Nodeless superconductivity in the infinite-layer electron-doped cuprate superconductor $Sr_{0.9}La_{0.1}CuO_2$

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(Received 21 March 2008; revised manuscript received 25 April 2008; published 19 May 2008)

We report on measurements of the in-plane magnetic penetration depth λ_{ab} in the infinite-layer electron-doped high-temperature cuprate superconductor $Sr_{0.9}La_{0.1}CuO_2$ by means of muon-spin rotation. The observed temperature and magnetic field dependences of λ_{ab} are consistent with the presence of a substantial *s*-wave component in the superconducting order parameter, in good agreement with the results of tunneling, specific heat, and small-angle neutron scattering experiments.

DOI: 10.1103/PhysRevB.77.184512 PACS number(s): 74.72.Jt, 74.25.Jb, 76.75.+i

I. INTRODUCTION

The symmetry of the superconducting energy gap is one of the essential issues for understanding the mechanism of high-temperature superconductivity. For hole-doped high-temperature cuprate superconductors (HTS's) it is commonly accepted that the superconducting energy gap has *d*-wave symmetry, as indicated, e.g., by tricrystal experiments,[1] although overwhelming evidence is now accumulating for a multicomponent [$(d+s)$-wave] gap.[2–4] For electron-doped HTS's, however, no consensus has been reached on this issue so far. A number of experiments, including angular resolved photoemission,[5,6] scanning superconducting quantum interference device (SQUID) measurements,[7] Raman scattering,[8] and magnetic penetration depth studies,[9,10] point to a *d*-wave, or, more generally, to a nonmonotonic *d*-wave gap (with the gap maximum between the nodal and antinodal points on the Fermi surface) in electron-doped $Nd_{2-x}Ce_xCuO_4$ and $Pr_{2-x}Ce_xCuO_4$. On the contrary, a state corresponding to an *s*-or a nonmonotonic *s*-wave gap symmetry was reported for similar compounds and infinite-layer $Sr_{1-x}La_xCuO_2$ in tunneling,[11,12] Raman scattering,[13] penetration depth,[14] small-angle neutron scattering,[15] and specific heat[16] studies. Results of Biswas *et al.*[17] and Skinta *et al.*[18] suggest that there is a *d*- to *s*-wave transition across optimal doping in $Nd_{2-x}Ce_xCuO_4$ and $Pr_{2-x}Ce_xCuO_4$. The two-gap picture was also introduced in Refs. 19 and 20 in order to explain the unusual behavior of the magnetic penetration depth and Raman spectra.

Here we report a study of the in-plane magnetic field penetration depth λ_{ab} in the infinite-layer electron-doped superconductor $Sr_{0.9}La_{0.1}CuO_2$ by means of transverse-field muon-spin rotation (TFμSR). This compound belongs to the family of electron-doped HTS's $(\mathrm{Sr}L)\mathrm{CuO}_2$ (L=La,Sm,Nd,Gd) with the so-called infinite-layer structure.[21,22] It has the simplest crystal structure among all HTS's, consisting of an infinite stacking of CuO_2 planes and (Sr, L) layers. The charge reservoir block, commonly present in HTS's, does not exist in the infinite-layer structure. It also has stoichiometric oxygen content without vacancies or interstitial oxygen,[23] which is a general problem for most of the electron- and hole-doped HTS's. In the present study $\lambda_{ab}^{-2}(T)$ was reconstructed from the measured temperature dependences of the μSR linewidth by applying the numerical calculations of Brandt.[24] λ_{ab} was found to be almost field independent, in contrast to the strong magnetic field dependence observed in the hole-doped HTS's, suggesting that there are no nodes in the superconducting energy gap of $Sr_{0.9}La_{0.1}CuO_2$. The temperature dependence of λ_{ab}^{-2} was found to be well described by anisotropic *s*-wave as well as by two-gap models ($d+s$ and $s+s$). In the case of the two-gap $d+s$ model the contribution of the *d*-wave gap to the total superfluid density was found to be of the order of 15%. Our results imply that a substantial *s*-wave component in the superconducting order parameter is present in $Sr_{0.9}La_{0.1}CuO_2$, in agreement with previously reported results of tunnelling,[12] specific heat,[16] and small-angle neutron scattering experiments.[15]

II. EXPERIMENTAL DETAILS

Details on the sample preparation for $Sr_{0.9}La_{0.1}CuO_2$ can be found elsewhere.[25] The TFμSR experiments were performed at the πM3 beamline at the Paul Scherrer Institute (Villigen, Switzerland). The sintered $Sr_{0.9}La_{0.1}CuO_2$ sample was field cooled from above T_c to 1.6 K in a series of fields ranging from 50 mT to 0.64 T.

In the transverse-field geometry, the local magnetic field distribution $P(B)$ inside the superconducting sample in the mixed state, probed by means of TFμSR, is determined by the values of the coherence length ξ and the magnetic field penetration depth λ. In extreme type-II superconductors

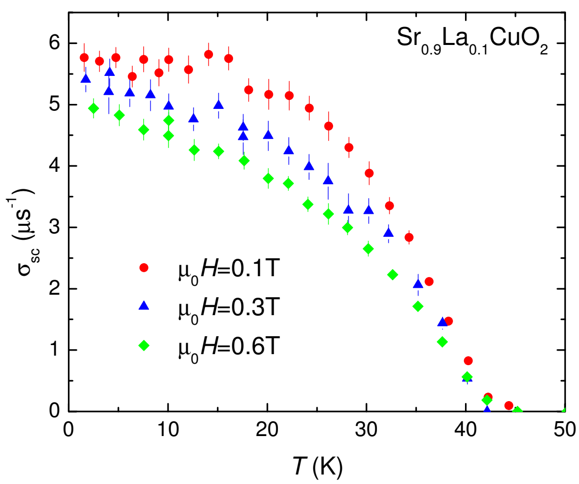

FIG. 1. (Color online) Temperature dependences of $\sigma_{sc} \propto \lambda_{ab}^{-2}$ of $Sr_{0.9}La_{0.1}CuO_2$ measured in magnetic fields of 0.1, 0.3, and 0.6 T.

($\lambda \gg \xi$) $P(B)$ is almost independent of ξ and the second moment of $P(B)$ becomes proportional to $1/\lambda^4$.[24,26] In order to describe the asymmetric local magnetic field distribution $P(B)$ in the superconductor in the mixed state, the analysis of the data was based on a two-component Gaussian fit of the μSR time spectra:[27]

$$P(t) = \sum_{i=1}^{2} A_i \exp(-\sigma_i^2 t^2/2)\cos(\gamma_\mu B_i t + \phi). \quad (1)$$

Here A_i, σ_i, and B_i are the asymmetry, the relaxation rate, and the mean field of the ith component, $\gamma_\mu = 2\pi \times 135.5342$ MHz/T denotes the muon gyromagnetic ratio, and ϕ is the initial phase of the muon-spin ensemble. The total second moment of the μSR line was derived as[27]

$$\langle \Delta B^2 \rangle = \frac{\sigma^2}{\gamma_\mu^2} = \sum_{i=1}^{2} \frac{A_i}{A_1 + A_2}\left[\frac{\sigma_i^2}{\gamma_\mu^2} + \left(B_i - \frac{A_i B_i}{A_1 + A_2}\right)^2\right]. \quad (2)$$

The superconducting part of the square root of the second moment σ_{sc} was further obtained by subtracting the contribution of the nuclear moments σ_{nm} measured at $T > T_c$ as $\sigma_{sc}^2 = \sigma^2 - \sigma_{nm}^2$.[28]

The following issue is important for the interpretation of the experimental data: The sample used in the experiment was a nonoriented sintered powder. In this case an effective averaged penetration depth λ_{eff} can be extracted. However, in highly anisotropic extreme type-II superconductors (like HTS's) λ_{eff} is dominated by the in-plane penetration depth so that $\lambda_{eff} \simeq 1.31\lambda_{ab}$.[29]

III. RESULTS AND DISCUSSION

Figure 1 shows the temperature dependences of $\sigma_{sc} \propto \lambda_{ab}^{-2}$ measured after field-cooling the sample from far above T_c in $\mu_0 H = 0.1$, 0.3, and 0.6 T. Two features are clearly seen: (i) In the whole temperature region (from $T \simeq 1.6$ K up to T_c) $\sigma_{sc}(T,H)$ decreases with increasing field; and (ii) the curvature of $\sigma_{sc}(T)$ changes with field. With decreasing temperature σ_{sc} at $\mu_0 H = 0.1$ T first increases and then becomes T independent for $T \lesssim 15$ K, while σ_{sc} for both $\mu_0 H = 0.3$ and 0.6 T increases continuously in the whole range of temperatures. It is also seen that the low-temperature slope of $\sigma_{sc}(T)$ is larger at the higher fields.

The decrease of σ_{sc} with increasing magnetic field is caused by the overlapping of the vortices by their cores, leading to a reduction of the field variance in the superconductor in the mixed state. As shown by Brandt,[24] at magnetic inductions $B/B_{c2} \lesssim 0.1$ (B_{c2} is the second critical field), overlapping of vortex cores may be neglected and vortex properties are independent of the applied magnetic field. Only the vortex density is changed. At higher magnetic inductions vortices start to overlap with their cores. Consequently, not only the vortex density, but also the properties of the individual vortices, become magnetic field dependent. Therefore, the different temperature behavior of σ_{sc} can be explained by the fact that the vortex core size, which is generally assumed to be equal to the coherence length $\xi = (\Phi_0/2\pi B_{c2})^{0.5}$, increases with increasing temperature (decreasing B_{c2}). Thus higher temperature [bigger reduced field $b(T) = B/B_{c2}(T)$] would correspond to larger overlapping vortex cores, leading to a stronger reduction of σ_{sc}. Consequently, the correcting factor between σ_{sc} and the magnetic field penetration depth is *not constant*, as is generally assumed, but depends on the reduced field b:[24]

$$\sigma_{sc}(b) \;\; (\mu s^{-1}) = A(b)[\lambda_{eff}^{-2} \;\; (nm^{-2})]. \quad (3)$$

Here $A(b)$ is a correcting factor, which for superconductors with a Ginzburg-Landau parameter $\kappa = \lambda/\xi \geq 5$ in the range of fields $0.25/\kappa^{1.3} \lesssim b \leq 1$ can be obtained analytically as $A(b) = 4.83 \times 10^4 (1-b)[1 + 1.21(1-\sqrt{b})^3]$ μs^{-1} nm^2.[24]

Equation (3) requires that, in order to derive λ from measured $\sigma_{sc}(T)$, the temperature dependence of B_{c2} must be taken into account. Since in our experiments the measured λ_{eff} is determined by the in-plane component of the magnetic penetration depth λ_{ab} (see above), $B_{c2}(T)$ has to be measured with the magnetic field applied parallel to the crystallographic c axis. The temperature dependence of $B_{c2}^{\parallel c}$ presented in Fig. 2(a) was obtained from measurements of the reversible magnetization on the c-axis-oriented powder by using the Landau-Ott scaling approach.[30] The solid line represents a fit of $B_{c2}^{\parallel c}(T)$ using the Werthamer-Helfand-Hohenberg (WHH) model[31] with $B_{c2}^{\parallel c}(0) = 10.7(2)$ T. Note that $B_{c2}^{\parallel c}(0) = 10.7(2)$ T obtained in the present study is close to that reported in the literature.[32,33]

The correcting factor A [see Fig. 2(b)] was calculated within the framework of the numerical Ginzburg-Landau approach developed by Brandt[24] in the following way. First, A was calculated as a function of the reduced field $b = B/B_{c2}$. The inset in Fig. 2(b) shows $A(b) = \sigma\lambda^2$ [see Eq. (3)] for a superconductor with $\kappa = 15$, 20, and 25. Vertical lines correspond to $b = B/B_{c2}(0)$ for the fields used in the present study (0.05, 0.1, 0.3, and 0.64 T). Second, by using the calculated WHH form of $B_{c2}(T)$ [red solid line in Fig. 2(a)], values of $A(T, B = \text{const})$ for each particular field B were reconstructed. Figure 2(b) shows the calculated $A(T,B)$ at $B = 0.05$, 0.1, and 0.64 T for $\kappa = 15$, 20, and 25. Note that the influence of κ is only important at the lowest magnetic field ($B = 0.05$ T), while at the higher fields the effect of κ becomes almost

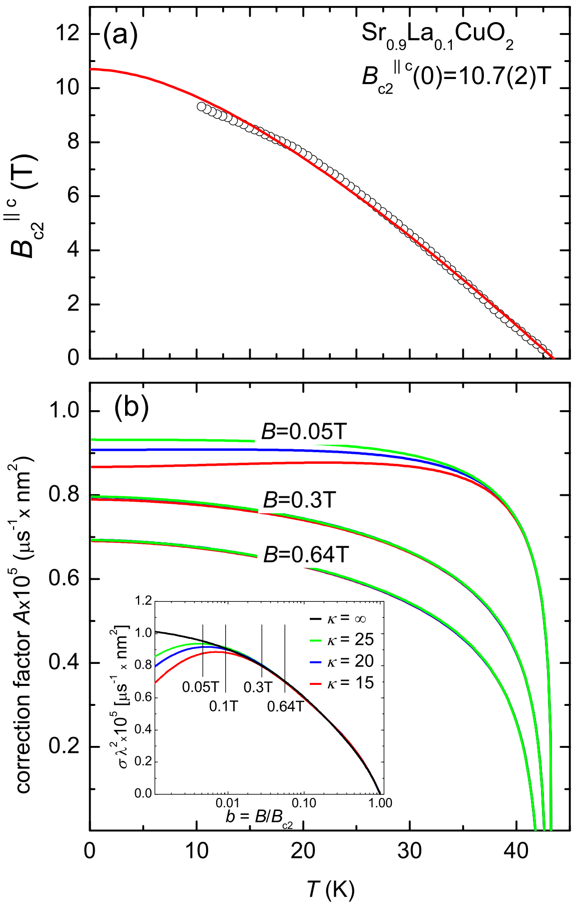

FIG. 2. (Color online) (a) Temperature dependence of the second critical field $B_{c2}^{\parallel c}(T)$ of $Sr_{0.9}La_{0.1}CuO_2$ obtained from measurements of the reversible magnetization on c-axis-oriented powder by using the Landau-Ott approach (Ref. 30). The solid line is the fit of the Werthamer-Helfand-Hohenberg (WHH) model (Ref. 31) to the data with $B_{c2}^{\parallel c}(0)=10.7(2)$ T. (b) Temperature dependence of the correcting factor A at B=0.05, 0.1, and 0.6 T for $\kappa=\lambda/\xi=15$ (green line), 20 (blue line), and 25 (red line). The inset shows $A(b)=\sigma\lambda^2$ obtained from numerical calculations using the model of Brandt (Ref. 24).Vertical lines mark the values of $B/B_{c2}(0)$ for B=0.05, 0.1, 0.3, and 0.64 T.

negligible. This allows us to estimate the absolute value of $\lambda_{\rm eff}$ at low temperatures, which was found to be $\lambda_{\rm eff} \simeq 120$ nm, in good agreement with previous results.[34] Bearing in mind that $\lambda_{\rm eff}\simeq 1.31\lambda_{ab}$ (Ref. 29) and $B_{c2}^{\parallel c}=10.7$ T [see Fig. 2(a)], the Ginzburg-Landau parameter was estimated to be $\kappa\simeq 17$. This value of κ was used in the following calculations.

The normalized $\lambda_{ab}^{-2}(T,B)/\lambda_{ab}^{-2}(0,B)$ curves reconstructed from the measured $\sigma_{\rm sc}(T,B)$ are shown in Fig. 3. The inset shows $\lambda_{ab}(0,B)$ obtained from the linear extrapolation of $\lambda_{ab}^{-2}(T,B)$ at $T<10$ K to zero temperature. It should be emphasized here that it makes no sense to introduce any unique value of λ for each particular magnetic field, since only the zero-field value of λ has a physical meaning.[35] The same statement holds for the temperature behavior of λ^{-2}. This implies that if the theory used to reconstruct $\lambda(T)$ from $\sigma_{\rm sc}(T)$ measured in various magnetic fields is correct, then the corresponding $\lambda(T)$ should *coincide*. The data presented

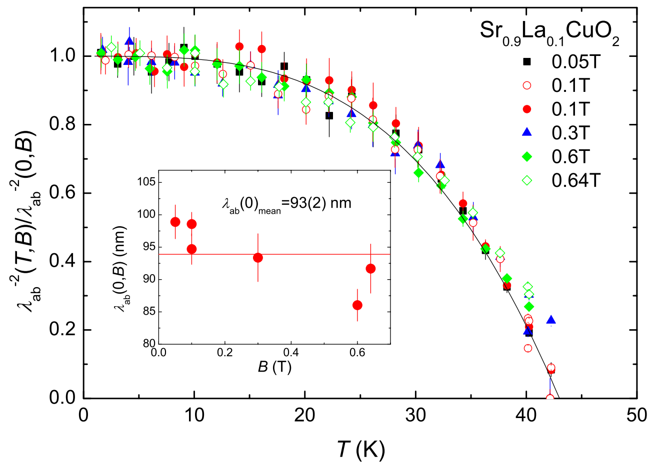

FIG. 3. (Color online) $\lambda_{ab}^{-2}(T,B)$ of $Sr_{0.9}La_{0.1}CuO_2$ normalized to its value at T=0 reconstructed from $\sigma_{\rm sc}(T)$ measured at μ_0H=0.05, 0.1, 0.3, 0.6, and 0.64 T. The solid line is a guide to the eye. The inset shows $\lambda_{ab}(0)$ as a function of the applied filed.

in Fig. 3 reveal that this is the case. The normalized $\lambda_{ab}^{-2}(T,B)/\lambda_{ab}^{-2}(0,B)$ curves merge together and the value of λ_{ab} at T=0, presented in the inset, is almost independent of the magnetic field. The small increase of $\lambda_{ab}(0)$ at low fields can be explained by the pinning effects which can lead to a reduction of the second moment of $P(B)$ in a powder HTS's at fields smaller than 0.1 T (see, e.g., Refs. 28 and 36).

Figure 4 shows $\lambda_{ab}^{-2}(T)$ reconstructed by means of the above described procedure from $\sigma_{\rm sc}(T)$ for one representative field value B=0.1 T. The data in Fig. 4 were analyzed by using single-gap and two-gap models, assuming that the superconducting energy gaps have the following symmetries: d wave (a), anisotropic d wave (b), $d+s$ wave (c), s wave (d), anisotropic s wave (e), and $s+s$ wave (f). The analysis for both the d-wave and the isotropic s-wave gaps was made in the clean ($\xi\gg l$, where l is the meanfree path) and the dirty ($\xi\ll l$) limits. The cases of anisotropic s- and d-wave gaps, as well as $(s+s)$- and $(d+s)$-wave gap symmetries were analyzed in the clean limit only.

In the clean limit the temperature dependence of the magnetic penetration depth λ was calculated within the local (London) approximation ($\lambda\gg\xi$) by using the following equation:[2,37]

$$\left.\frac{\lambda^{-2}(T)}{\lambda^{-2}(0)}\right|_{\rm clean}=1+\frac{1}{\pi}\int_0^{2\pi}\int_{\Delta(T,\varphi)}^{\infty}\left(\frac{\partial f}{\partial E}\right)\frac{E\,dE\,d\varphi}{\sqrt{E^2-\Delta(T,\varphi)^2}}. \tag{4}$$

Here $\lambda^{-2}(0)$ is the zero-temperature value of the magnetic penetration depth, $f=[1+\exp(E/k_BT)]^{-1}$ is the Fermi function, φ is the angle along the Fermi surface ($\varphi=\pi/4$ corresponds to a zone diagonal), and $\Delta(T,\varphi)=\Delta_0\tilde{\Delta}(T/T_c)g(\varphi)$ (Δ_0 is the maximum gap value at T=0). The temperature dependence of the gap is approximated by $\tilde{\Delta}(T/T_c)=\tanh\{1.82[1.018(T_c/T-1)]^{0.51}\}$.[38] The function $g(\varphi)$ describes the angular dependence of the gap and is given by $g^s(\varphi)=1$ for the s-wave gap, $g^d(\varphi)=|\cos(2\varphi)|$ for the d-wave

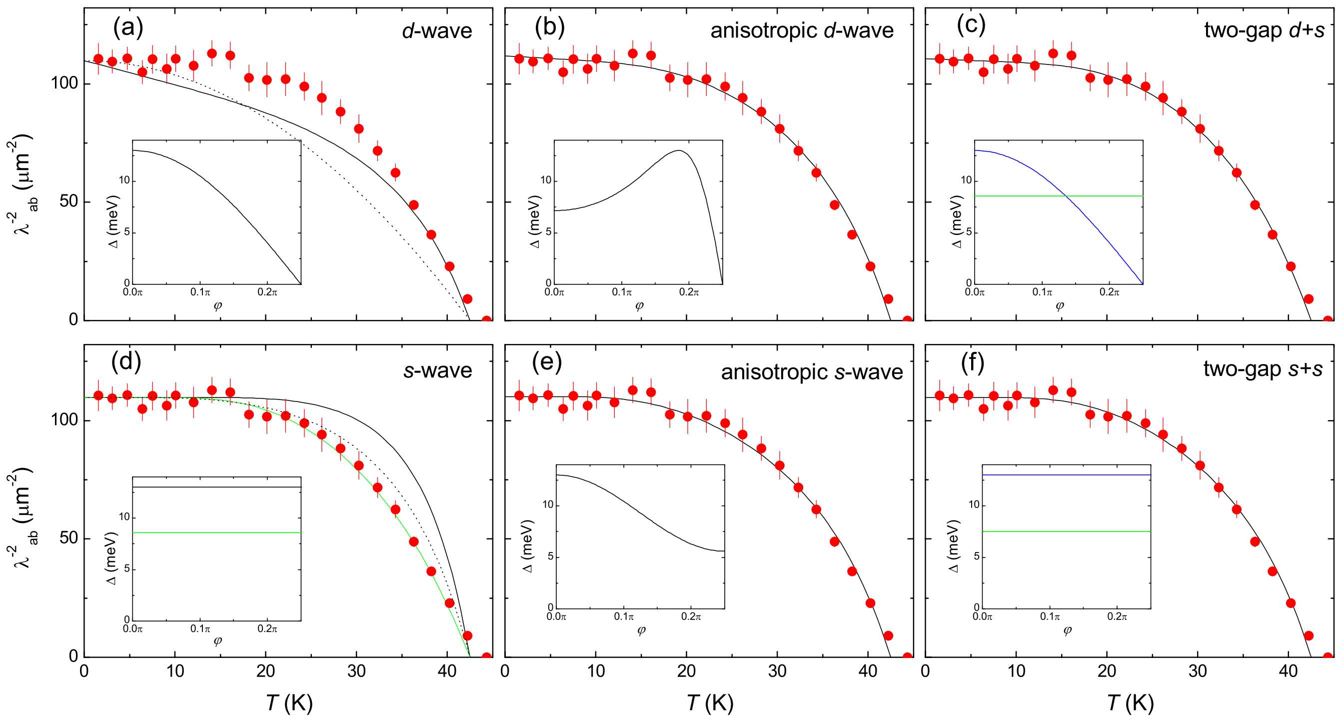

FIG. 4. (Color online) Temperature dependence of λ_{ab}^{-2} of $Sr_{0.9}La_{0.1}CuO_2$ reconstructed from $\sigma_{sc}(\mu_0 H=0.1\ \mathrm{T})$. The curves were obtained within the following models of the gap symmetries: (a) clean (solid line) and dirty *d* wave (dotted line), (b) anisotropic *d* wave, (c) two-gap *d*+*s*, (d) clean (solid black, $\Delta_0=13$ meV, and solid green, $\Delta_0=8.6$ meV) and dirty *s* wave (dotted line), (e) anisotropic *s* wave, (f) two-gap *s*+*s*. The corresponding angular dependences of the gaps are shown in the insets.

gap, $g^{s\,\mathrm{an}}(\varphi)=(1+a\cos 4\varphi)/(1+a)$ for the anisotropic *s*-wave gap,[11] and $g^{d\,\mathrm{an}}(\varphi)=3\sqrt{3a}\cos 2\varphi/2(1+a\cos^2\varphi)^{3/2}$ for the anisotropic *d*-wave gap.[39]

In the dirty limit $\lambda^{-2}(T)$ was obtained via[37]

$$\left.\frac{\lambda^{-2}(T)}{\lambda^{-2}(0)}\right|_{\mathrm{dirty}\ s\ \mathrm{wave}}=\frac{\Delta(T)}{\Delta(0)}\tanh\left(\frac{\Delta(T)}{2k_B T}\right), \tag{5}$$

and assuming the power law dependence

$$\left.\frac{\lambda^{-2}(T)}{\lambda^{-2}(0)}\right|_{\mathrm{dirty}\ d\ \mathrm{wave}}=1-\left(\frac{T}{T_c}\right)^n \tag{6}$$

with the exponent $n\equiv 2$,[40] for *s*- and *d*-wave gaps, respectively.

The two-gap calculations [($d+s$)- and ($s+s$)-wave] were performed within the framework of the so called α model assuming that the total superfluid density is a sum of two components:[2,38]

$$\frac{\lambda^{-2}(T)}{\lambda^{-2}(0)}=\omega\frac{\lambda^{-2}(T,\Delta_1)}{\lambda^{-2}(0,\Delta_1)}+(1-\omega)\frac{\lambda^{-2}(T,\Delta_2)}{\lambda^{-2}(0,\Delta_2)}. \tag{7}$$

Here $\Delta_1(0)$ and $\Delta_2(0)$ are the zero-temperature values of the large and the small gap, respectively, and ω ($0\leq\omega\leq 1$) is the weighting factor which represents the relative contribution of the larger gap to λ^{-2}.

The results of the analysis are presented in Fig. 4. The angular dependences of the gaps [$\Delta_0 g(\varphi)$] are shown in the corresponding insets. The maximum value of the gap $\Delta_0=13$ meV was kept fixed in accordance with the results of tunneling experiments.[12] It is obvious that simple *d*- and *s*-wave approaches with $\Delta_0=13$ meV cannot describe the observed $\lambda_{ab}^{-2}(T)$ [see Figs. 4(a) and 4(d)]. All other models, such as, e.g., both anisotropic *d* and *s* wave, and both two-gap *d*+*s* and *s*+*s* [Figs. 4(b), 4(c), 4(e), and 4(f)], as well as the single *s*-wave model with $\Delta_0=8.6$ meV [solid green line in Fig. 4(a)] describe the experimental data reasonably well. The results of the analysis are summarized in Table I.

Based on our analysis we cannot differentiate between the models presented in Figs. 4(b), 4(c), 4(e), and 4(f), and the isotropic *s*-wave model with $\Delta_0=8.6$ meV [solid green line in Fig. 4(d)]. We believe, however, that from our consideration we can exclude the case of an anisotropic *d*-wave gap [see Fig. 4(b)]. The reasons are the following. (i) The model of Brandt,[24] used to reconstruct $\lambda_{ab}^{-2}(T)$ from the measured $\sigma_{sc}(T)$, is strictly valid for conventional superconductors only. The presence of nodes in the gap makes the electrody-

TABLE I. Summary of the gap analysis of the temperature dependences of λ_{ab}^{-2} for $Sr_{0.9}La_{0.1}CuO_2$. The meaning of the parameters is explained in the text.

Model	$\Delta_{0,1}$ (meV)	a	$\Delta_{0,2}$ (meV)	ω
Clean limit *s*	8.6			
Anisotropic *d*	13[a]	3.1		
Anisotropic *s*	13[a]	0.4		
Two-gap *d*+*s*	13[a]		9.0	0.15
Two-gap *s*+*s*	13[a]		7.5	0.35

[a]From tunneling experiments (Ref. 12).

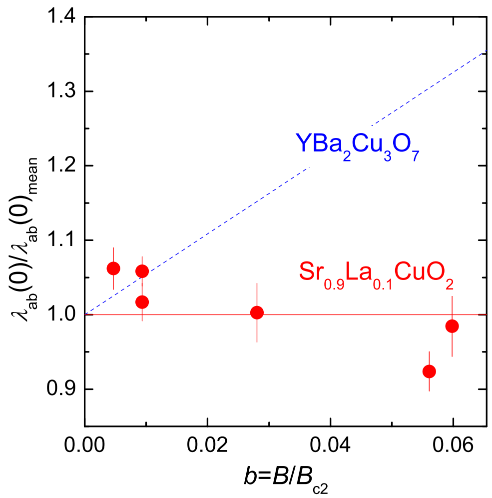

FIG. 5. (Color online) $\lambda_{ab}(0)$ of $Sr_{0.9}La_{0.1}CuO_2$ normalized to its mean value as a function of reduced magnetic field $b=B/B_{c2}$. The dashed blue line corresponds to $\lambda_{ab}(0,b)$ for $YBa_2Cu_3O_{7-\delta}$ from Ref. 43.

namics of the mixed state highly nonlocal, leading to an additional source for decreasing the superfluid density due to excitation of the quasiparticles along the nodal directions (see, e.g., Ref. 41 and references therein). In this case the proportionality coefficient A [see Eq. (3) and Fig. 2(b)] depends much more strongly on the reduced magnetic field and, as a consequence, on temperature than is expected in the case of conventional superconductors. However, the normalized $\lambda_{ab}^{-2}(T,B)/\lambda_{ab}^{-2}(0,B)$ evaluated at various fields merge together (see Fig. 3), implying that $Sr_{0.9}La_{0.1}CuO_2$ can be, indeed, described within the model developed for conventional isotropic superconductors.[24] (ii) The nonlocal and nonlinear response of a superconductor with nodes in the gap to the magnetic field leads to the fact that $\lambda_{ab}(0)$, evaluated from μSR experiments, becomes magnetic field dependent and increases with increasing field.[42] This behavior was observed in various hole-doped HTS's.[41,43,44] For comparison, in Fig. 5 we show $\lambda(0,b)$ for $Sr_{0.9}La_{0.1}CuO_2$ measured here and the one obtained by Sonier *et al.*[43] for hole-doped $YBa_2Cu_3O_{7-\delta}$. Whereas $\lambda_{ab}(0)$ for $YBa_2Cu_3O_{7-\delta}$ strongly increases with magnetic field, for $Sr_{0.9}La_{0.1}CuO_2$ it is almost field independent.

Now we want to point to differences between $Sr_{0.9}La_{0.1}CuO_2$ studied here and $Nd_{2-x}Ce_xCuO_4$ and $Pr_{2-x}Ce_xCuO_4$, electron-doped superconductors for which the d-wave symmetry of the order parameter was reported based on the results of angular resolved photoemission,[5,6] scanning SQUID,[7] Raman scattering,[8] and magnetic penetration depth studies.[9,10] First, the reversible magnetization experiments of Kim *et al.*[32] reveal that the zero-temperature c-axis coherence length in $Sr_{0.9}La_{0.1}CuO_2$ is longer then the c-axis lattice parameter. This suggests that the superconducting order parameter of one CuO_2 plane overlaps with those of neighboring CuO_2 planes for all temperatures below T_c. This would imply that the superconducting properties of the infinite-layer system are nearly three dimensional, as opposed to the quasi-two-dimensional nature of all other cuprates. Second, the superconductivity in the infinite-layer compound appears to be sensitive to the type of impurity. While nonmagnetic Zn has little effect on T_c for up to 3% concentration, strong suppression of T_c already occurs with 1% of Ni, and nearly complete suppression of T_c occurs with only 2% of Ni.[25] Thus, the response of the infinite-layer compound to impurities is different from that of other HTS's and it is more similar to the one observed in conventional superconductors. Third, isotropic s-wave gap symmetry of $Sr_{0.9}La_{0.1}CuO_2$ was reported by Chen *et al.*[12] based on the analysis of tunneling data. This statement was further confirmed in specific heat[16] and small-angle neutron scattering experiments.[15]

IV. SUMMARY AND CONCLUSIONS

Muon-spin rotation measurements were performed on the electron-doped high-temperature cuprate superconductor $Sr_{0.9}La_{0.1}CuO_2$ ($T_c\simeq 43$ K). $\lambda_{ab}^{-2}(T)$ was reconstructed from the measured temperature dependence of the μSR linewidth by using numerical calculations of Brandt.[24] The main results are as follows. (i) The absolute value of the in-plane magnetic penetration depth λ_{ab} at $T=0$ was found to be $\lambda_{ab}(0)=93(2)$ nm. (ii) λ_{ab} is independent of the magnetic field, in contrast to the strong magnetic field dependence observed in hole-doped HTS's. This suggests that there are no nodes in the superconducting energy gap of $Sr_{0.9}La_{0.1}CuO_2$. (iii) The temperature dependence of λ_{ab}^{-2} was found to be inconsistent with isotropic s-wave as well as with a d-wave symmetry of the superconducting energy gap, in both the clean and the dirty limits. (iv) $\lambda_{ab}^{-2}(T)$ is well described by anisotropic s-wave and two-gap models ($d+s$ and $s+s$). In the case of the two-gap $d+s$ model, the contribution of the d-wave gap to the total superfluid density was estimated to be $\simeq 15\%$. To conclude, a substantial s-wave component in the superconducting order parameter is present in $Sr_{0.9}La_{0.1}CuO_2$, in agreement with previously reported results of tunneling,[12] specific heat,[16] and small-angle neutron scattering experiments.[15]

ACKNOWLEDGMENTS

This work was partly performed at the Swiss Muon Source (SμS), Paul Scherrer Institute (PSI, Switzerland). The authors are grateful to A. Amato and D. Herlach for providing the instrumental support, and D. G. Eshchenko for help during the μSR experiments. This work was supported by the Swiss National Science Foundation and the K. Alex Müller Foundation, and in part by the SCOPES Grant No. IB7420-110784 and the EU Project CoMePhS.

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