Expansion of Vortex Cores by Strong Electronic Correlation in La$_{2-x}$Sr$_x$CuO$_4$ at Low Magnetic Induction

R. Kadono,¹ W. Higemoto,¹ A. Koda,¹ M. I. Larkin,² G. M. Luke,² A.T. Savici,² Y. J. Uemura,² K. M. Kojima,³ T. Okamoto,³ T. Kakeshita,³ S. Uchida,³ T. Ito,⁴ K. Oka,⁴ M. Takigawa,⁵ M. Ichioka,⁵ and K. Machida⁵

¹Institute of Materials Structure Science, Research Organization (KEK), Tsukuba, Ibaraki 305-0801, Japan
²Physics Department, Columbia University, New York, NY10027, USA
³Department of Superconductivity, University of Tokyo, Tokyo 113-8656, Japan
⁴National Institute of Advanced Industrial Science and Technology, Tsukuba, Ibaraki 305-8562, Japan
⁵Department of Physics, Okayama University, Okayama 700-8530, Japan

(Dated: November 14, 2018)

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The structural and electronic properties of the flux line lattice (FLL) state in high-$T_c$ cuprates have been extensively studied by microscopic techniques, including scanning tunneling spectroscopy (STS), small angle neutron scattering (SANS), nuclear magnetic resonance (NMR), and muon spin rotation ($\mu$SR). Rich information has been provided not only on the superconductivity itself, but also on the unique FLL state as a ‘vortex matter’ real-

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FLL was determined by the in-plane penetration depth to magnetic field. The initial muon spin polarization was perpendicular to the crystal c-axis after setting the magnetic field at every field point (i.e., superconducting state, as the crystal flow cryostat and cooled from a temperature above Tc). Some results on LB15 were previously reported. The specimen was loaded onto a He-flow cryostat and cooled from a temperature above Tc in measurements with an applied field above 0.4 T, which do not exhibit such additional relaxation due to random flux pinning. In the present study, we analyzed data obtained only for H = 0.2 T which do not exhibit such additional relaxation in order to avoid any complication due to magnetism. We will report on high-field μSR measurements in all of these specimens in a separate paper. In addition, we observed an additional component undergoing a fast spin relaxation which became more prominent at higher fields. In zero field, μSR results from the specimen LA15 exhibited slightly enhanced spin relaxation below ~10 K, which was virtually absent in LB15. The spin relaxation rate in LA15 at 3.5 K was greater than that at 41 K by 0.026(2) μs⁻¹, indicating a weak remnant magnetism. In addition, we observed an additional component undergoing a fast spin relaxation in LA15 in measurements with an applied field above 0.4 T, which became more prominent at higher fields. In the present study, we analyzed data obtained only for H = 0.2 T which do not exhibit such additional relaxation, in order to avoid any complication due to magnetism. We will report on high-field μSR measurements in all of these specimens in a separate paper, and discuss the possible relation to field-induced magnetism suggested by neutron diffraction. Since muons stop randomly along the length scale of FLL, the time evolution of complex muon polarization \( \mathcal{P}(t) \) provides a random sampling of the internal field distribution, \( \mathbf{B}(r) = (0, 0, B(r)) \),

\[
\mathcal{P}(t) = P_x(t) + iP_y(t) = \exp(-\sigma_p^2/2) \int_{-\infty}^{\infty} n(B) \exp(i\gamma_B t - i\psi) dB,
\]

where \( P_{x,y}(t) \) is proportional to the time-dependent \( \mu^+ - e^- \) decay asymmetry \( A_{x,y}(t) \) deduced from the corresponding sets of positron counters \( N_{x,y}(t) \) after subtracting constant backgrounds estimated from data at earlier times \( t < 0 \),

\[
A_x(t) = N_0(t)/N_{\pi/2}(t) = A_0 P_x(t),
\]

\[
A_y(t) = \alpha N_{\pi/2}(t)/N_{3\pi/2}(t) = A_0 P_y(t),
\]

with \( \alpha = [{N_0(0)N_{3\pi/2}(0)}]/[{N_\pi(0)N_{\pi/2}(0)}] \), \( A_0 \) is the normalized instrumental asymmetry, \( \sigma_p \) is the additional relaxation due to random flux pinning, \( n(B) \) is the spectral density for the internal field defined as a spatial average of the delta function, \( \gamma_B \) is the muon gyromagnetic ratio \( (=2\pi\times135.53 \text{ MHz/T}) \), and \( \psi \) is the initial phase. These equations indicate that the real amplitude of the Fourier-transformed muon precession signal corresponds to \( n(B) \) with an appropriate correction of \( \sigma_p \). In the modified London model, \( B(r) \) is approximated as the sum of the magnetic induction from isolated vortices to yield

\[
B(r) = B_0 \sum_K \frac{e^{-iK \cdot r}}{1 + K^2 \lambda^2} F(K, \xi_v),
\]

where \( B_0 \) is the radial distribution of the supercurrent density, \( \xi_v \) is the coherence length for the internal field defined as a spatial average of the Fourier-transformed muon precession signal, \( \lambda \) is the in-plane penetration depth to magnetic field, and \( \xi_v \) is the coherence length of the delta function. In the present study, we analyzed data obtained only for H = 0.2 T which do not exhibit such additional relaxation in order to avoid any complication due to magnetism. We will report on high-field μSR measurements in all of these specimens in a separate paper, and discuss the possible relation to field-induced magnetism suggested by neutron diffraction.
where $K$ are the vortex reciprocal lattice vectors, $B_0$ ($\sim H$) is the average internal field, $A$ is the London penetration depth, and $F(K, \xi_v)$ is a nonlocal correction term with $\xi_v$ being the cutoff parameter for the magnetic field distribution; it must be stressed that $\xi_v$ is a parameter to describe the electromagnetic response of a vortex, thereby not necessarily equivalent to the core radius $\rho_v$. Considering small anisotropy predicted by the theory, we assumed an isotropic vortex core and associated field distribution near the vortex center.

**TABLE I:** A comparison of the central field between points above and below the superconducting transition temperature $T_c$, where the external field was set to $6 \ T$ above $T_c$ and then field-cooled. The relative precision of the field ($|\Delta f/f|$) was better than $10^{-4}$ for all cases. Note that the effect of field inhomogeneity and/or the Knight shift is larger at higher fields. Therefore, the above number gives an upper bound for the uncertainty of the central field.

| Samples  | $f = \gamma_0 B_0/2\pi$ (MHz) | $|\Delta f/f|$ |
|----------|-------------------------------|-------------|
| LA15     | 818.9905(22)                  | 818.9902(66) | 0.0(9) $\times 10^{-5}$ |
| LB13     | 813.0765(15)                  | 813.0533(41) | 2.9(5) $\times 10^{-5}$ |
| LB15     | 812.8405(17)                  | 812.7686(30) | 8.9(3) $\times 10^{-5}$ |
| LB19     | 813.0679(12)                  | 813.0384(19) | 3.6(3) $\times 10^{-5}$ |

As summarized in Table I, the relative change of $B_0$ due to field inhomogeneity and/or the Knight shift was less than $10^{-4}$ throughout measurements below 50 K, and therefore $B_0$ was fixed to the value determined in the normal state ($T > T_c$) for the analysis of data below $T_c$. This well-defined $B_0$, together with the strongly asymmetric feature of $n(B)$ against $B_0$, allowed us to deduce the physical parameters including $\lambda$ and $\xi_v$ without much ambiguity. Moreover, we have found that the deduced values of $\xi_v$ (and thereby $\rho_v$) were virtually independent of the apex angle ($\theta$) of FLL, while $\lambda$ showed a considerable dependence on $\theta$. This feature can be readily understood by considering the fact that the change in the apex angle has least effect on the field distribution of single vortex near the center (corresponding to the high-field end of $n(B)$), while it does modify the overlap of field distribution between the vortices. Based on this robustness of the analysis result against $\theta$, we assumed triangular FLL ($\theta = \pi/3$) throughout the following data analysis for simplicity. We have also found that the $\mu$SR spectra in LSCO at lower fields are much better reproduced by the Lorentzian cutoff, $F(K, \xi_v) = \exp(-\sqrt{2}K \xi_v)$, compared with the conventional Gaussian cutoff, $F(K, \xi_v) = \exp(-K^2 \xi_v^2/2)$. Figure 1 shows a typical example of the measured decay asymmetry, $A_2(t) \propto P_2(t)$, observed at 0.2 T in specimen LA15 together with the result of a fitting analysis. The reduced chi-square $\chi^2/N_f$ (where $N_f$ denotes the number of degrees of freedom) is close to the ideal value of $\approx 1$ when the Lorentzian cutoff is applied to Eq. 2, while it becomes much worse as $\chi^2/N_f \approx 1.61$ with the Gaussian cutoff. A similar tendency was commonly observed in all other samples. This observation is consistent with the result of a theoretical analysis which showed that the Lorentzian cutoff is indeed a better approximation at the low-field limit. Accordingly, all of the $\mu$SR time spectra were analyzed by comparing the data with the time evolution calculated by Eqs. 1 of the Lorentzian cutoff, where $A_0$, $\psi$, $\sigma_p$, $\xi_v$, and $\lambda$ were free parameters while $B_0$ and $\theta$ were always fixed.

While the fitting analysis of the $\mu$SR data is performed entirely in time domain, one can reconstruct $B(r)$ by Eq. 4 using the physical parameters deduced from the fitting analysis. An example of the reconstructed field distribution, $B(r) = B_z(r)$, along the radial direction from the core center to a saddle point is shown in Fig. 2a, together with the corresponding $n(B)$ (Fig. 2b) and supercurrent density, $J(r) = |\text{rot} B(r)|$ (Fig. 2c). As shown in Fig. 2d, the core radius defined by $J(J_{\text{rot}}) = J_{\text{max}}$ (where $J_{\text{max}}$ denotes the maximum of $J(r)$) is considerably larger than the magnetic cutoff parameter $\xi_v$, indicating the need for a special precaution in interpreting $\xi_v$ directly as the core radius. However, the result of a data analysis for various fields/temperatures indicates that $\rho_v$ is always proportional to $\xi_v$. It has been estimated that $\xi_v \approx 23 \ \AA$ for an optimally doped sample ($H_{c2} \approx 62 \ \text{T}$ for $x = 0.15$). Provided that $\rho_v$ is independent of the field and determined by $\xi_v\text{GL} = 0.6 \sim 0.8 \xi_v\text{GL}$ as predicted by theory, then $\xi_v$ must be smaller than 13 $\AA$ to reproduce the corresponding core radius (see Fig. 2c). However, an attempt to fit the data assuming $\xi_v = 13 \ \AA$ completely fails to reproduce the data; the reduced chi-square $\chi^2/N_f$ was 2.76 at its best when other parameters (i.e., $A_0$, $\psi$, $\sigma_p$, and $\lambda$) were set free to minimize $\chi^2/N_f$. This is primarily because, as evident in Fig. 2d, the peak of $n(B)$ shifts significantly to a lower field, which cannot be compensated by any other parameters within the current model. Note that $B_0$ is determined with a relative precision better than $10^{-4}$ so that the shift of the peak in Fig. 2d ($\approx 0.2 \ \text{mT}$) is readily discernible. Thus, it is inferred from these results that $\rho_v$ is about three times as large as $\xi_v\text{GL}$ near the lower critical field. This tendency is qualitatively in line with the theoretical prediction, although the magnitude is yet to be explained.

The deduced vortex core radius for samples LB13, LB15, and LB19, as a function of the temperature, are shown in Figs. 3a-c together with the results for LA15. It exhibits a slight increase with decreasing temperature for $x = 0.13$ and 0.15, while an opposite tendency is observed for $x = 0.19$. No such anomaly has been reported for YBCO. Note that the conventional theory predicts a behavior opposite to what is actually observed in the former cases (see Fig. 3a). The values deduced from the data in LA15 are in good agreement with those in LB15, indicating that the influence of the remnant magnetism found in LA15 is not significant at $H = 0.2 \ \text{T}$ and $T \geq 15 \ \text{K}$. The penetration depth extrapolated to 0 K was 2559(50) $\AA$ in LA15, 2446(6) $\AA$, 2460(4) $\AA$, and 2051(7) $\AA$ in LB13, LB15, and LB19, respectively. The
relatively large error for LA15 is due to the influence of the remnant magnetism below 10 K. We also note that \( \sigma_p \) was typically 0.4~0.5 \( \mu s^{-1} \) at the lowest temperature in all specimens, showing a common tendency of gradual decrease with increasing temperature. This is understood as the thermal depinning of FLL from random pinning centers.

For understanding the \( T \)-dependence of \( \rho_v \), we performed a model calculation based on Bogoliubov-de Gennes theory, where both the \( d \)-wave superconductivity and the spatially modulated antiferromagnetic (AF) spin correlations are simultaneously considered by incorporating a pairing interaction to the standard Hubbard model for two-dimensional square lattice\(^{24}\). As can be seen in Fig. 3b, the core radius at which the superconducting order parameter reaches 60\% of the bulk value, is strongly enhanced when the electronic correlation is present (i.e., \( U/t > 0 \), with \( U \) and \( t \) being the respective on-site Coulomb energy and transfer matrix element to the nearest neighboring sites in the Hubbard model), while it obeys the prediction by Kramer and Pesch when \( U/t = 0 \). The \( T \)-dependence of \( \rho_v \) for \( U/t > 0 \) qualitatively agrees with those observed for \( x = 0.13 \) and 0.15. In the model, this feature is due to AF spin correlations induced at the core sites, leading to a reduction of the localization length \( \chi \), which was typically 0.4~0.5 \( \mu m \) at the lowest temperature (i.e., \( \sigma_p \)).

In order to obtain the extrapolated value \( \rho_{v}(0) \), the data in Figs. 3a–3c were analyzed by the \( \chi^2 \)-minimization method using an analytical approximation of Fig. 3d,

\[
\rho_v(T) = \rho_v(0)(1+c_1(T)+c_2(\exp(-T)-1)+c_3(\exp(T)-1))
\]

under a condition \( c_1 < c_2 < 0 < c_3 \) with \( T = T/T_c \). The result of fitting analysis is summarized in Table II. The values of \( \rho_v(0) \) are plotted against the Sr concentration \( x \) in Fig. 4. (The data of LA15 were not used for this analysis because they were available only for \( T \geq 15 \) K, leading to large uncertainty in the fitting.) As the hole doping progresses with increasing \( x \), the core size decreases monotonically. We note that almost equivalent result was obtained by an independent analysis using a linear function \( \rho_v(T) = \rho_v(0)(1+cT) \) for \( T < 0.5T_c \).

We shall compare this behavior with those of other relevant parameters. Firstly, \( \rho_v(0) \) and \( T_c \) have no simple correlations, since \( T_c \) shows a maximum at \( x = 0.15 \). Second, the superconducting gap, \( \Delta \), does not scale with \( \rho_v(0) \), if \( \Delta \) scales with \( T_c \). In contrast, \( \rho_v(0) \) in YBCO is larger for the lower \( T_c \), thus following the simple scaling law\(^{10}\). It has been noticed in LSCO that the gap (or pseudo-gap) increases with decreasing hole concentration in the optimum-to-underdoped region. In this sense, the (pseudo-) gap value and \( \rho_v(0) \) vary along with each other. This behavior is opposite to that expected in a simple BCS superconductor, where the coherence length is inversely proportional to the gap value as \( \xi = hv_F/\pi \Delta \), where \( v_F \) denotes the Fermi velocity. In \( La_{2-x}Sr_xCuO_4 \), the Fermi velocity has been known to have little dependence on \( x \). It is interesting to note that the model calculation shown in Fig. 3d appears to simulate the doping dependence if we assume that the incipient AF correlations become weaker (\( U/t \to 0 \)) as the doping proceeds. A similar result has been reported for the correlation length, \( \xi_{Zn}(x) \), in Zn-doped \( La_{2-x}Sr_xCuO_4 \) over which the pairing is suppressed\(^{22}\).

Finally, we point out the potential link between our results and those of recent neutron scattering in LSCO where a field-induced quasistatic antiferromagnetism has been suggested\(^{12,13,14}\). Although it has not been clearly identified as being due to the vortex cores, they found that a long-range AF correlation is recovered in the mixed state under moderate magnetic fields of a few Tesla. A similar situation is also suggested in YBCO\(^{23,24}\), BSCCO\(^{25}\), and \( YBa_2Cu_3O_7 \). The enhanced vortex core radius at lower fields in LSCO may be interpreted as a precursor of such a quasistatic correlation. Besides our model calculation, theories including those based on the \( t-J \) model\(^{27,28}\) or SO(5) symmetry\(^{29}\) also predict the de-

| Samples | \( \rho_v(0) \) (A) | \( c_1 \) | \( c_2 \) | \( c_3 \) | \( \chi^2/N_f \) |
|---------|----------------|--------|--------|--------|-------------|
| LB13    | 120.9 (4.6)    | -6.36(3) | -2.69(4) | 2.65(2) | 7.5/5       |
| LB15    | 95.0(7.3)      | -6.42(5) | -2.68(6) | 2.72(4) | 9.8/7       |
| LB19    | 73.3(11.5)     | -5.51(6) | -2.14(9) | 2.91(4) | 18.5/3      |

TABLE II: Parameter values for Eq. (6) deduced by fitting data in Fig. 3a–3c.
magnetic field or impurity atoms. When the superconducting order parameter is suppressed at the center of vortices, one would expect the emergence of AF spin correlations around the vortex cores, which may be strongly dynamical at lower fields.

In summary, we have found in LSCO that the vortex core radius at 0.2 T is about three times as large as that estimated from $H_{c2}$. The core size increases monotonically with decreasing Sr doping with an unusual temperature dependence for $x = 0.13$ and 0.15. Compared with the cases of other cuprates, these features more strongly suggest a possibility that $\rho_v$ is influenced by the two-dimensional AF correlations in this system.

We would like to thank the staff of TRIUMF for their technical support during the experiment. This work was partially supported by a Grant-in-Aid for Scientific Research on Priority Areas and a Grant-in-Aid for Creative Scientific Research from the Ministry of Education, Culture, Sports, Science and Technology of Japan, NSF Grant DMR-01-02752 and CHE-01-17752 at Columbia, and by NSERC of Canada.

\textsuperscript{1} Also at School of Mathematical and Physical Science, The Graduate University for Advanced Studies

\textsuperscript{1} Present address: Department of Physics and Astronomy, McMaster University, Hamilton, ON L8P4N3, Canada

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FIG. 4: Doping dependence of $\rho_v$, extrapolated to 0 K at $H = 0.2$ T and that of the superconducting transition temperature ($T_c$) determined by the Meissner effect. The lines are guides for the eye.
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