Direct determination of the collective pinning radius in the cuprate superconductor $Bi_2Sr_2CaCu_2O_{8+\delta}$

M. I. Dolz, A. B. Kolton, and H. Pastoriza
Centro Atómico Bariloche, Comisión Nacional de Energía Atómica, R8402AGP San Carlos de Bariloche, Argentina
(Received 21 August 2009; revised manuscript received 4 November 2009; published 3 March 2010)

We study the onset of the irreversible magnetic behavior of vortex matter in micron-sized $Bi_2Sr_2CaCu_2O_{8+\delta}$ single crystals by using silicon micro-oscillators. We find an irreversibility line appearing well below the thermodynamic Bragg-glass melting line at a magnetic field which increases both with increasing the sample radius and with decreasing the temperature, paradoxically implying the existence of a reversible vortex solid. We show that at this irreversibility line, the sample radius can be identified with the crossover length between the Larkin and the random-manifold regimes of the vortex-lattice transverse roughness. Our method thus allows to determine, as a function of temperature and applied magnetic field, the minimum size of a vortex system that can be collectively pinned, or the so-called three-dimensional weak collective pinning Larkin radius, in a *direct way*.

DOI: 10.1103/PhysRevB.81.092502 PACS number(s): 74.25.Uv, 74.25.Dw, 74.72.−h, 85.85.+j

The frustrating competition between elasticity and disorder is at the root of the universal glassy behavior displayed by many physical systems. These include periodic systems such as charge-density waves[1,2] and Wigner crystals[3] or interfaces, such as magnetic[4–6] and ferroelectric[7] domain walls, liquid menisci,[8] and fractures.[9] Among these systems, superconducting vortex lattices provide an exceptional test ground for the theory of elastic objects embedded in a random quenched environment since their density and interactions can be experimentally tuned.

Larkin and Ovchinikov[10] demonstrated the unavoidable impact of arbitrarily weak disorder on an otherwise perfect elastic lattice. To produce a measurable pinning however, a weak disorder has to act in a minimum region of space in order to compete with the elasticity which dominates the physics at smaller scales. It is the finiteness of the so-called Larkin lengths which fundamentally explains the mere existence of pinning and measures its effective strength on the extended system.[11] In the modern elastic theory, designed to correctly describe the large-scale static and dynamical universal behavior of elastic manifolds, the Larkin lengths are the fundamental input for making quantitative predictions for a given experimental system. Determining the Larkin lengths, in general, and for a vortex lattice in a high-T_c superconductor, in particular, remains a difficult challenge however. Indirect empirical estimates based on transport properties such as the critical current or the creep barriers are an alternative but they spoil precise comparisons between experiments and theory. Recently, an experimental finite-size analysis was applied to determine the characteristic dynamical length, predicted by the (bulk) elastic theory, controlling the domain-wall creep motion in ferromagnetic nanowires.[6] Here we report a finite-size study in micron-sized superconductors showing how collective pinning arises from weak disorder at the so-called vortex-matter Larkin radius that we determine in a *direct* way.

The elastic theory characterizes the translational order by the roughness function $W(\mathbf{r})=\langle[u(\mathbf{r})-u(\mathbf{0})]^2\rangle$, with $u(\mathbf{r})$ the vortex displacement field with respect to the perfect lattice. For the Bragg-glass (BG) phase,[12] expected for weak pinning at low enough temperatures, a logarithmic growth of W is predicted at large distances. At short distances however, we have a *Larkin* regime where displacements grow as $W\sim r^{4-d}$, with $d=3$ the internal dimension of the elastic manifold. A crossover to the random-manifold (RM) regime at a distance $\mathbf{r}_c$ occurs when displacements are comparable to the pinning force range r_p, such that $W(\mathbf{r}_c)\sim r_p^2$. For a superconductor with vortices directed along the direction $\hat{z}$ of an external magnetic field, this crossover defines the longitudinal $(L_c\equiv|\mathbf{r}_c\cdot\hat{z}|)$ and transverse $(R_c\equiv|\mathbf{r}_c-L_c\hat{z}|)$ Larkin lengths. Besides describing a geometrical crossover, the Larkin lengths also determine the size of the minimum bundle of vortices that can be individually pinned by the quenched disorder.[10]

We fabricated $Bi_2Sr_2CaCu_2O_{8+\delta}$ (BSCCO) samples with a procedure similar to that of Wang *et al.*[13] Disks of radii $R_s=6.75$, 13.5, and 25 μm, and $d=1$ μm of thickness and critical temperatures $T_c=84.2$, 85.3, and 89.6 K, respectively, were made up by lithography and ion etching and then glued to high-Q silicon torsional micro-oscillators.[14,15] Scanning electron microscope images of two samples are shown in Fig. 1.

When a external magnetic field is applied perpendicular to the superconducting planes and the torsional axis, the change in the resonant frequency $\Delta\nu_r$ of the oscillator is proportional to the applied field times the magnetization of the sample and is independent of the excitation amplitude in the linear-response regime.[16]

In Fig. 2, we plot the results obtained for $R_s=13.5$ μm under two different protocols. In the field-cooled (FC) pro-

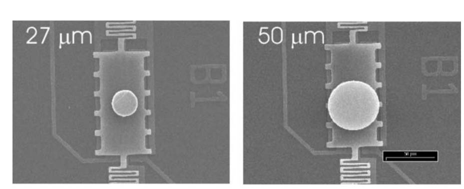

FIG. 1. Scanning electron microscope images of two samples mounted in the silicon micro-oscillators. Left (right) sample of radius 13.5(25) μm. Scale bar: 50 μm.

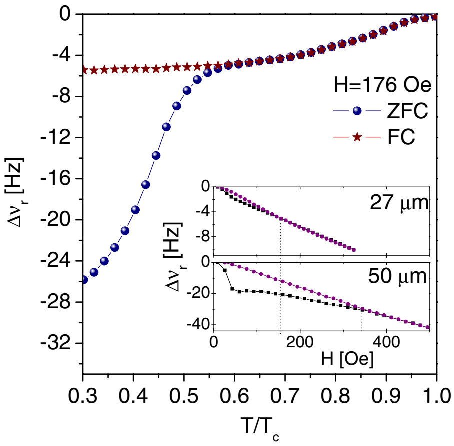

FIG. 2. (Color online) Change in the resonant frequency of the oscillator as a function of temperature ($t=T/T_C$) for the sample of radius 13.5 μm with an magnetic field of 176 Oe applied perpendicularly to the superconducting planes in a ZFC (dots) and FC (stars) experiment. Inset: $\Delta\nu_r$ as a function of H at $T/T_C=0.58$ for two sample diameters. These measurements allow the determination of the size-dependent irreversibility line.

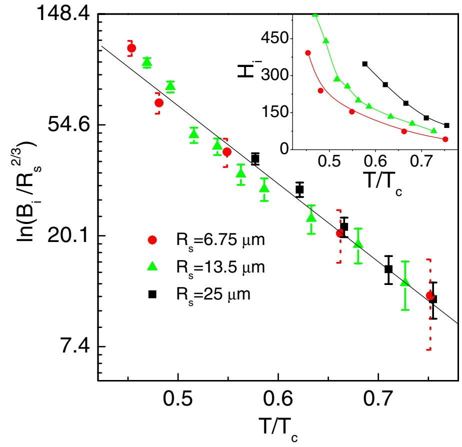

FIG. 3. (Color online) Scaling of the irreversibility field (Ref. 23) through the identification of the Larkin radius with the sample radius at the onset of irreversibility. Inset: nonscaled data for samples of different radii.

tocol, we cool the sample below its critical temperature at an applied field of 176 Oe while registering $\Delta\nu_r$ (upper curve). In the zero-field-cooled (ZFC) protocol, we cool the sample at zero field up to the lowest temperature, apply the same field as before, and then measure $\Delta\nu_r$ while warming up the sample (lower curve). The onset of irreversibility can be defined at the merging of both curves. Similar data can be obtained from fields loops at constant temperature as shown in the inset of Fig. 2. Two features can be readily observed: the clear size dependence of the irreversibility line and the wide spanning in temperature of the reversible state compared to that of bulk samples.[17] Phenomenologically, reversibility is reached when thermal fluctuations overcome the strongest pinning mechanism present in the sample. It has been argued that in this material, geometrical[18,19] or surface barriers[20] were responsible for irreversibility. Several aspects of the data point against these as the cause of the irreversibility. Geometrical and surface barriers decreases as the aspect ratio (diameter/thickness) increases.[21] Our data shows the opposite behavior, irreversibility is enhanced as the sample aspect ratio grows as can be directly seen in the inset of Fig. 2. Moreover, our data does not comply with the predictions given for the temperature dependence of the geometrical barrier and its scaling with the first penetration field.[21] As it can observed in the inset of Fig. 2, it does not comply with the expected shape of the magnetization loops for a surface barrier characterized by a zero magnetization (i.e., $\Delta\nu_r \approx 0$ in our case) in the returning field branch.[22] We show, in the following, that these puzzling finite-size effects can be explained, however by a different and more fundamental mechanism.

The observation that the irreversibility lines we find appear well below the melting line suggests the intriguing existence of a reversible vortex solid. Since such a conclusion is inconsistent with the fact that pinning (and thus irreversibility) is always relevant at large enough length scales,[10–12] we shall analyze the onset of irreversibility as the finite-size crossover at the vortex-lattice Larkin length. For an applied field parallel to the c axis, and neglecting the compression modulus contribution, we can use the Larkin-Ovchinikov perturbative result,[10]

$$W(\mathbf{r}) \approx r_p^2 \epsilon^4 \left[\frac{a_0}{l_c}\right]^3 \left[\frac{R^2}{\lambda^2} + \frac{a_0^2 L^2}{\lambda^4}\right]^{1/2}, \qquad (1)$$

where the so-called single-vortex collective pinning length l_c absorbs the effective pinning strength,[11] $a_0 = (\frac{2\phi_0}{\sqrt{3}B})^{1/2}$ is the lattice constant, λ the penetration length, and ϵ the anisotropy parameter.[11] At zero temperature, $r_p \equiv \xi$ for point impurities, with ξ the vortex core radius. At high temperatures however, fast futile thermal vortex motion induces a growth in r_p, thus effectively smoothing the microscopic disordered potential. Equation (1), which is valid for $L_c \geq L > \lambda^2/\epsilon a_0$ and $R_c \geq R > \lambda/\epsilon$ (we neglect the dispersivity in the tilt modulus[10,11]), yields

$$R_c \approx \frac{\lambda}{\epsilon^4}\left[\frac{l_c}{a_0}\right]^3, \quad L_c \approx \frac{\lambda}{a_0} R_c \qquad (2)$$

for the transverse and longitudinal Larkin lengths, respectively. Pinning, metastability and thus irreversibility (and the failure of the perturbation theory) sets in at these length scales. Above R_c and/or L_c, glassy properties are manifested. In principle, reversible behavior can be thus recovered in samples of dimensions $L_s \times R_s$, such that $R_s < R_c$ and $L_s < L_c$, provided that they still contain a large number of vortices $R_s/a_0 \gg 1$. Assuming that this is the situation in our samples, we can get the radius-dependent irreversibility line $B_i(T,R_s)$ from the equality $R_s = R_c(B_i,T)$, assuming that $L_c > L_s$. Using Eq. (2) we get, $B_i(T,R_s) \approx \phi_0[\epsilon^4/\lambda]^{2/3} l_c(T)^{-2} R_s^{2/3}$, where we have made the temperature dependence of the different parameters explicit (we neglect the temperature dependence of λ and use its average value). In Fig. 3, we show that the predicted scaling $B_i \sim R_s^{2/3}$ produces a good collapse of the irreversibility point as a function of temperature, for different R_s. This supports our identification of the sample size R_s at the onset of irreversibility with the Larkin radius, although the condition $R_s > \lambda/\epsilon$ is not strictly satisfied for all our samples. Note also that since $R_s \gg L_s$, the assumption $L_c = R_c\lambda/a_0 > L_s$ is automatically satisfied, as $R_s/L_s > a_0/\lambda$ for our measurements. In Fig. 3, we also show that our results can be well described by the expression, $B_i(R_s,T) \sim R_s^{2/3} \exp(-2T/T_0)$, with a characteristic temperature $T_0 \approx 25$ K. Interestingly, Wang *et al.*[13] have reported size effects at the second magnetization peak in BSCCO controlled by the same exponential temperature dependence, with a characteristic temperature of 22.5 K, very close to our value. In our calculations, the temperature dependence of B_i is exclusively attributed to the parameter l_c, as $l_c(T) \sim \exp(T/T_0)$. In order to grasp the physical meaning of this result, we can assume that l_c represents [as it indeed does in Eq. (1) at zero temperature] the Larkin length of an isolated vortex at finite T. An exponential sensitivity $\exp[CT^{\alpha}]$ is consistent with the marginality of the pinning of an elastic string in a three-dimensional disordered medium, and it has been predicted, with C a constant and α an exponent which depends on the precise nature of the disorder correlator function.[11,24] In particular, the value $\alpha=1$ has been predicted[24] for high-T_c superconductors for vortex displacements u satisfying $\xi < u < \lambda$, suitable for our case. More interestingly, the value of T_0 we get is very close to the one observed in creep[25] ($\sim$20 K), ac-transverse permeability[26–28] ($\sim$22 K), and critical current measurements[27] ($\sim$20 K) in samples of the same material but with radii one and two orders of magnitude bigger than ours, using the expected relations of these different quantities with l_c.[11] Being $T_0 \sim U_{pc}$, with U_{pc} the single pancake pinning energy,[25] the anomalies observed near T_0 are commonly attributed to the crossover between a strong zero-dimensional pinning regime (when l_c becomes smaller than the layer spacing and thus pancakes pin individually) and a weak three-dimensional (3D) pinning regime (see Kierfield[29] for a recent discussion). The temperature dependence of l_c, and the fact that for $T > T_0$, we can use 3D weak collective pinning [Eq. (2)], both relying on bulk pinning properties rather than on border effects, support our identification of the Larkin radius.

Our empirical estimate $R_c(B,T) \approx R_s[B/B_i(R_s,T)]^{3/2} \sim B^{3/2}\exp(T/2T_0)$, for the weak nondispersive pinning regime, implies that in the phase-space region we analyze, very big BSCCO samples are necessary to achieve the vortex-matter thermodynamic limit for $T_0 < T \lesssim T_m$, with T_m the BG melting temperature. This is relevant for the predicted crossover in the vortex-lattice roughness, from the RM to the asymptotic BG regime[12] at the characteristic scale $R_a = R_c(a_0/r_p)^{1/\zeta}$, with $\zeta \sim 0.2$ the random-manifold roughness exponent.[30,31] If we can roughly identify the renormalized pinning range with the thermally induced displacement, $r_p^2 \sim \langle u^2 \rangle_{th}$,[11] the Lindemann criterion with constant $c_L \sim 0.2$ lead us to an upper bound for the pinning range, $r_p^2 \leq c_L^2 a_0^2$. We thus get $R_a \gtrsim R_c c_L^{-1/\zeta} \sim 10^3 R_c$. If we evaluate this expression at T and $B \sim B_i(T,R_s)$, where we have shown that $R_c = R_s$ for our three samples, we find that R_a is on the order of 1 cm. Our naive estimate thus indicates that the asymptotic logarithmic growth, characteristic of the BG phase, can be only achieved in huge samples in the region of the phase space we analyze. This striking result seems to be however consistent with magnetic decoration experiments[30] displaying the random-manifold roughness up to distances $R \approx 80a_0$, for which $W(R) \approx 0.05a_0^2$ in the range $B \approx 70-120$ G. Note that the naive extrapolation of the latter to R_a, such that $W(R_a) = a_0^2$ gives $R_a \sim O(mm)$, in fair quantitative agreement with our previous estimate. These results indicate the remarkable possibility of detecting, in normal samples, the crossover from the RM to the BG regime at temperatures $T_0 < T$ below the irreversibility line as a finite-size crossover when $R_a(B,T) \approx R_s$. This would provide an independent experimental tool, different from neutron diffraction,[32] magnetic decorations,[30] or creep measurements,[25] to test the predicted geometrical features of the BG phase[12] in these materials.

We have experimentally determined, in a direct way, the minimum size of an elastic vortex lattice that can be pinned by weak disorder, by analyzing finite-size crossover effects in micron-sized high-T_c superconductors. This kind of study, complemented with micron-scale transport measurements can lead to a better understanding of the rich multiscale physics of pinned vortex lattices, and of the universal properties they share with other pinned elastic manifolds.

We thank G. Nieva for providing us with the BSCCO raw crystals, and S. Bustingorry, V. Correa, F. de la Cruz, D. Domínguez, and T. Giamarchi for very useful discussions. M.I.D., A.B.K., and H.P. are CONICET researchers.

[1] G. Grüner, Rev. Mod. Phys. **60**, 1129 (1988).
[2] S. Brazovskii and T. Nattermann, Adv. Phys. **53**, 177 (2004).
[3] E. Y. Andrei, G. Deville, D. C. Glattli, F. I. B. Williams, E. Paris, and B. Etienne, Phys. Rev. Lett. **60**, 2765 (1988).
[4] S. Lemerle, J. Ferre, C. Chappert, V. Mathet, T. Giamarchi, and P. Le Doussal, Phys. Rev. Lett. **80**, 849 (1998).
[5] P. J. Metaxas, J. P. Jamet, A. Mougin, M. Cormier, J. Ferre, V. Baltz, B. Rodmacq, B. Dieny, and R. L. Stamps, Phys. Rev. Lett. **99**, 217208 (2007).
[6] K.-J. Kim, J.-C. Lee, S.-M. Ahn, K.-S. Lee, C.-W. Lee, Y. J. Cho, S. Seo, K.-H. Shin, S.-B. Choe, and H.-W. Lee, Nature (London) **458**, 740 (2009).
[7] T. Tybell, P. Paruch, T. Giamarchi, and J. M. Triscone, Phys. Rev. Lett. **89**, 097601 (2002).
[8] S. Moulinet, A. Rosso, W. Krauth, and E. Rolley, Phys. Rev. E **69**, 035103(R) (2004).
[9] L. Ponson, D. Bonamy, and E. Bouchaud, Phys. Rev. Lett. **96**, 035506 (2006).
[10] A. I. Larkin and Y. N. Ovchinnikov, J. Low Temp. Phys. **34**, 409 (1979).
[11] G. Blatter, M. V. Feigel'man, V. B. Geshkenbein, A. I. Larkin, and V. M. Vinokur, Rev. Mod. Phys. **66**, 1125 (1994).

[12] T. Giamarchi and P. Le Doussal, Phys. Rev. B **52**, 1242 (1995).
[13] Y. M. Wang, M. S. Fuhrer, A. Zettl, S. Ooi, and T. Tamegai, Phys. Rev. Lett. **86**, 3626 (2001).
[14] C. A. Bolle, V. Aksyuk, F. Pardo, P. L. Gammel, E. Zeldov, E. Bucher, R. Boie, D. J. Bishop, and D. R. Nelson, Nature (London) **399**, 43 (1999).
[15] M. I. Dolz, W. Bast, D. Antonio, H. Pastoriza, J. Curiale, R. D. Sánchez, and A. G. Leyva, J. Appl. Phys. **103**, 083909 (2008).
[16] M. Dolz, D. Antonio, and H. Pastoriza, Physica B **398**, 329 (2007).
[17] H. Pastoriza, M. F. Goffman, A. Arribere, and F. de la Cruz, Phys. Rev. Lett. **72**, 2951 (1994).
[18] D. Majer, E. Zeldov, and M. Konczykowski, Phys. Rev. Lett. **75**, 1166 (1995).
[19] E. Zeldov, A. I. Larkin, V. B. Geshkenbein, M. Konczykowski, D. Majer, B. Khaykovich, V. M. Vinokur, and H. Shtrikman, Phys. Rev. Lett. **73**, 1428 (1994).
[20] M. S. James, S. T. Stoddart, S. J. Bending, S. Aukkaravittayapun, P. J. King, and M. Henini, Phys. Rev. B **56**, R5771 (1997).
[21] E. H. Brandt, Phys. Rev. B **60**, 11939 (1999).
[22] A. M. Campbell and J. E. Evetts, Adv. Phys. **21**, 199 (1972).
[23] Our approximation $\mathbf{B} \sim \mathbf{H}$ gives a correction factor of about 10% in the values obtained for R_c.
[24] M. Müller, D. A. Gorokhov, and G. Blatter, Phys. Rev. B **64**, 134523 (2001).
[25] M. Nideröst, A. Suter, P. Visani, A. C. Mota, and G. Blatter, Phys. Rev. B **53**, 9286 (1996).
[26] M. F. Goffman, J. A. Herbsommer, F. de la Cruz, T. W. Li, and P. H. Kes, Phys. Rev. B **57**, 3663 (1998).
[27] V. F. Correa, J. A. Herbsommer, E. E. Kaul, F. de la Cruz, and G. Nieva, Phys. Rev. B **63**, 092502 (2001).
[28] V. F. Correa, G. Nieva, and F. de la Cruz, Phys. Rev. Lett. **87**, 057003 (2001).
[29] J. Kierfeld, Phys. Rev. B **69**, 144513 (2004).
[30] P. Kim, Z. Yao, C. A. Bolle, and C. M. Lieber, Phys. Rev. B **60**, R12589 (1999).
[31] S. Bogner, T. Emig, and T. Nattermann, Phys. Rev. B **63**, 174501 (2001).
[32] T. Klein, I. Joumard, S. Blanchard, J. Marcus, R. Cubitt, T. Giamarchi, and P. Le Doussal, Nature (London) **413**, 404 (2001).