JAIST Repository
https://dspace.jaist.ac.jp/

Title	Anomalous vortex liquid-to-glass transition line in twinned Yba_2Cu_3O_〈7-δ〉
Author(s)	Naito, Tomoyuki; Iwasaki, Hideo; Nishizaki, Terukazu; Kobayashi, Norio
Citation	Physical Review B, 70(1): 014515-1-014515-6
Issue Date	2004-07-29
Type	Journal Article
Text version	publisher
URL	http://hdl.handle.net/10119/4611
Rights	T. Naito, H. Iwasaki, T. Nishizaki, and N. Kobayashi, Physical Review B , 70(1), 2004, 014515-1-014515-6. Copyright 2004 by the American Physical Society. http://link.aps.org/abstract/PRB/v70/e014515
Description	

Japan Advanced Institute of Science and Technology

Anomalous vortex liquid-to-glass transition line in twinned $YBa_2Cu_3O_{7-\delta}$

Tomoyuki Naito and Hideo Iwasaki
School of Materials Science, Japan Advanced Institute of Science and Technology (JAIST), Tatsunokuchi 923-1292, Japan

Terukazu Nishizaki and Norio Kobayashi
Institute for Materials Research, Tohoku University, Sendai 980-8577, Japan

(Received 21 November 2003; revised manuscript received 10 March 2004; published 29 July 2004)

The vortex phase transition in twinned $YBa_2Cu_3O_{7-\delta}$ single crystals is studied by measuring the resistivity under conditions in which the correlated disorders compete with point disorders. When a magnetic field is applied exactly parallel to the twin planes, we find a kink in the obtained liquid-to-glass transition line near 80 kOe and a nonuniversality in critical behavior. By introducing a matchinglike effect for planar defects, we show that both anomalous phenomena can be explained by the scenario that a mixture of Bose glass and vortex glass appears above the magnetic field of the kink, and Bose glass appears below it.

DOI: 10.1103/PhysRevB.70.014515 PACS number(s): 74.25.Fy, 74.25.Qt, 74.72.Bk

I. INTRODUCTION

It is well known that in the mixed state of high-temperature superconductors, a vortex phase is determined by the competition among three energies, i.e., vortex-vortex interaction, pinning, and thermal energies.[1] When the vortex-vortex interaction dominates the configuration of vortices, the vortex liquid solidifies into the vortex lattice via a first-order phase transition; that phase transition in clean $YBa_2Cu_3O_{7-\delta}$ (YBCO) single crystals has been established by a number of studies.[2] When the pinning energy exceeds the elastic energy of vortices (i.e., in a dirty system), the vortex solid phase changes from a lattice to a glassy phase, such as vortex glass[3] or Bose glass.[4] Fisher *et al.*[3] theoretically predicted that in the presence of randomly distributed point defects, the vortex liquid undergoes a second-order phase transition into vortex glass; this has been observed experimentally in thin films[5] and single crystals.[6] Since the contribution of point disorders on the vortex pinning is enhanced with an increasing magnetic field, this phase can be observed even in clean samples such as untwinned YBCO.[7,8] Columnar defects are randomly introduced into a crystal by irradiating heavy ions. When a magnetic field is applied parallel to the columnar defects, they accommodate the vortices and, as proposed by Nelson and Vinokur,[4] a glassy phase, Bose glass, appears; the transition between Bose glass and liquid is also second order. From observing small critical exponents in a previous study we pointed out the possibility that this phase appears in twinned YBCO.[8] Later, Grigera *et al.*,[9] using extensive scaling analysis, claimed that this transition occurs in the same materials. Many other reports support the existence of a Bose glass phase, but those experiments were performed in relatively low magnetic fields where the point disorders do not act as effective pinning centers. Maiorov and Osquiguil[10] studied the glassy phase in twinned YBCO in magnetic fields up to 180 kOe and found a crossover magnetic field $H_{cr} \sim 40$ kOe, where the angular dependence of the transition temperature abruptly changes; they attributed this phenomenon to a matching effect caused by columnar-type defects, such as the ends of twin planes (TPs) between the electrodes and/or the screw dislocations. The authors claimed that below H_{cr} the Bose glass phase characterized by such linear defects exists and above H_{cr} a glassy phase, neither vortex glass nor Bose glass, exists; furthermore the vortex glass phase was not observed. Therefore, it seems that the effect of point disorders on the Bose glass phase is still an unsolved problem.

In this paper, we study the phase transition of vortices in twinned YBCO under the condition that both correlated and point disorders are effecting the pinning. When the magnetic field is applied exactly parallel to the TPs, upon cooling the vortex liquid solidifies into a glassy phase via a second-order phase transition. We find two anomalies: the obtained liquid-to-glass transition line shows a kink near 80 kOe, and the critical exponent is nonuniversal. By assuming a matching-like effect for planar defects and considering the magnitude of the critical exponent, we discuss the glassy phase in twinned YBCO in relatively wide magnetic fields.

II. EXPERIMENT

YBCO single crystals were grown by a self-flux method using Y_2O_3 crucibles.[11] We prepared slightly overdoped samples[12] by adjusting the oxygen content, annealing the as-grown crystals at 450 °C for 168 h in 1 atm flowing oxygen gas. Twin planes having a mosaic pattern and lying in the [110] and $[1\bar{1}0]$ directions in the *ab* plane are contained in the crystals, as confirmed by a polarized light microscope and transmission electron microscopy (TEM). The superconducting transition temperature T_c, defined at the zero resistivity, is 91.8 K. The sample dimensions are $1.2\times0.9\times0.08$ mm^3. In-plane resistivity is measured by a standard direct current four-probe method under a current density $J = 1-3$ A/cm^2, giving linear resistivity, as functions of the temperature, magnetic field H up to 170 kOe, and the angle θ between the magnetic field and the TPs. The experimental setup is schematically illustrated in the left inset of Fig. 1.

III. RESULTS

The configuration of $H \parallel$ TPs ($\theta = 0°$) is obtained by measuring the angular dependence of the resistivity $\rho(\theta)$, as

 ©2004 The American Physical Society

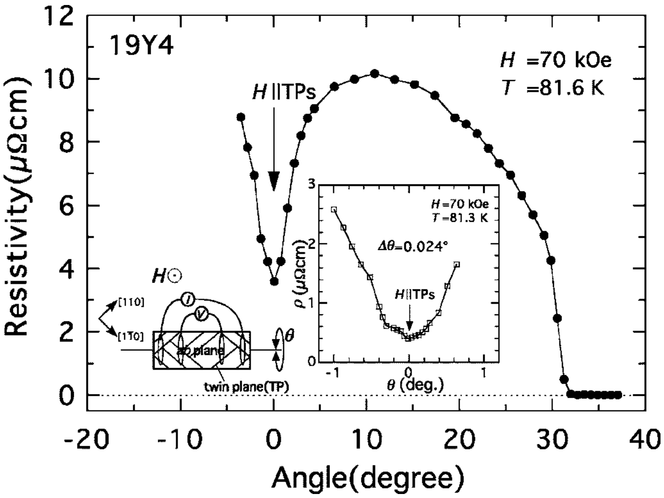

FIG. 1. Angular dependence of resistivity $\rho(\theta)$ in a relatively wide angle region at $H=70$ kOe and $T=81.6$ K. Middle inset: The $\rho(\theta)$ curve for $T=81.3$ K in the vicinity of $\theta=0°$ with a very small angle step. Left inset: The schematic experimental configuration. The directions of twin planes in the *ab* plane are [110] and $[1\bar{1}0]$.

shown in Fig. 1. The tip of the dip structure corresponds to the $H\parallel$TPs,[13] which is determined with an angle interval of 0.024°. Figure 2 demonstrates the temperature dependence of the resistivity $\rho(T,\theta)$ for both $\theta=0°$ and 5.9° in various magnetic fields up to 170 kOe. For $\theta=0°$, upon cooling the resistivity abruptly drops at a certain temperature, T_{TP}, defined as the temperature at which $\rho(T,5.9°)$ starts to deviate from $\rho(T,0°)$. This is the result of the vortex pinning caused by the TPs.[13] To more strictly determine the T_{TP}, the $\rho(\theta)$ curves should be taken as a function of the temperature, as reported by Figueras *et al.*;[14] however, we believe that our definition is also reasonable to estimate T_{TP} roughly. Below the T_{TP} the resistivity shows a gradual temperature dependence and the

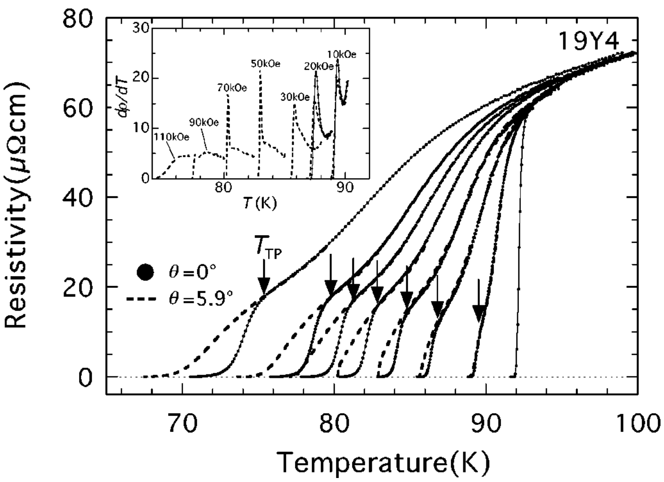

FIG. 2. Temperature dependence of resistivity for both $\theta=0°$ and 5.9° in magnetic fields of $H=170$, 110, 90, 70, 50, 30, 10, and 0 kOe (from left to right). Inset: Temperature dependence of $d\rho/dT$ for $\theta=5.9°$ at indicated fields; at $H=10$ and 20 kOe it is plotted also for $\theta=0°$.

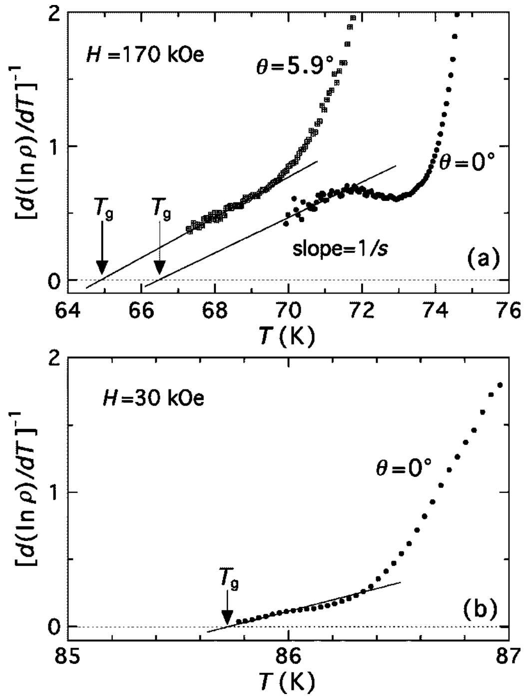

FIG. 3. $[d(\ln\rho)/dT]^{-1}$ vs T plot for magnetic fields (a) $H=170$ kOe and for both $\theta=0°$ and 5.9° and (b) $H=30$ kOe and for $\theta=0°$.

transition to zero resistivity is continuous. For $\theta=5.9°$, although this position is located in the dip structure of the $\rho(\theta)$, the dropping found for $\theta=0°$ cannot be observed in the $\rho(T)$ curves, indicating that the vortices are mostly released from the TPs. Thus, the vortex phase transition for $\theta=5.9°$ is expected to be similar to that in the untwinned sample. Actually, the $\rho(T,5.9°)$ curves show a jump in the intermediate magnetic fields and only a gradual temperature dependence with the continuous transition in high magnetic fields above 110 kOe. Here, the discontinuity in the resistivity transition is qualitatively evaluated from the $d\rho/dT$ peak. As shown in the inset of Fig. 2, the $d\rho/dT$ peak survives up to 90 kOe and fully disappears at 110 kOe, therefore we consider that the discontinuous transition occurs from 10 to 90 kOe. Although resistivity is not a thermodynamic physical property, it has been established that its jump represents a first-order vortex lattice melting transition.[2] In this study, the melting temperature T_m is defined as the temperature at $\rho=7.8\times10^{-3}$ $\mu\Omega$ cm, corresponding to a voltage criterion of 1 nV.

The continuous behavior in the $\rho(T)$ curve indicates that a second-order vortex liquid-to-glass transition occurs. Both vortex glass[3] and Bose glass[4] theories predicted the following relation, $\rho\propto(T-T_g)^s$, near a glass transition temperature T_g, where s is the critical exponent. Thus, we can find a linear relationship in the plot $[d(\ln\rho)/dT]^{-1}$ against T, where the T_g is given by the crosspoint of the extrapolation of the linear section with the T axis and s by the inverse of its slope, as demonstrated in Figs. 3(a) and 3(b). In Fig. 3(a), such plots are shown for both $\theta=0°$ and 5.9° at $H=170$ kOe. A minimum is found at around $T=73$ K for $\theta=0°$, which is attributed to the resistivity drop at T_{TP}. Al-

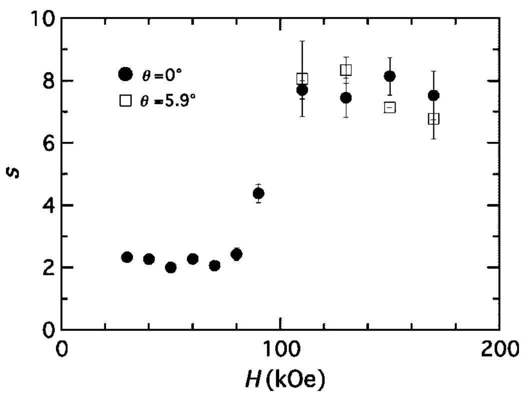

FIG. 4. Magnetic-field dependence of the critical exponent for $\theta=0°$ and 5.9°.

though the behaviors of the two plots differ above the critical region, they have a similar slope in the linear part. As shown in Fig. 3(b), the plot for $H=30$ kOe and $\theta=0°$ has a similar shape to that for $H=170$ kOe, but the slope is gentle. We note that the linear part cannot be found for both $H=10$ and 20 kOe. Figure 4 shows the magnetic-field dependence of the critical exponent $s(H)$ for both $\theta=0°$ and 5.9°. The $s(H)$ curve for $\theta=0°$ shows nonuniversal behavior and is basically separated into two distinct universal classes. On the other hand, the s values for $\theta=5.9°$ are mostly constant.

Figure 5 represents the vortex phase diagram of twinned YBCO; H_{TP}, H_g, and H_m correspond to T_{TP}, T_g, and T_m, respectively. For $\theta=0°$, the liquid-to-glass transition line $H_g(T,0°)$ shows a nonsmooth temperature dependence with a kink at about 80 kOe; the magnetic field at the kink is denoted as $H_g^{kink}(0°)$. We will present a detailed discussion of the glassy phase for $\theta=0°$ in the following sections. The first-order melting transition line $H_m(T,5.9°)$ for $\theta=5.9°$ is disrupted at 90 kOe and is well fitted by the relation predicted by Houghton *et al.*,[15] $H_m(T)=H_0(1-T/T_c)^n$, with reasonable parameters of $H_0=1102$ kOe and $n=1.34$. On the other hand, the $H_g(T,5.9°)$ line appears at 110 kOe. Both features mean that the upper critical point of the melting line $H_m^{ucp}(5.9°)$ exists between 90 and 110 kOe, which is consistent with the results for an untwinned YBCO with the same doping condition.[12] As confirmed in electron-irradiated untwinned YBCO,[16] such a change of the vortex phase transition strongly indicates that the point disorders act as an effective pinning center above 110 kOe. Furthermore, as already pointed out in Fig. 2, for $\theta=5.9°$ we can virtually neglect the pinning effect due to TPs, since they do not act as a correlated disorder. Thus, the glassy phase for $\theta=5.9°$ is vortex glass; this is also supported by the fact that the obtained s values are consistent with reported values ($s=6-8$).[5–8]

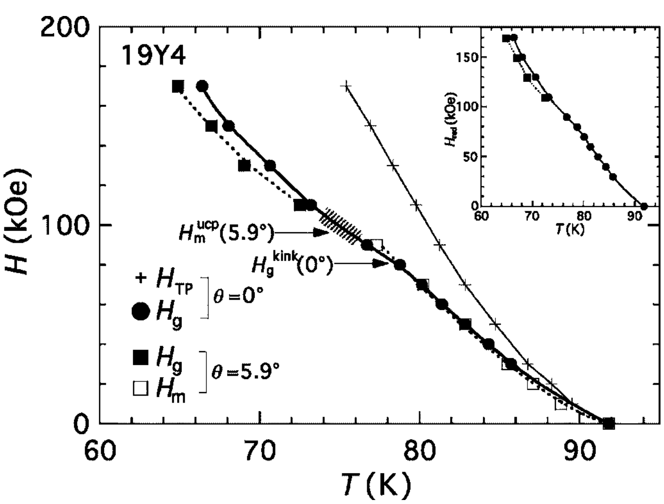

FIG. 5. Magnetic field vs temperature plane for twinned YBCO. H_{TP} is the boundary of the TPs pinning. H_g and H_m show the second-order vortex liquid-to-glass and the first-order vortex lattice melting transitions, respectively. The lines are guides for the eyes except that of $H_m(T)$, which is a fitting curve based on the vortex lattice melting theory (see text). Inset: T_g is plotted as a function of the reduced-field H_{red}; the lines are guides for the eyes.

IV. DISCUSSION

The kink structure of the $H_g(T,0°)$ line and nonuniversal critical exponent strongly indicate that the different glass phases exist above and below ~80 kOe. Since both point and correlated disorders exist in twinned YBCO, to clarify a vortex solid phase we must consider the competition among three energies, i.e., elastic E_{el}, point-disorder pinning E_p, and TPs pinning E_{TP} energies; however, it is difficult to estimate these energies quantitatively. Thus, we shall discuss a nature of the glass phase qualitatively based on the specific features of the experimental results.

First, an origin of the $H_g^{kink}(0°)$ is considered. Since the $H_g^{kink}(0°)$ is located below the $H_m^{ucp}(5.9°)$, we can discard a scenario that the point disorders cause it. A twinned YBCO has a limited number of TPs, so that there is usually a limit to the number of vortices that can be trapped on TPs. It is well known that in crystals with columnar defects the so-called matching effect occurs in the irreversibility line[17] or Bose glass phase transition line[18] at the matching field H_ϕ, where the vortex density is equivalent to the defect density. Although TP is a planar object and some vortices are naturally trapped onto a single TP, as has actually been observed by a Bitter pattern technique[19] and by scanning tunneling spectroscopy,[20] we roughly discuss such an effect. Taking a TEM image, we ascertain that the distance between neighboring TPs, d_{TP}, is 470–630 Å; this corresponds to the matching field of $H_\phi(=\phi_0/d_{TP}^2)=5.2-9.5$ kOe, which is simply estimated by assuming a square lattice state,[20] where ϕ_0 is the flux quantum. These values are considerably inconsistent with the $H_g^{kink}(0°)$ value, ~80 kOe. Here, we consider the relation between the number of vortices on a couple of TPs, n_v^{TP}, and that between them [i.e., in a twin-free (TF) region],
n_v^{TF}. When a vortex spacing a_0 is equivalent to $d_{TP}/3$, *ideally* we can obtain the condition of $n_v^{TP}=n_v^{TF}$, and it is expected that the vortex phase transition begins to be affected by the vortices in the TF region, i.e., by the pinning due to point disorders. Considering the obtained d_{TP}, this matching-like effect is expected to occur in magnetic fields of 46–86 kOe.

These values support the occurrence of the matching-like effect near 80 kOe, i.e., the $H_g^{\mathrm{kink}}(0°)$ originates from the relation between TP and vortex densities.

In magnetic fields between 30 and 80 kOe, pinning effects due to point disorders are negligible, since the first-order melting transition occurs for $\theta=5.9°$ in this regime, i.e., $E_p < E_{\mathrm{el}}$. The steep decrease in the $\rho(T,0°)$ below T_{TP} upon cooling is evidence that the TPs pinning is very effective; besides this range is below the $H_g^{\mathrm{kink}}(0°)$, i.e., $E_{\mathrm{el}} < E_{\mathrm{TP}}$. Therefore, it is valid to think that the glass phase in this region is Bose glass characterized by the planar defects. We now turn to the behavior of the critical exponent. The obtained s values are almost constant, ~ 2, and are rather smaller than those in vortex glass transition ($s=6-8$);[5–8] this universal behavior means that an identical glass phase exists in this region. It was reported that a similar magnitude of s is observed in the Bose glass phase transition on materials with correlated disorders, such as twinned YBCO ($s=2.8\pm0.2$)[9] or YBCO with columnar defects ($s=2.4\pm0.2$).[21] Furthermore, Lidmar and Wallin[22] claimed that the results of a Monte Carlo simulation also give similarly a small critical exponent, $s=2.6$. Thus, the small s values strongly support the existence of the Bose glass here. One may point out that discussing a kind of glassy phase in terms of the magnitude of s is hazardous, however, such small s values have not been observed in the vortex glass transition.[5–8] Consequently, we conclude that the glass phase in the intermediate-field regime is Bose glass. We note that Bose glass in this regime is comprised of the vortices on the TP and those not pinned by TP but strongly correlated with the vortices on TP, considering the matching-like picture.

The magnitude of the critical exponent at 90 kOe does not belong to both universal classes below 80 kOe and above 110 kOe. Since this field is located just above the $H_g(T,0°)$, it is expected that the vortices not affected by the pinning potential of TPs begin to appear at temperatures even below the Bose glass transition temperature, i.e., a vortex liquid exists between the Bose glass regions which exist on and very near the TPs. This picture resembles the following confined geometry. Marchetti and Nelson[23] theoretically studied the Bose glass transition in patterned geometry such that a region with a low irradiation dose is sandwiched between the heavier irradiated regions, where the Bose glass transition temperature in the channel region is assumed to be lower than that in both sides. Combining an inhomogeneous Bose glass scaling theory with the hydrodynamic description of vortex flow in the liquid regime, they predicted the resistivity for viscous flow of vortices in the confined channel as

$$\rho(T,L) = \rho_f(T)[1-(2\ell/L)\tanh(L/2\ell)], \qquad (1)$$

where ρ_f is the bulk resistivity without confined geometry, L is the channel width, and ℓ is the viscous length. For $\ell \ll L$, Eq. (1) represents the bulk behavior, i.e. $\rho(T,L)=\rho_f(T)\sim(T-T_g)^s$, while for $\ell \gg L$ it can be described as $\rho(T,L)\sim L^2(T-T_g)^{s+2}$. The most attractive feature is the fact that when the condition, $\ell \gg L$, is realized the critical exponent changes from s to $s+2$. This change seems to be consistent with the jump in the $s(H)$ between 80 and 90 kOe observed here, indicating that the glassy phase at $H=90$ kOe is still Bose glass. This magnetic field, 90 kOe, is located slightly above the $H_g^{\mathrm{kink}}(0°)$ and at around the $H_m^{\mathrm{ucp}}(5.9°)$, its special situation produces the vortices that are not affected directly by pinning effects due to TPs and point disorders, and enables it to cause the confined geometry. We note that this confined model cannot explain the glass transition above 110 kOe, judged from the magnitude of s; because the free vortices that are not trapped by TPs form the glassy phase regardless of the viscosity due to the Bose glass walls, as will be shown in the next section.

Considering the behaviors of the $\rho(T)$ curves shown in Fig. 2, we find that the TP pinning is effective up to 170 kOe. However, it seems too simplistic to conclude that the glass phase above 110 kOe is Bose glass, for the following reasons. First, the fact that this regime is above the $H_m^{\mathrm{ucp}}(5.9°)$ indicates that the effect of point defects on the pinning is remarkable, i.e., $E_p > E_{\mathrm{el}}$. Second, the critical exponents above 110 kOe have a universal class which is different from that below 80 kOe. Third, the position of the $H_g(T,0°)$ line above 110 kOe is lower than the extrapolation of that line below the $H_g^{\mathrm{kink}}(0°)$. These features obviously indicate that the glassy phase above 110 kOe is not conventional Bose glass. In addition, the s values of ~ 8 fit in the range of those reported in the vortex glass transition.[5–8] However, the glassy phase for $\theta=0°$ in the high-field regime also does not seem to be conventional vortex glass. The $H_g(T,0°)$ line has a higher position toward the $H_g(T,5.9°)$ line. If the glassy phases for $\theta=0°$ and 5.9° are the same, the discrepancy between the two $H_g(T)$ lines should be explained by the effective mass model.[24] This model gives the reduced field, $H_{\mathrm{red}}(\theta)=H(\cos^2\theta+\gamma^{-2}\sin^2\theta)^{1/2}$, where γ is the anisotropy parameter defined as the ratio of the coherence length in the ab-plane, ξ_{ab}, with that along the c axis, ξ_c. In the inset of Fig. 5, the glass transition temperatures for both $\theta=0°$ and 5.9° are plotted as a function of the reduced field H_{red} with $\gamma=6$.[25] One can easily find that the effective mass model cannot describe that discrepancy. The rather elevated $H_g(T,0°)$ line is thought to come from the phenomenon that the thermal fluctuation toward the vortices in TF regions is reduced by both Bose glass walls, therefore the high-field glass phase is not perfectly free from the Bose glass parts. Summarizing the discussions earlier, we can conclude that the vortices between the neighboring TPs dominate critical behavior and vortex glass transition in the TF region occurs almost independent of the viscosity due to the Bose glass walls. In other words, the vortex glass and the Bose glass coexist.

Maiorov and Osquiguil[10] found that in twinned YBCO the shape of the cusp in the angular dependence of the glass transition temperature, notated as $T_{\mathrm{BG}}(\theta)$, and a critical exponent s_x, which directly relates to the dynamic critical exponent ν,[4,22] abruptly change at a crossover field $H_{\mathrm{cr}} \sim 40$ kOe. The authors claimed that both changes originate from the matching effect of the columnar-type defects, such as the ends of the TPs between the electrodes and/or the screw dislocations. Their conclusions are as follows: below H_{cr} the solid phase is Bose glass characterized by columnar-type defects, on the other hand, above H_{cr} a glass phase having a distinct structure from Bose glass appears, although the TPs pinning is effective. Thus, the mechanism giving our $H_g^{\mathrm{kink}}(0°)$ obviously differs from that in their H_{cr}. The authors found that another field-independent critical exponent, $s_y[=\nu(z-2)]$, which is the same with s in our study, where z is the static critical exponent. Moreover, it shows a small value, ~2.5, giving a possibility that the glass phase above H_{cr} is Bose glass. However, their interpretation is basically founded upon the nonlinearity of the $T_{\mathrm{BG}}(\theta)$ curves, which we cannot discuss more deeply here; we note that our Bose glass in the intermediate-field regime consists of two kinds of vortices, pinned and unpinned by TPs, so that in terms of the theoretical definition it may also be *novel*. On the other hand, their low-field Bose glass phase, characterized by pinning due to linear defects, is not found in our crystals. A possible reason for this absence is that the TP density and/or the number of the screw dislocations in our crystals are rather smaller than those in theirs; the authors[10] also pointed out that the H_{cr} would disappear in high-quality crystals. The most remarkable difference between our and their results is the fact that the vortex glass was not found in their study. They used optimally doped samples ($T_c \sim 93$ K); we do not know why the vortex glass is absent because, in untwinned crystals with the same doping condition, the vortex glass transition occurs above ~90 kOe.[26] We speculate that their crystals have an extremely small TP spacing, i.e., dense TP density. By assuming the matching-like picture, it may be possible that the effect of point disorders on the pinning is covered, even in high magnetic fields.

Finally, we consider the vortex phase below 20 kOe. Absence of the linear part in the $[d(\ln \rho)/dT]^{-1}$ vs T plot indicates that the solid phase is not well-defined glass. As seen in the inset of Fig. 2, at 10 kOe the $d\rho/dT$ peak temperatures for both $\theta=0°$ and 5.9° are almost the same, yielding the possibility that for $\theta=0°$ there is a first-order melting transition accompanied by pinning due to the TPs, i.e., the Bragg-Bose-glass phase proposed by Giamarchi and Le Doussal[27] appears. In this field around H_ϕ, the vortices in TF regions hardly exist, and the vortices on the TPs would easily move along the boundary owing to the sparseness of them. Therefore, it is reasonable to think that the vortex structure is dominated by the vortex-vortex interaction, i.e., forming a lattice structure is allowed. Kwok *et al.*[28] also reported that the first-order melting transition accompanied by TP pinning is observed in twinned YBCO with extremely low TP density, indicating that such a first-order transition would be caused by the vortices in the TF regions. Thus, our findings are basically different from their results. The n_v^{TF} at 20 kOe is $\sim n_v^{\mathrm{TP}}/2$, thus the scenario outlined above can still be applicable, although the difference between the $d\rho/dT$ peak temperatures for both $\theta=0°$ and 5.9° is ~0.2 K. Further study is needed to clarify the vortex state in low magnetic fields.

V. CONCLUSION

We have studied the vortex phase transition in twinned YBCO in magnetic fields applied exactly parallel to the TPs. We find that the vortex phase transition is a second-order liquid-to-glass transition and its transition line $H_g(T)$ has a kink at around 80 kOe; the critical behavior is also nonuniversal above and below that magnetic field. By considering the effectiveness of point disorders on pinning and the magnitude of the critical exponent, we conclude that, except below 20 kOe, pure Bose glass appears until the magnetic field exceeds a threshold value of ~80 kOe; above 110 kOe it changes to inhomogeneous glass which is a mixture of Bose glass and vortex glass. Our findings indicate that a matching-like effect for planar defects occurs not at the so-called matching field, but at a specific magnetic field at which the number of vortices untrapped by defects exceeds that of the trapped vortices.

ACKNOWLEDGEMENTS

The authors would like to thank S. Awaji for valuable discussions. They also thank H. Komatsu and I. Kimura for supplying liquid helium, K. Higashimine for operating the TEM, and J. A. Steeh for reviewing and editing this manuscript. The experiments were in part performed at the High Field Laboratory for Superconducting Materials, Institute for Materials Research, Tohoku University. One of the authors (T.N.) was financially supported by a Grant-in-Aid for Scientific Research (Grant No. 15740208) from the Ministry of Education, Culture, Sports, Science and Technology, Japan.

[1] G. Blatter, M. V. Feigelman, V. B. Geshkenbein, A. I. Larkin, and V. M. Vinokur, Rev. Mod. Phys. **66**, 1125 (1994).
[2] For review, T. Nishizaki and N. Kobayashi, Supercond. Sci. Technol. **13**, 1 (2000), and references therein.
[3] M. P. A. Fisher, Phys. Rev. Lett. **62**, 1415 (1989); D. S. Fisher, M. P. A. Fisher, and D. A. Huse, Phys. Rev. B **43**, 130 (1991).
[4] D. R. Nelson and V. M. Vinokur, Phys. Rev. Lett. **68**, 2398 (1992); Phys. Rev. B **48**, 13060 (1993).
[5] R. H. Koch, V. Foglietti, W. J. Gallagher, G. Koren, A. Gupta, and M. P. A. Fisher, Phys. Rev. Lett. **63**, 1511 (1989).
[6] P. L. Gammel, L. F. Schneemeyer, and D. J. Bishop, Phys. Rev. Lett. **66**, 953 (1991).
[7] H. Safar, P. L. Gammel, D. A. Huse, D. J. Bishop, W. C. Lee, J. Giapintzakis, and D. M. Ginsberg, Phys. Rev. Lett. **70**, 3800 (1993).
[8] T. Naito, T. Nishizaki, and N. Kobayashi, Physica C **293**, 186 (1997); we note that in these experiments the magnetic field is *nearly* parallel to the TPs without precise angle alignment.
[9] S. A. Grigera, E. Morre, E. Osquiguil, C. Balseiro, G. Nieva, and F. de la Cruz, Phys. Rev. Lett. **81**, 2348 (1998).
[10] B. Maiorov and E. Osquiguil, Phys. Rev. B **64**, 052511 (2001).
[11] T. Naito, T. Nishizaki, Y. Watanabe, and N. Kobayashi, in *Advances in Superconductivity IX*, edited by S. Nakajima and M. Murakami (Springer-Verlag, Tokyo, 1997), p. 601.

[12] T. Nishizaki, K. Shibata, T. Naito, M. Maki, and N. Kobayashi, J. Low Temp. Phys. **117**, 1375 (1999).
[13] S. Fleshler, W.-K. Kwok, U. Welp, V. M. Vinokur, M. K. Smith, J. Downey, and G. W. Crabtree, Phys. Rev. B **47**, 14 448 (1993).
[14] J. Figueras, T. Puig, and X. Obradors, Phys. Rev. B **67**, 145503 (2003).
[15] A. Houghton, R. A. Pelcovits, and A. Subdo, Phys. Rev. B **40**, 6763 (1989).
[16] T. Nishizaki, T. Naito, S. Okayasu, A. Iwase, and N. Kobayashi, Phys. Rev. B **61**, 3649 (2000).
[17] L. Civale, A. D. Marwick, T. K. Worthington, M. A. Kirk, J. R. Thompson, L. Krusin-Elbaum, Y. Sun, J. R. Clem, and F. Holtzberg, Phys. Rev. Lett. **67**, 648 (1991).
[18] A. W. Smith, H. M. Jaeger, T. F. Rosenbaum, W. K. Kwok, and G. W. Crabtree, Phys. Rev. B **59**, R11665 (1999).
[19] G. J. Dolan, G. V. Chandrashekhar, T. R. Dinger, C. Feild, and F. Holtzberg, Phys. Rev. Lett. **62**, 827 (1989).
[20] K. Shibata, M. Maki, T. Nishizaki, and N. Kobayashi, Physica C **392–396**, 323 (2003).
[21] R. J. Olsson, W.-K. Kwok, L. M. Paulius, A. M. Petrean, D. J. Hofman, and G. W. Crabtree, Phys. Rev. B **65**, 104520 (2002).
[22] J. Lidmar and M. Wallin, Europhys. Lett. **47**, 494 (1999).
[23] M. C. Marchetti and D. R. Nelson, Phys. Rev. B **59**, 13624 (1999); Physica C **330**, 105 (2000).
[24] G. Blatter, V. B. Geshkenbein, and A. I. Larkin, Phys. Rev. Lett. **68**, 875 (1992).
[25] This value is determined by scaling the $\rho(\theta)$ curves for various magnetic fields up to 150 kOe by the redued field.
[26] K. Shibata, T. Nishizaki, T. Sasaki, and N. Kobayashi, Phys. Rev. B **66**, 214518 (2002).
[27] T. Giamarchi and P. Le Doussal, Phys. Rev. B **55**, 6577 (1997).
[28] W. K. Kwok, J. Fendrich, U. Welp, S. Fleshler, J. Downey, and G. W. Crabtree, Phys. Rev. Lett. **72**, 1088 (1994);W. K. Kwok, J. A. Fendrich, C. J. van der Beek, and G. W. Crabtree, *ibid.* **73**, 2614 (1994).