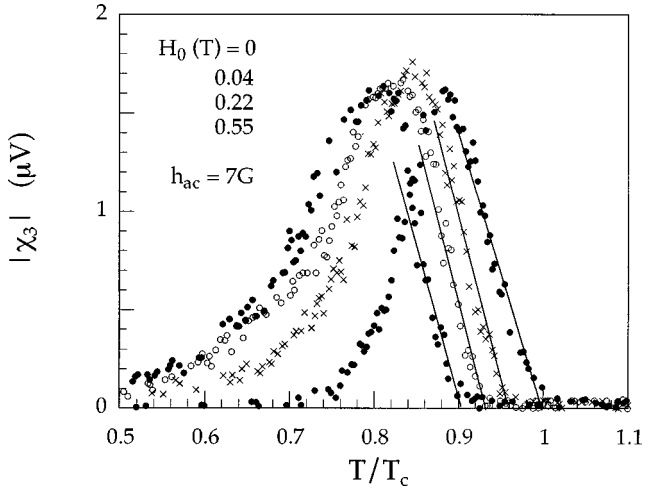

FIG. 3. Amplitude of the third-harmonic susceptibility $|\chi_3|$ for a (K,Ba)BiO_3 single crystal (sample A); the linear extrapolation to zero marks the onset of nonlinearity in the magnetic response, i.e., the vortex-glass transition temperature (in the low-frequency limit).

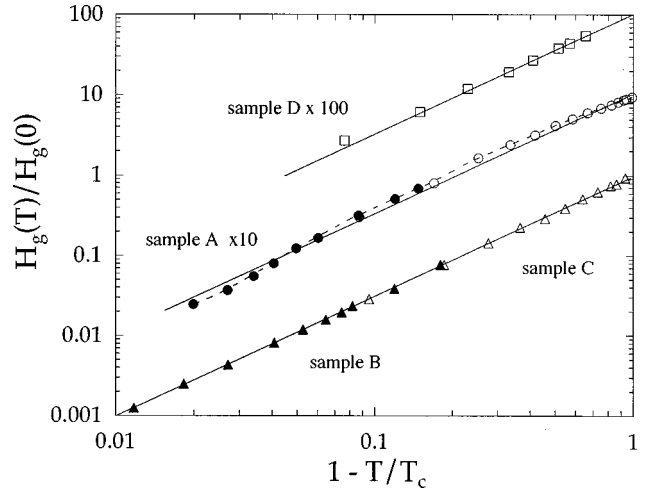

FIG. 4. Temperature dependence of the vortex-glass transition temperature deduced from transport (open symbols) and third-harmonic susceptibility measurements (solid symbols). $H_g(T)$ can be well described by a simple power law $H_g(T)\sim(1-T/T_c)^{1.5}$ in the entire temperature range (solid lines). Small deviations can be observed for sample A (the dotted line is a guide to the eye).

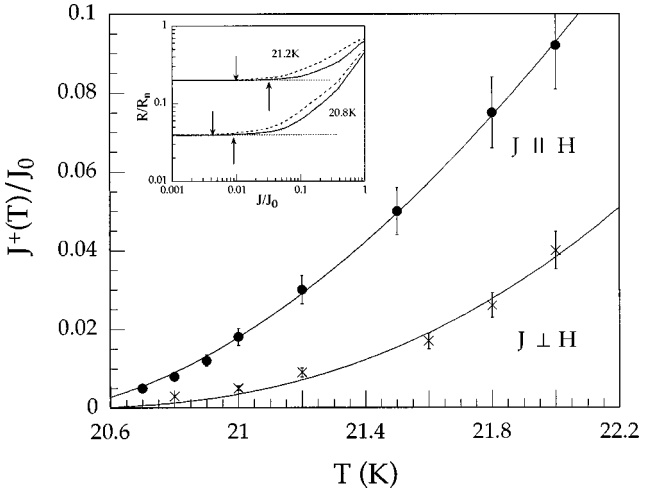

FIG. 5. Temperature dependence of the current $J^+(T)$ separating the Ohmic from the power law behaviors in the E-J characteristic as shown in the inset for $H\|J$ (solid lines) and $H\perp J$ (dotted lines) at $H_0=1000$ G for two temperatures ($J_0=8\times10^4$ A/m^2). $J^+(T)$ is indicated by the arrows in the inset (the R values for $J\|H$ have been shifted upwards for an easier comparison with $J\perp H$).

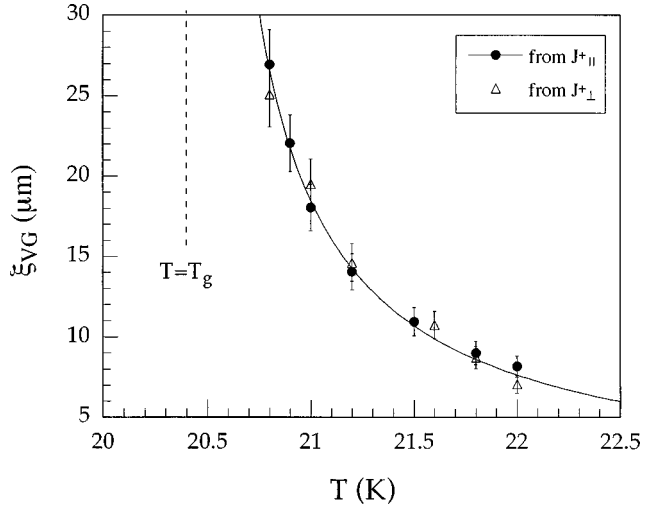

FIG. 6. Temperature dependence of the vortex-glass correlation length. The solid line is a fit to the data using a $\xi_{vg}\sim1/(T/T_g-1)^{\nu}$ law with $\nu=0.9\pm0.1$.

then proposed by Yeshurun and Malozemoff[6] who assumed that the activation energy (U_c) scales like $H_c^2a_0^2\xi\sim(1-T/T_c)^{1.5}/H$ close to T_c and suggested that the critical current may fall below the limit of detectability for $U_c\sim kT$ (H_c is the thermodynamic critical field and a_0 the intervortex spacing, $\sim1/\sqrt{H}$). Such an explanation can, however, not be used to describe our vortex-glass transition line far away from T_c. Alternatively, Blatter and Ivlev[23] suggested that the exponent $n=1.5$ may be used as an approximation of the vortex-lattice melting line for a nonzero suppression parameter S [Eq. (2)]. However, here again this approximation would only be valid close to T_c (and in any case only to describe a first-order transition line). The origin of this exponent $n=1.5$ thus remains an open question which obviously needs further theoretical investigation.

A striking feature is the absence in our system of the fast increase of $H_g(T)$ that is usually observed in cuprates below $\sim T_c/2$. This crossover has been attributed by Moloni *et al.*[11] in Y-Ba-Cu-O thin films to a change in the fluctuation regime. Indeed, they assumed that the glass fluctuations coexist with 3D XY fluctuations close to T_c whereas LLL fluctuations become predominant at lower temperatures, i.e., at higher fields. Even though the exponent $n=3/2$ is rather close to the 4/3 value expected in the 3D XY regime,[12] it is hardly believable that those fluctuations may dominate the entire phase diagram in our system. Alternatively, this crossover has been interpreted as the signature of a transition from weakly correlated vortices to rigid bundles by Almasen *et al.*[7] The transition field is expected to correspond to a vortex spacing of the order of $\lambda/8$ (or even $\lambda/4$). Taking $\lambda\sim1500$ Å we would expect to find a crossover field of the order of 2 T in clear disagreement with our experimental data.

The absence of any crossover in our system would thus rather suggest that the change of slope in the T_g vs H curve of the cuprates is related to a change in the dimensionality of the vortex system as proposed by several authors.[8–10] However, the field-independent ν and z values observed in Ref. 11 do not corroborate this explanation. The origin of this crossover temperature in cuprates thus also remains an open question.

IV. INFLUENCE OF THE ORIENTATION OF *J* vs *H*

In order to perform transport measurements for different J vs H orientations, current probes have been attached on the opposite sides of small cubic samples and two voltage probes on the top of the sample. The magnetic field (up to 3 T) has then been applied either along the J direction, at 45° or perpendicular to it. As shown in the inset of Fig. 2 ($H_0=1$ T), in the Ohmic regime, the resistivity presents a very similar temperature dependence whatever the orientation of the magnetic field. Similarly the same z value has been deduced from the E-J characteristics at T_g (see, for instance, Fig. 2 for the $J\|H$ case). The absence of any anisotropy in the scaling exponents is consistent with the vortex-glass picture, i.e., isotropic pinning induced by point disorder.

Within the vortex-glass transition model, the vortex-glass correlation length ξ_{vg} can be deduced from the crossover current $J^+(T)$ above which the E-J characteristics become non-Ohmic (see inset of Fig. 5 and dotted line in Fig. 2). $J^+(T)$ is reached when the Lorentz force energy applied on a bundle of size ξ_{vg} is of the order of the thermal energy kT and is thus given by[12]

$$J^+\phi_0\xi_{vg}^{d-1}\langle\sin(\theta)\rangle\sim kT, \tag{3}$$

where ϕ_0 is the flux quantum ($h/2e$) and θ the angle between J and the flux line. The Lorentz force energy will obviously depend on the orientation of J vs H and as expected $J^+(T)$ is smaller for $J\perp H$ than it is for $J\|H$ (see Fig. 5). The existence of a critical regime for $J\|H$ is directly related to the wandering of the vortex line which leads to a nonzero angle between the local flux density vector and the current density even in the $J\|H$ configuration. Rotational experiments[27] indicate that this average flux density vector makes an angle $\sim20°-30°$ with the external field in

(K,Ba)BiO_3 samples (at low fields). The $J_\perp^+/J_\parallel^+$ ratio is thus expected to be of the order of 2–3 in good agreement with our experimental value. As shown in Fig. 6, ξ_{vg} [deduced from Eq. (3)] rises up rapidly as the vortex-glass transition temperature is approached varying by a factor ~ 5 between 20.8 K and 22 K ($T_g \sim 20.4$ K). At our lowest temperature, ξ_{vg} is very large, ~ 30 μm$\sim 200a_0$, as already observed by Gammel *et al.*[28] in Y-Ba-Cu-O single crystals. Finally, the exponent ν can then be directly deduced from the temperature dependence of $\xi_{vg} \sim 1/(T/T_g - 1)^\nu$. We get $\nu = 0.9 \pm 0.1$ (see dotted line in Fig. 6) in good agreement with the R vs H data.

V. CONCLUSION

In summary, we have presented a detailed characterization of the vortex-glass transition in the cubic (K,Ba)BiO_3 superconductor. The transport and third-harmonic susceptibility data can be very well described by the vortex-glass transition theory with $\nu(z-1) = 3.9 \pm 0.3$ independent of the magnetic field down to 0.3 K. The vortex-glass transition line can be fitted by a simple power law $H_g \sim (1 - T/T_c)^{1.5}$ in the entire temperature range. This line does not present the fast increase of $H_g(T)$ that is usually observed in cuprates at low temperature. As expected for the glass transition in the presence of weak point disorder the characteristic exponents ν and z do not depend on the orientation of the magnetic field compared to the current. A very large (i.e., several μm) vortex-glass correlation length has been deduced from the crossover current separating the Ohmic and power law regimes in the E-J characteristics.

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