On the origin of the irreversibility line in thin $YBa_2Cu_3O_{7-\delta}$ films with and without columnar defects

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Abstract

We report on measurements of the angular dependence of the irreversibility temperature $T_{irr}(\theta)$ in $YBa_2Cu_3O_{7-\delta}$ thin films, defined by the onset of a third harmonic signal and measured by a miniature Hall probe. From the functional form of $T_{irr}(\theta)$ we conclude that the origin of the irreversibility line in unirradiated films is a dynamic crossover from an unpinned to a pinned vortex liquid. In irradiated films the irreversibility temperature is determined by the trapping angle.

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I. INTRODUCTION

The origin of the irreversibility line (IRL) in the field-temperature ($H - T$) phase diagram of high-temperature superconductors (HTS) is intriguing and still a widely discussed topic [1–9]. Experimentally, this line is defined as the border line at which the magnetic response of the sample changes from irreversible to reversible. In HTS, large fluctuations and relatively weak pinning lead to a rich $H - T$ phase diagram with a variety of dynamic and static transitions which can be responsible for the appearance of magnetic reversibility [3–5,10–12].

Thus, a thorough experimental investigation of the IRL is important for the understanding of the vortex-lattice behavior in superconductors in general and of the mechanisms responsible for the onset of irreversible magnetic response, in particular.

Several models, like thermally activated depinning [2,4,13–14], vortex-lattice melting [15–21] and a transition from vortex glass to vortex fluid [23–27], were proposed to identify the origin of the IRL in HTS. Also, attention was given to the possibility of pinning in the vortex-liquid phase [2,10,11,26] and to different dissipation mechanisms above the melting line [3,13,27–29]. Irreversibility due to geometrical [1,8] or surface barriers [22] have also been proposed, but these mechanisms are less probable in thin $YBa_2Cu_3O_{7-\delta}$ films with strong pinning. The irreversibility line may be affected by sample-dependent properties such as the nature and density of pinning centers and by intrinsic or extrinsic anisotropy. For example, in superconductors with columnar defects, the irreversibility line may either be identified with the Bose-glass transition [3,29,37], or related to the concept of a trapping angle [38].

The configuration of the columnar defects is also very important, since it affects possibility for different types of depinning mechanism. A splayed configuration, for example, inhibits creep from columnar defects [3,40]. Similarly ”crossed” defects (i.e. defects at angles $\pm \theta$) were shown to act collectively, i.e., they introduce unidirectional anisotropy such that the current density reaches its maximum for magnetic field directed in a mid angle between defects [41,12].

Experimentally, the situation is even more complex, since different techniques (magne-
tization loops, field-cool vs. zero-field-cool DC magnetization, peak in the imaginary part of the first harmonic etc.) yield different IRLs \[24,43\]. To a great extent, the reliability of the determination of IRL depends on the criterion for the onset of the irreversibility. We determine the irreversibility temperature at given DC field by the onset of third harmonic in the \textit{ac} response, which, we believe, is one of the most reliable methods for contactless determination of the IRL \[44\]. In most experiments \(T_{\text{irr}}\) is measured as a function of the external field \(H\). This information is insufficient to distinguish between different models for the origin of the irreversibility. Additional information, like the frequency dependence of the IRL \[24,45\] or its angular variation \[18,20,21,32,34,36,46\], is needed.

In this paper we report on a study of the angular dependence of the irreversibility temperature \(T_{\text{irr}}(\theta)\) in thin \(YBa_2Cu_3O_{7-\delta}\) \((YBCO)\) films before and after irradiation with Pb ions.

**II. EXPERIMENTAL**

The 1500 Å \(YBCO\) films were ”sandwiched” between \(SrTiO_3\) layers \[17\]. First, a 500 Å layer of \(SrTiO_3\) was deposited on a \(MgO\) substrate. Then, the \(YBCO\) film was laser ablated on top of the \(SrTiO_3\) and finally, the \(YBCO\) was covered by a protective 300 Å layer of \(SrTiO_3\). All three samples have the same lateral dimensions of 100 × 500 \(\mu m^2\). One film, denoted as \(REF\), was used as a reference sample. The other two, \(UIR\) and \(CIR\), were irradiated at GANIL with \(2 \times 10^{11} \text{ ions/cm}^2\) \(5.8 \text{ GeV}\) Pb ions along the \(c-\text{axis}\) and along \(\theta = \pm 45^\circ\), respectively. \((UIR \text{ and } CIR \text{ stand for } ”\text{uniform irradiation}” \text{ and } ”\text{crossed irradiation}”\), respectively). The superconducting transition temperatures, measured by a \textit{Quantum Design SQUID} susceptometer and defined as the onset of the Meissner expulsion in a DC field of 5 \(G\), are \(T_c \approx 89 \text{ K}\) for the samples \(REF\) and \(UIR\) and 88 \(K\) for \(CIR\).

For the \textit{ac} measurements reported below we used a miniature \(80 \times 80 \mu m^2\) InSb Hall-probe, which was positioned in the center of the sample. The 1 \(G\) \textit{ac} magnetic field, always parallel to the \(c-\text{axis}\), was induced by a small coil surrounding the sample. An external \textit{dc}
magnetic field, up to $H_a = 1.5\ Tesla$, could be applied at any direction $\theta$ with respect to the $c-$ axis. In our experiments DC magnetic field was always turned on at a fixed angle at $T > T_c$ and then the $ac$ response was recorded during sample cooling. The irreversibility temperature, $T_{irr}(\theta)$ is defined as the onset of the third harmonic signal in the $ac$ response measured by the Hall probe [44]. This procedure was repeated for various $dc$ fields and at various angles $\theta$ of the field with respect to the $c-$axis.

III. RESULTS

Fig. 1 presents measurements of $V_3$, the third harmonic in the $ac$ response, versus temperature $T$, during field-cooling at 1 Tesla for the sample REF at various angles between 0 and 90°. Apparently, as the angle $\theta$ increases the whole $V_3$ curve shifts to higher temperatures and becomes narrower. The onset of irreversibility, $T_{irr}(\theta)$, is defined by the criterion $V_3^{onset} = 0.05$ in the units of Fig. 1.

Fig. 2 exhibits typical $T_{irr}(\theta)$ data for the unirradiated sample REF, measured at two values of the external field: 0.5 Tesla and 1 Tesla. Both curves exhibit a shallow minimum around $\theta = 0$ and they reach their maximum values for $H$ along the $ab$ plane, at angles $\theta = \pm 90^\circ$. We also measured the frequency dependence of $T_{irr}$ for the same values of magnetic field. As shown in Fig. 3, the slope $\partial T_{irr}/\partial \ln (f)$ is larger for larger field.

The sample irradiated along the $c-$axis exhibits additional feature - a peak around $\theta = 0$. This is clearly shown in Fig. 4 where we compare $T_{irr}(\theta)$ at $H = 1\ Tesla$ for the samples REF and UIR. As discussed below, this peak is a signature of the unidirectional magnetic anisotropy induced by the columnar defects. Intuitively, one would therefore expect two peaks, along $\theta = \pm 45^\circ$, for the third sample, CIR, crossed-irradiated at $\theta = \pm 45^\circ$. Instead, we find one strong peak around $\theta = 0$, similar to that found in BSCCO crystals [42]. This is demonstrated in Fig. 5 where we compare $T_{irr}(\theta)$ at $H = 0.5\ Tesla$ for this sample (CIR) and for the unirradiated sample (REF). We argue below that the peak around $\theta = 0^\circ$ is a result of a collective action of the crossed columnar defects, and its origin is the same as
that for unidirectional enhancement of critical current density observed in BSCCO crystals [11,12].

IV. ANALYSIS

The "true" irreversibility temperature $T_0$ is defined as a temperature above which the pinning vanishes. Such disappearance of pinning can be of static (true phase transition), as well as of dynamic origin (gradual freezing, pinning in liquid). In practice, one determines the irreversibility temperature $T_{irr}(\Delta)$ as the temperature above which the critical current density is less than some threshold value $\Delta$. Therefore, by definition, $T_0 = \lim_{\Delta \to 0} (T_{irr}(\Delta))$. The apparent current depends on temperature $T$, magnetic field $B$ and the frequency $f$ of the exciting field which defines a characteristic time-scale $1/f$ for the experiment. By solving the equation $j(T,B,f) = \Delta$ with respect to $T$ one finds the experimental irreversibility temperature $T_{irr}$ for constant $B$ and $f$. In the following we argue that in our experiments the measured $T_{irr}$ is a good approximation of $T_0$. In order to estimate $T_{irr}$ we employ a general form for the apparent current density in the vicinity of the irreversibility line (IRL) [4–6,48]:

$$j(T,B,f) \propto j_c(0) \left( \frac{1-T/T_0}{B/B_0} \right)^\beta \left( \frac{f}{f_0} \right)^\gamma$$

where the parameters $B_0$ and $f_0$ are temperature independent (Eq. 1 is thus valid only in a narrow temperature interval near the IRL and for fields larger than $H_{c1}$). From Eq. 1 we get:

$$T_{irr} = T_0(B) \left( 1 - \frac{\Delta}{j_c(0)} \left( \frac{B}{B_0} \right)^\beta \left( \frac{f_0}{f} \right)^\gamma \right)^\frac{1}{\alpha}$$

(2)

Inserting reasonable numerical estimates: $j_c(0) \simeq 10^7$ A/cm$^2$, $\Delta \simeq 100$ A/cm$^2$ for our experimental resolution, $B_0 \simeq 10^3$ G, $B \simeq 10^4$ G, $\beta \simeq 1$ [45], $\gamma \simeq 1$ [6], $f \simeq 10^2$ Hz, and $f_0 \simeq 10^7$ Hz [6], we get from Eq. 2: $T_{irr} = T_0(B) \left( 1 - 0.005^{1/\alpha} \right)$. Thus, with 0.5% accuracy we may say that $T_{irr}$, the measured onset of the third harmonic component in the
ac response, marks some "true" irreversibility crossover line \( T_0 (B) \). The nature of this line \( T_0 (B) \) is our main interest, since, as discussed in the introduction, it is directly related to the pinning properties of vortex lattice in type-II superconductors at high temperatures.

A. Unirradiated YBCO film

We turn now to consider the effect of the intrinsic anisotropy on \( T_{irr} (\theta) \). Following the anisotropic scaling approach proposed by Blatter et al. \cite{49,50}, we replace \( T \) by \( \varepsilon T \) and \( B \) by \( B_{eff} = \varepsilon \theta B \), where \( \varepsilon \theta = \sqrt{\cos^2 (\theta) + \varepsilon^2 \sin^2 (\theta)} \) and \( \varepsilon \approx 1/7 \) is the anisotropy parameter for YBCO. It should be emphasized that we can use this scheme only in the case of intrinsic anisotropy \( \varepsilon = \sqrt{m_{ab}/m_c} \), where \( m_c \) and \( m_{ab} \) denote the effective masses of the electron along the \( c \)-axis and in the \( ab \)-plane, respectively. In the case of some extrinsic magnetic anisotropy, (columnar defects or twin planes), the critical current depends on the angle not only via the effective magnetic field \( B_{eff} \), but also because of this extrinsic anisotropy.

As we have already indicated in the Introduction, there are several possible origins for a crossover from irreversible to reversible magnetic behavior in unirradiated samples. We exclude the vortex-glass to vortex fluid transition as a possible origin for the IRL, because this transition was shown to occur at temperatures lower than the onset of dissipation \cite{24,26,51}. The thermal depinning temperature increases with increase of field \( T_{dp} \propto \sqrt{B} \) and, therefore is excluded as well. Vortex-lattice melting transition is believed to be responsible for the appearance of reversibility \cite{13,14,18,19}. The explicit angular dependence of \( T_m \) was derived by Blatter et al. \cite{13,14,18,19} using their scaling approach:

\[
T_m (\theta) \approx 2\sqrt{\pi \varepsilon \varepsilon_0 c_L^2} (\Phi_0/B\varepsilon_0)^{1/2} \approx \frac{c_L^2 T_c}{\sqrt{\beta_m G\dot{\gamma}}} \left( 1 - \frac{T_m}{T_c} \right) \left( \frac{H_{c2} (0) \xi^3 (0)}{\varepsilon \dot{\gamma} B} \right)^{1/2},
\]

where \( \Phi_0 \) is the flux quantum, \( \xi \) is the coherence length, \( \beta_m \approx 5.6 \) is a numerical factor, estimated in \cite{13}, \( c_L \approx 0.1 \) is the Lindemann number, \( G\dot{\gamma} = (T_c/\varepsilon H_{c2} (0) \xi^3 (0))^2 /2 \) is the Ginzburg number, and \( H_{c2} (0) \) is the linear extrapolation of the upper critical field from \( T_c \) to zero. Solving Eq. 3 with respect to \( T_m \) we get:
\[
T_m (\theta) \simeq \frac{T_c}{1 + \left( \frac{\beta_m G_i}{c_L^4 H_{c2}(0)} \right) \left( \varepsilon_0 B \right)^{\frac{1}{2}}} \equiv \frac{T_c}{1 + C \sqrt{\varepsilon_0 B}}.
\]

Eq. 4 predicts that the melting temperature decreases as \( B_{\text{eff}} \) increases. This is due to the fact that the inter-vortex distance \( a_0^2 \propto 1/B_{\text{eff}} \) decreases faster than the characteristic amplitude of fluctuations \( \langle u^2(B_{\text{eff}}, T_m) \rangle_{\text{th}} \propto 1/\sqrt{B_{\text{eff}}} \). Therefore, the condition for the vortex-lattice melting \( \langle u^2(B_{\text{eff}}, T_m) \rangle_{\text{th}} \simeq c_{L}^2 a_0^2 \) implies larger melting temperatures for smaller effective fields, i.e., for larger angles. In agreement with this prediction, the experimental data of Fig. 2 show that \( T_{\text{irr}} \) increases with the angle, i.e., decreases with \( B_{\text{eff}} \). The solid lines in Fig. 2 are fits to equation 4. From this fit we get \( C \simeq 0.0005 \). However, a reasonable estimate of \( C \simeq \sqrt{\beta_m G_i/c_L^4 H_{c2}(0)} \) yields \( C \simeq 0.01 \), where we take \( H_{c2}(0) = 5 \cdot 10^6 \, \text{G}, \, c_L = 0.1, \, G_i = 0.01, \) and \( \beta_m = 5.6 \) [5]. Also, Yeh et al. showed that the onset of irreversibility occurs above the melting temperature (Ref. [29], Fig. 4). In addition, the important effect of the frequency (see Fig. 3) is not included in Eq. 4.

We discuss now another possibility for the onset of the irreversibility, namely, pinning in the vortex liquid (for a discussion see Ch. VI in Blatter et al. [5] and references therein). Any fluctuation in the vortex structure in the liquid state has to be averaged over the characteristic time scale for pinning \( t_{\text{pin}} \). In the absence of viscosity the only fluctuations in the liquid state are thermal fluctuations, which have a characteristic time \( t_{\text{th}} \ll t_{\text{pin}} \). (As shown in [3] \( t_{\text{pin}}/t_{\text{th}} \propto j_0/j_c \), where \( j_0 \) is the depairing current). Thus, such a liquid is always unpinned. The situation is different for a liquid with finite viscosity. In this case there exists another type of excitations in the vortex structure, i.e. plastic deformations with a characteristic time scale \( t_{\text{pl}} \). The energy barrier, corresponding to plastic deformation is shown to be [3][12]:

\[
U_{\text{pl}} \simeq \gamma \varepsilon_0 a_0 \simeq \gamma \left( \frac{H_{c2}}{4G_i} \right)^{\frac{1}{2}} (T_c - T) B^{-1/2}.
\]

where \( \gamma \) is a coefficient of the order of unity. The corresponding characteristic time scale is:

\[
t_{\text{pl}} \sim t_{\text{th}} \exp \left( U_{\text{pl}}/T \right)
\]
Thus, depending on the viscosity, $t_{pl}$ can be smaller or larger than $t_{pin}$. In the latter case, after averaging over a time $t_{pin}$, the vortex structure remains distorted and such a liquid shows irreversible magnetic behavior. Thus, on the time scale of $t_{pin}$ the distorted vortex structure is pinned. The crossover between pinned and unpinned liquid occurs at temperature $T_k$ where the characteristic relaxation time for pinning $t_{pin}(T)$ becomes comparable to that for plastic motion $t_{pl}(T)$. Thus, using Eqs. 3 and 5 we obtain:

$$T_k = \frac{T_c}{1 + \frac{1}{\gamma} \left( \frac{4Gi}{H_{c2}(0)} \right)^{\frac{1}{2}} \ln \left( t_{pin}/t_{th} \right) \sqrt{B}}$$

(7)

Finally, using the anisotropic scaling [49] we may rewrite Eq. 4 for $f_{pin} < f < f_{th}$ as:

$$T_{irr}(\theta) = T_k(\theta) = \frac{T_c}{1 + \frac{1}{\gamma} \left( \frac{4Gi}{H_{c2}(0)} \right)^{\frac{1}{2}} \ln \left( f_{th}/f \right) \sqrt{B} \sqrt{\varepsilon_\theta}} \equiv \frac{T_c}{1 + A \sqrt{\varepsilon_\theta B}}$$

(8)

with $f_{th} \equiv 1/t_{th}$ and $f_{pin} \equiv 1/t_{pin}$. Note the apparent similarity with the expression for the melting temperature, Eq. 4. The numerical estimate for $A = \frac{1}{\gamma} \left( \frac{4Gi}{H_{c2}(0)} \right)^{\frac{1}{2}} \ln \left( f_{th}/f \right)$ gives: $A \approx 10^{-4} \ln \left( f_{th}/f \right) / \gamma$. This is in agreement with the value found from the fit (solid line in Fig. 2) for $H_{c2}(0) = 5 \cdot 10^6 G$, $Gi = 0.01$, $f_{th} \sim 10^{10} Hz$, and $\gamma \approx 4$.

To further confirm that in our $YBCO$ films the most probable physical mechanism for the onset of irreversibility is a dynamic crossover from unpinned to pinned vortex liquid we discuss now the frequency dependence of $T_{irr}$. Equation 8 has a clear prediction for the frequency dependence of $T_{irr}$. To see it directly we may simplify it by using the experimentally determined smallness of value of the fit parameter $A \approx 0.0005$, which allows to expand Eq. 8 (for not too large fields) as

$$T_k \approx T_c \left( 1 - \frac{1}{\gamma} \left( \frac{4Gi}{H_{c2}(0)} \right)^{\frac{1}{2}} \ln \left( f_{th}/f \right) \sqrt{\varepsilon_\theta B} \right)$$

(9)

which results in a linear dependence of $T_{irr}$ upon $\ln(f)$ and a slope $S \equiv \partial T_{irr}/\partial \ln(f) \approx \frac{T_c}{\gamma} \left( \frac{4Gi}{H_{c2}(0)} \right)^{\frac{1}{2}} \sqrt{\varepsilon_\theta B} = T_c A \sqrt{\varepsilon_\theta B} \ln \left( f/f_{th} \right)$. Note that the slope is proportional to $\sqrt{B}$. This is indeed confirmed by the experimental data, as is demonstrated by the solid lines in Fig. 2. From this fit we get $S/\sqrt{B} = 0.004$ and we can independently verify the parameter $A$
appeared in Eq. 8 $A = S/ \left( T c \sqrt{\varepsilon \theta} B \ln (f/f_{th}) \right) = 0.0008$, which is in an agreement with the value obtained above.

We note that the approximated expression for the frequency dependence of $T_{irr}$, Eq. 8, is valid in the whole experimentally accessible range of magnetic field since Eq. 8 predicts a maximum in the slope $S$ at $B_{max} = \left( A \sqrt{\varepsilon \theta} \right)^{-2} \approx 400$ Tesla for the experimental parameters. This value is, of course, beyond the experimental limits and, probably, even exceeds $H_{c2}$.

Another support for the onset of the irreversibility in a vortex liquid is the $ac$ field amplitude dependence of the IRL. In both, thermal-activated (TAFF) and pure flux-flow (FF) regimes the I-V curves are linear and the onset of the third harmonic is due to a change in the slope (from $\rho_{FF}$ to $\rho_{TAFF}$). In this case we expect the amplitude dependence for this onset. Contrary, at the melting transition the onset of irreversibility is sharp and is not expected to depend upon the amplitude of the $ac$ field. In our experiments we find a pronounced amplitude dependence of the IRL, thus confirming the above scenario.

### B. Irradiated YBCO films

For the irradiated films the situation is quite different. The models for $T_{irr}(\theta)$ in unirradiated films cannot explain the experimental features exhibited in Figs. 4 and 5, in particular the increase in $T_{irr}$ in the vicinity of $\theta = 0$. Such discrepancy can only be due to the angular anisotropy introduced by columnar defects, i.e., the angle-dependent pinning strength. It was shown, both theoretically [5,30] and experimentally [37], that for magnetic field oriented along the defects the irreversibility line is shifted upward with respect to the unirradiated system. Thus, our results in Fig. 3 suggest that the measured $T_{irr}(\theta)$ is a superposition of the angular variation of $T_{irr}$ in unirradiated film (denoted in this section as $T_{irr^{REF}}$) and the anisotropic enhancement of the pinning strength due to irradiation.

We can estimate the latter contribution by employing the concept of a "trapping angle" $\theta_t$, the angle between the external field and the defects at which vortices start to be partially trapped by columnar tracks. (For a schematic description, see Fig. 43 in Blatter et al. [4]).
As we show in Appendix A,

\[ \tan(\theta_t) \approx \sqrt{2\varepsilon_r/\varepsilon_l} \]  

(10)

where \( \varepsilon_r(T) \) is the trapping potential of a columnar defect and \( \varepsilon_l \) is the line tension. In the experiment we cool down at a fixed \( \theta \), and the onset of irreversibility must occur when \( \theta = \theta_t(T) \), provided that the temperature is still larger than \( T_{irr}^{REF}(\theta) \). Otherwise, the onset occurs at \( T_{irr}^{REF} \). This defines the condition for the irreversibility temperature \( T_{irr} \) for angles \( \theta \leq \theta_c \equiv \theta_t(T_{irr}^{REF}(\theta_t)) \) ≃ 50° in our case:

\[ \tan(\theta) = \tan(\theta_t(T_{irr})) \approx \sqrt{2\varepsilon_r/\varepsilon_l} \]  

(11)

At high-temperatures \( \varepsilon_r(T) \propto \exp\left(-T/\tilde{T}_{dp}\right) \), where \( \tilde{T}_{dp} \) is the depinning energy [5]. Thus, we can write for \( T_{irr} \):

\[ T_{irr}(\theta) = \begin{cases} T_{irr}^{REF}(\theta) - D \ln(C|\tan(\theta)|) & \theta \leq \theta_c \\ T_{irr}^{REF}(\theta) & \theta > \theta_c \end{cases} \]  

(12)

where \( D \) and \( C \) are constants. This expression is in an agreement with our results shown in Fig. 4 (solid line). We note, however, some discrepancy in the vicinity of \( \theta = 0 \), where we find quite weak dependence of \( T_{irr} \) on angle. We explain this deviation by considering the influence of relaxation, which, in the case of parallel defects depends on angle. The relaxation rate is maximal, when vortices are aligned along the defects and retains its normal ”background” value for perpendicular direction [52]. Vortex, captured by a defect, can nucleate a double kink which slides out resulting in a displacement of a vortex on a neighboring column. In our irradiated samples the defect lattice is very dense (the matching field \( B_\phi = 4 \) Tesla, i.e., distance between columns \( d \approx 220 \) Å) and such double-kink nucleation an easy process. Thus, the irreversibility temperature should be shifted down around \( \theta = 0 \) as compared to the ”ideal”, non-relaxed value, Eq. [12]. This explains the reduction in \( T_{irr} \) in Fig. 4.

We may now conclude that in irradiated films, for angles less than the critical angle \( \theta_c \), the irreversibility line is determined by the trapping angle \( \theta_t \). The Bose-glass transition can
probably only be found for small angles within the lock-in angle $\theta_L \leq 10^\circ$. This conclusion is also indirectly confirmed in [53].

As was pointed out in the introduction, crossed defects should hinder the relaxation due to forced entanglement of vortices. Thus, the irreversibility temperature is expected to be closer to that predicted by Eq. 12. Fig. 3 shows a good agreement of the experimental data with Eq. 12 (solid line). To explain why defects crossed at large angle act collectively and force unidirectional magnetic anisotropy, we follow here the approach outlined in [41], and extend that description to account for arbitrary orientation of the external field with respect to the crossed columnar defects and to the $c-$axis. In Ref. [41] the authors consider possible motion of vortices in a "forest" of crossed defects for field oriented along the $c-$axis. In our case of a dense lattice we may exclude from consideration free kink sliding and consider only depinning from the intersections. We sketch in Fig. 6 the two limiting situations: (a) the external field is parallel to one subsystem of the columnar defects ($\theta = 45^\circ$) and (b) the external field is oriented along the $c-$axis, between crossed columns ($\theta = 0$). In case (a), Fig. 6a, vortices can depin just by nucleation the single kinks which are sliding from intersection to intersection, or, by nucleation of super-kinks resulting in a kind of motion, similar to a variable-range hopping. This type of thermally assisted vortex depinning does not cost any additional energy on vortex bending. Another situation arises for field along the $c-$axis, Fig. 6b. Now vortices can depin only via nucleation of multiple half-loops, which characteristic size depends upon current density. This results in additional barrier for vortex depinning, which even diverges at zero current [4]. As a result, the relaxation rate is anisotropic, i.e, it is suppressed when the external field is oriented along the mid direction between the two subsystems of the crossed columnar defects. This is just opposite to a situation in uniformly irradiated samples.
V. CONCLUSIONS

We presented angle-resolved measurements of the irreversibility temperature in unirradiated $YBa_2Cu_3O_{7-\delta}$ film and in two films with columnar defects, induced by $6 \, GeV$ Pb-ions irradiation, either parallel to the $c$–axis or 'crossed' in $\theta = \pm 45^\circ$. We find that in the unirradiated film the transition from irreversible to reversible state occurs above the melting line and marks the crossover from pinned to unpinned vortex liquid. In irradiated films, within the critical angle $\theta_c \simeq 50^\circ$, the irreversibility line is determined by the temperature dependent trapping angle. For larger angles $T_{irr}$ is determined by the intrinsic anisotropy via the effective field. The formulae for $T_{irr}(\theta)$ for both unirradiated and irradiated films are given. We also discuss the possible influence of anisotropic enhancement in relaxation rate which leads to a smearing of the expected cusp at $\theta = 0$ in the $T_{irr}(\theta)$ curve in the uniformly irradiated film. Finally, we demonstrate the collective action of crossed columnar defects, which can lead to suppression of relaxation and enhancement of pinning strength along the mid direction.

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APPENDIX A

We describe here the derivation of our Eq. (10), which differs slightly from the analogous Eq. (9.173) of Blatter et al. [5]. We derive it using exactly the same approach (and notions) as in [5], but, in view of the experimental situation, avoid assumption of small angles, which allows Blatter et al. to approximate \( \tan (\theta) \approx \sin (\theta) \approx \theta \). In order to estimate the trapping angle one has to optimize the energy change due to the vortex trapping by columnar defects. This energy is written as [5]:

\[
\varepsilon (r, \theta) = \varepsilon_l \left[ r + \left( d^2 + \left( \frac{d}{\tan (\theta)} - r \right)^2 \right) - \frac{d}{\sin (\theta)} \right] - r \varepsilon_r \tag{1A}
\]

where \( r (\theta) \) is the length of the vortex segment, trapped by a defect; \( d \) is the distance between the columns; \( \varepsilon_l \) is the line tension; and \( \varepsilon_r \) is the trapping potential of the defects. The variation of Eq. (1A) with respect to \( r \) at fixed an angle \( \theta \) defines the angular dependence of \( r (\theta) \). The trapping angle \( \theta_t \) can be found by solving the equation \( r (\theta_t) = 0 \). This results in

\[
\tan (\theta_t) = \frac{\sqrt{\varepsilon_r (2\varepsilon_l - \varepsilon_r)}}{\varepsilon_l - \varepsilon_r} \tag{2A}
\]

which, at sufficiently small \( \varepsilon_r \), can be approximated as

\[
\tan (\theta_t) = \sqrt{\frac{2\varepsilon_r}{\varepsilon_l}} + O \left( \frac{\varepsilon^3}{2} \right) \tag{3A}
\]

Apparently, at very small angles we recover the original result of Ref. [5]. In the paper, for the sake of simplicity, we use expression Eq. (3A) instead of the full expression Eq. (2A). However, as noted above we cannot limit ourselves to small angles and, generally speaking, the trapping angle may be quite large (\( \theta_t \approx 40^\circ \) in our case). The error due to use of Eq. (1A) can be estimated as follows: at \( \theta \approx 40^\circ \) Eq. (2A) gives \( \varepsilon_r / \varepsilon_l \approx 0.24 \), whereas Eq. (3A) gives \( \varepsilon_r / \varepsilon_l \approx 0.35 \), which is suitable for our implication of Eq. (3A) since we consider exponential decrease of \( \varepsilon_r \). Also, as shown in [5] in a system with anisotropy \( \varepsilon \), the trapping angle is enlarged by a factor of \( 1/\varepsilon \).
FIGURES

FIG. 1. The third harmonic signal $V_3$ versus temperature during field-cooling at 1 Tesla for sample $REF$ at $\theta = 0$, 10º, 30º, 40º, 60º, 80º, 90º.

FIG. 2. The irreversibility temperature in the unirradiated sample $REF$ at two values of the external field: $H = 0.5$ and 1 Tesla. The solid lines are fits to Eq. 8.

FIG. 3. The frequency dependence of $T_{irr}$ in the unirradiated sample $REF$ at two values of the external field: $H = 0.5$ and 1 Tesla. The solid lines are fits to Eq. 3.

FIG. 4. The irreversibility temperature for two samples: $REF$ (unirradiated - open circles) and $UIR$ (irradiated along the $c-$ axis sample, - filled circles) at $H = 1$ Tesla. Solid lines are fits to Eq. 8 and Eq. 12, respectively.

FIG. 5. The irreversibility temperature for two samples: $REF$ (unirradiated - open circles) and $CIR$ (irradiated along $\theta = \pm 45^o$ - filled circles) at $H = 0.5$ Tesla. Solid lines are fits to Eq. 8 and Eq. 12, respectively.

FIG. 6. Schematic description of a possible depinning modes of a vortex line in the case of crossed columnar defects; (a) magnetic field is directed along $\theta = 45^o$; (b) magnetic field is along $\theta = 0$. 
Prozorov et al.  Fig. 1
Prozorov et al.  Fig. 2

![Graph showing $T_{irr}$ (K) vs. $\theta$ for 0.5 Tesla and 1 Tesla magnetic fields. The graph plots $T_{irr}$ against $\theta$ with data points and curves for 0.5 Tesla and 1 Tesla fields.](image-url)
Prozorov et al. Fig. 3

0.5 Tesla

1 Tesla
Prozorov et al. Fig. 4
0.5 Tesla

\[ \frac{T_{\text{irr}}}{T_c} \]

\( \theta \)

Prozorov et al. Fig. 5
Prozorov et al.