Thermal conductivity of La$_{2-x}$Sr$_x$CuO$_4$ (0.05 ≤ x ≤ 0.22)

J-Q Yan, J-S Zhou and J B Goodenough

Texas Materials Institute, ETC 9.102, The University of Texas at Austin, 1 University Station, C2201, Austin, TX 78712, USA
E-mail: jqyan@mail.utexas.edu

New Journal of Physics 6 (2004) 143
Received 29 June 2004
Published 26 October 2004
Online at http://www.njp.org/
doi:10.1088/1367-2630/6/1/143

Abstract. A systematic study of the thermal conductivity of single-crystal samples of the La$_{2-x}$Sr$_x$CuO$_4$ superconductive system and comparison with that of La$_{1.60-x}$Nd$_{0.40}$Sr$_x$CuO$_4$ (x = 0.12, 0.15 and 0.20) have demonstrated that this measurement is a useful indirect probe of mechanisms that suppress phonons. The data distinguish octahedral-site rotational or charge-order fluctuations above a structural order–disorder transition and two-phase fluctuations associated with locally cooperative atomic displacements within two-dimensional sheets that segregate hole-rich and hole-poor regions in a mixed-valent system. The former suppress phonons in both the basal plane and along the apical axis, whereas the latter only suppress phonons in the basal plane. The data support a spinodal phase segregation below room temperature into the parent and the superconductive phases in the underdoped compositional range as well as the superconductive and metallic phases in the overdoped compositions; they also support the existence of locally cooperative bond-length fluctuations in the normal state of the superconductive phase that prevent the formation of a percolative matrix capable of supporting phonons. Restoration of the phonons below $T_c$ signals a long-range, dynamic ordering of the bond-length fluctuations which implies stabilization of a travelling charge-density wave with possible hybridization of electrons and phonons below $T_c$ to give heavy vibrons that pair in the superconductive phase.
1. Introduction

The system \( \text{La}_{2-x}\text{Sr}_{x}\text{CuO}_4 \) (LSCO), \( 0 \leq x \leq 0.3 \), is the simplest \( p \)-type copper oxide exhibiting superconductivity. The Ruddlesden–Popper \((n = 1)\) structure consists of \( \text{CuO}_2 \) sheets alternating with \( \text{(LaO)}_2 \) rocksalt layers on traversing the \( c \)-axis. It has no charge reservoir layer between the two \( \text{LaO} \) sheets of the rock-salt layers, which makes the fraction of holes per \( \text{Cu} \) atom in the \( \text{CuO}_2 \) sheets unambiguously \( p = x \) provided the oxygen stoichiometry is maintained. Moreover, it spans the entire superconductive compositional range between the antiferromagnetic insulator, parent phase \((x = 0)\) and the overdoped, metallic and nonsuperconductive phase \((x > 0.28)\). Also, the \( \text{La}_2\text{CuO}_4+\delta \) system has the same parent phase, and stoichiometric single crystals can be grown by the floating-zone method in an infrared image furnace. It has, therefore, been extensively and intensively studied in a search for the origin of the high-\( T_c \) phenomenon in the copper oxides.

Although most theorists initially believed that electron–phonon interactions could be responsible for the formation of Cooper pairs and a phenomenological theory of spin–fluctuation interactions led to the successful prediction of d-wave pairing [1, 2], it was pointed out early that the cross-over from a Mott–Hubbard antiferromagnetic insulator to an itinerant-electron metal would be first-order and would, therefore, involve a phase segregation and/or bond-length fluctuations associated with a strong electron–lattice interaction [3]. Others [4, 5] recognized that a phase change from localized to itinerant electronic behaviour would lead to a quantum electrical point above which spin and bond-length fluctuations would give unusual physical properties. Kulic [6] has provided a critical review in which he cites evidence that electron–phonon interactions appear to play a more important role than the electron–spin interactions in the formation of Cooper pairs. However, his review does not consider the available data on thermal conductivity, perhaps because there has been no systematic study of the thermal conductivity of \( \text{La}_{2-x}\text{Sr}_{x}\text{CuO}_4 \) and related systems. In this paper, we attempt to fill this void.

2. Experimental details

Single crystals were grown by the travelling-solvent, floating-zone technique in an IR image furnace as reported elsewhere [7]. Before measurement of physical properties, the as-grown crystals were annealed for a long time in a flowing gas of selected \( P_{\text{O}_2} \) in order to (i) relieve any stress formed during crystal growth, cutting or polishing, (ii) homogenize the crystal and (iii) tune...
up the oxygen content. The physical properties are sensitive to the oxygen stoichiometry and also to chemical inhomogeneities. The hole concentration was established by the following measurements and their comparison with published data on stoichiometric samples of the same composition: (i) the superconductive transition temperature \( T_c \) as determined from the temperature dependence of the diamagnetic signal measured with a SQUID in a zero-field-cooled (ZFC) mode; (ii) lattice parameters from x-ray powder diffraction of pulverized single-crystal samples with Si as internal standard; (iii) the in-plane and c-axis thermoelectric power measured in a home-made set-up; (iv) the in-plane and c-axis electrical resistivity \( \rho(T) \) obtained by a four-probe method and a rectangular-bar-shaped sample. The in-plane and c-axis thermal conductivities \( \kappa_{ab}(T) \) and \( \kappa_c(T) \) were measured only when the above physical properties were consistent with well-characterized literature data. A steady-state heat-transfer method was used to measure thermal conductivity as described elsewhere [8]. We found that the quantitative values of our results varied somewhat with three experimental conditions: (i) how the sample is connected to the sink? (ii) how the thermocouples are attached to the surface? and (iii) what current is used to generate the heat flow? However, in all cases, the temperature dependences were similar. Therefore, all measurements were made under the same conditions; and in our discussion, the principal focus is on the temperature dependence of \( \kappa_{ab}(T) \) and \( \kappa_c(T) \).

3. Results

Figure 1 shows the measured \( T_c \) of our La\(_{2-x}\)Sr\(_x\)CuO\(_4\) samples and compares our single-crystal data with the polycrystalline data of Takagi et al [9]. The dips at \( x \approx 1/8 \) and \( x \approx 5/24 \) are characteristic of lattice instabilities associated with competitive stripe phases that become pinned with a little Zn doping [10].

Figure 1. The measured \( T_c \) as a function of Sr doping in La\(_{2-x}\)Sr\(_x\)CuO\(_4\). The dashed line shows the \( T_c \) of ceramic samples from [9].
Figure 2. The c-axis thermal conductivity $\kappa_c(T)$ for crystals of La$_{2-x}$Sr$_x$CuO$_4$: (a) for underdoped compositions, (b) overdoped compositions. Solid arrows in (a) indicate the LTO/HTT structural transition temperature $T_t$, while those in (b) indicate $T_c$. Inset in (a) shows the ZFC magnetization measured in a 35 Oe magnetic field.

Figure 2(a) shows $\kappa_c(T)$ data for underdoped samples $0.08 \leq x \leq 0.15$ and figure 2(b) for the overdoped samples $0.19 \leq x \leq 0.22$ of La$_{2-x}$Sr$_x$CuO$_4$. The measured $\kappa_c(T)$ is at least two orders of magnitude larger than the electronic contribution $\kappa_e(T)$ obtained from the electrical resistivity $\rho_c(T)$ with the Wiedemann–Franz law. Therefore, phonons are the dominant heat carriers determining $\kappa_c(T)$ for these quasi-two-dimensional (2D) electronic conductors. The arrows in figure 2(a) mark the temperature of a low-temperature orthorhombic (LTO) to high-temperature tetragonal (HTT) structural transition as determined by an independent neutron-diffraction measurement [11]. The LTO phase is derived from the HTT phase by cooperative...
rotation of the octahedral CuO$_6$ sites about a [110] axis. Attention is drawn to the fact that $\kappa_c(T)$ obeys closely a phononic $1/T$ law in the LTO phase, but it exhibits characteristics of an amorphous phase in a finite temperature interval above the LTO/HTT transition at $T_t$. Also noteworthy is the monotonic decrease with increasing $x$ of the low-temperature phononic peak, which becomes completely suppressed in the $x \geq 0.21$ samples where the HTT phase is stabilized to lowest temperatures. Moreover, the overdoped samples all exhibit a suppression of $\kappa_c(T)$ in the interval $60 \text{K} < T < 150 \text{K}$. These data are consistent with those of Nakamura et al [12], who reported, in addition, $\kappa_{ab}(T)$ and $\kappa_c(T)$ curves for $x = 0.30$ that showed a recovery of the low-temperature phononic peak in both $\kappa_{ab}(T)$ and $\kappa_c(T)$ in this metallic, non-superconductive sample.

Figure 3 shows $\kappa_{ab}(T)$ for (a) $0.05 \leq x \leq 0.10$, (b) $0.12 \leq x \leq 0.15$ and (c) $0.19 \leq x \leq 0.22$. The electronic component $\kappa_e(T)$ of the conductive $x = 0.22$ sample, as estimated from the Wiedemann–Frantz law and the measured $\rho_{ab}(T)$, is seen to contribute a few per cent to the measured $\kappa_{ab}(T)$. In the compositional range (figure 3(a)), $\kappa_{ab}(T)$ has a phonon-like behaviour as does $\kappa_c(T)$, but the peak at about 20 K is larger and $\kappa_{ab}(T)$ decreases more sharply with increasing temperature in the interval $T > 80 \text{K}$ than does $\kappa_c(T)$. The sharper drop in $\kappa_{ab}(T)$ signals the presence of an extra scattering mechanism in the $a$–$b$ plane that is absent along the $c$-axis in the underdoped samples. The extra scattering below 200 K becomes more evident in the compositional range (figure 3(b)) where a minimum in $\kappa_{ab}(T)$ is found near 50 K.

The values of $T_t$ for the LTO/HTT transition as determined by $\kappa_{ab}(T)$ agree with those determined from $\kappa_c(T)$, which shows that our crystals were homogeneous since the samples used for the two measurements were from different parts of the ingot. Above $T_t$, $\kappa_{ab}(T)$ shows a positive temperature coefficient characteristic of an amorphous material as was also noted for $\kappa_c(T)$. We also call attention to the partial suppression of the low-temperature phononic peak in the interval $0.10 \leq x \leq 0.15$. An upturn at $T_c$ in $\kappa_{ab}(T)$ with decreasing temperature signals that in-plane phonons are at least partially recovered in the superconductive phase.

In the overdoped compositional range (figure 3(c)), $0.19 \leq x \leq 0.22$, $\kappa_{ab}(T)$ has a temperature dependence similar to that of $\kappa_c(T)$, but with a larger magnitude. From the dashed line in figure 3(c), it is apparent that the electronic component $\kappa_e(T)$ cannot account for a suppression of $\kappa_{ab}(T)$, similar to that of $\kappa_c(T)$, in the interval $50 \text{K} < T < 150 \text{K}$. A similar additional scattering mechanism for in-plane and $c$-axis phonons is operative in the overdoped samples in the range $50 \text{K} < T < 150 \text{K}$.

4. Discussion

We discuss our data with reference to the phase diagram of figure 4 from [3]. The first point to note is that $T_N$ falls off significantly with increasing hole concentration $p = x$ from $T_N = 325 \text{K}$ at $x = 0$ to below 4 K at $x = 0.02$. The sharp decrease in $T_N$ with $x$ signals a remarkable disruption of the long-range magnetic order and hence any magnon contribution to the thermal conductivity $\kappa_{ab}(T)$ below $T_N$ in the CuO$_2$ sheets. This observation is consistent with the sensitivity of the higher-temperature hump in $\kappa_{ab}(T)$ of La$_2$CuO$_{4+\delta}$ to the oxygen stoichiometry, decreasing sharply with $\delta$ [7]. This disruption of the long-range magnetic order has been attributed to the formation above $T = T_F$ of one-hole correlation bags containing 5–6 Cu centres in a CuO$_2$ sheet rather than small polarons and to a spinodal phase segregation below $T_F$ into the parent phase and the superconductive phase (see below).
Figure 3. In-plane thermal conductivity $\kappa_{ab}(T)$ of La$_{2-x}$Sr$_x$CuO$_4$ single crystals. The superconducting transition temperature $T_c$ and the LTO/HTT structural transition temperature $T_t$ are indicated by solid arrows in (b) and (c).

Strikingly, both $\kappa_{ab}(T)$ and $\kappa_c(T)$ show a transition from phonon conduction to a glass-like thermal conductivity at the long-range LTO/HTT phase-transition temperature $T_t$ as determined by neutron diffraction. However, XAFS measurements [13] at the La site have determined that short-range local tilting of the CuO$_6$ octahedra persists well into the HTT phase and that the temperature-induced LTO/HTT transition is an order–disorder transition. Our $\kappa_{ab}(T)$ and $\kappa_c(T)$ data show that the short-range-cooperative tiltings of the octahedral sites fluctuate so as to destroy the phonon $\mathbf{q}$ vector as a good quantum number, thereby rendering both $\kappa_{ab}(T)$ and $\kappa_c(T)$
glass-like in a finite-temperature interval above $T_f$. Long-range ordering of the octahedral-site tilts below $T_f$ restores the phonons.

It is significant that short-range structural fluctuations are manifest in both $\kappa_{ab}(T)$ and $\kappa_c(T)$, whereas the electronic fluctuations that disrupt the magnons influence only $\kappa_{ab}(T)$ in this structure with quasi-2D electronic properties.

In the compositional range $0 < x \leq 0.10$, a temperature-independent thermoelectric power $\alpha(T)$ at temperatures $T > T_F$ was interpreted with the statistical term [14]

$$\alpha = \frac{\kappa}{e} \ln \left[ \frac{\beta(1 - c)}{c} \right],$$

in which $\beta = 2$ is the spin-degeneracy factor and $c = pN/N^* = pn$, where $n = 5.3$ is the average number of Cu centres in a correlation bag that move diffusively. The concept of a correlation bag within which the electrons are delocalized but are confined to a bag by a matrix of localized (strongly correlated) electrons represents a segregation into regions of localized and itinerant Cu-$(x^2 - y^2)^*$ electrons at the crossover from localized to itinerant electronic behaviour. The correlation bags do not collapse to small polarons due to the lowering of the kinetic energy of the cluster (bag) electrons and an elastic energy gain associated with cooperative pseudo-Jahn–Teller distortions at the several Cu centres of a bag from square to rhomboidal Cu–O bonding [15]. The correlation bags break up the long-range magnetic order in a CuO2 sheet, but they do not appear to suppress the phonons of the percolating localized-electron matrix in the LTO phase at temperature $T > T_F$.
In the system La$_2$CuO$_{4+y}$, the interstitial oxygens are mobile and, below room temperature, a phase segregation into the parent antiferromagnetic-insulator phase and a superconductive phase is accomplished by segregation of the interstitial oxygen, which is mobile to below room temperature, into a hole-poor and a hole-rich phase [16]. This phase segregation is responsible for the appearance of filamentary superconductivity in La$_2$CuO$_{4+y}$. The driving force for the phase segregation is clearly the first-order character of the transition from localized to itinerant electronic behaviour and the lowering of the kinetic energy of the electrons of the hole-rich superconductive phase in which the one-hole correlation bags coalesce into multi-hole correlation bags containing one Cu-$(x^2 - y^2)^*$ electron for every two Cu centres. Therefore, we have postulated that a similar phase segregation occurs below $T_F$ in La$_{2-x}$Sr$_x$CuO$_4$ with $0 \leq x \leq 0.10$; but in this case no ions can move diffusively below room temperature, so phase segregation can only take place by locally cooperative oxide-ion displacements that create shorter Cu–O bonds in the hole-rich regions and longer Cu–O bonds in the hole-poor regions.

The existence of short-range antiferromagnetic clusters in a low-temperature spin glass has been observed by several techniques with antiferromagnetic clusters persisting through the underdoped compositional range. In order to demonstrate that the suppression of $\kappa_{ab}(T)$, but not of $\kappa_c(T)$, below 200 K in the underdoped samples is due to the presence of a dynamic segregation between the parent and superconductive phases by locally cooperative atomic displacements, we compare our measurements of $\alpha^2(T)$ and $1/\kappa_{ab}(T)$ in figure 5. The mapping of $T_F$ as defined by the change from a temperature-independent $\alpha(T)$ and by a maximum in $1/\kappa_{ab}(T)$ is seen to hold well for $x = 0.05$, 0.06 and 0.08; the maximum in $1/\kappa_{ab}(T)$ for $x = 0.10$ is shifted by the approach of $T_c$ rather than by the onset of the suppression of $\kappa_{ab}(T)$ in the interval 50 K $< T < 250$ K. Therefore, we can correlate the transition at $T_F$ in $\alpha(T)$ with the onset of an additional phonon-scattering mechanism. This result provides strong indirect evidence for a dynamic spinodal phase segregation into hole-poor and hole-rich phases by locally cooperative atomic displacements. These phases may be the hole-poor parent phase and the hole-rich superconductive phase and/or be ordered into a competitive stripe phase that contains mobile hole-rich stripes alternating with hole-poor stripes. Since the atomic displacements responsible for this phase segregation are restricted to the CuO$_2$ sheets, only $\kappa_{ab}(T)$ is suppressed by this phase segregation. The greater suppression of $\kappa_{ab}(T)$ below 150 K in the 0.12 $\leq x \leq 0.13$ samples probably reflects the appearance of the competitive stripe phase that causes the suppression of $T_c$.

Recovery of the low-temperature phonon peak below $T_c$ is consistent with the report [17] of an abrupt increase in $\kappa(T)$ of YBa$_2$Cu$_3$O$_{7-\delta}$ on cooling through $T_c$; a higher $T_c$ allows a stronger manifestation of the phonon peak below $T_c$. We interpret the recovery of phonons below $T_c$ to reflect the establishment of long-range order of the bond-length fluctuations as would occur with a travelling charge-density wave. We note that the low-temperature peak is smaller in the 0.12 $\leq x \leq 0.13$ samples where $T_c$ is lowered by the appearance of a competitive stripe phase; therefore, the data do not support a mobile $x = 1/8$ stripe phase as the superconductive phase.

To investigate further the influence of an $x = 1/8$ stripe phase on the thermal conductivity, we turned to the La$_{1.60-\delta}$Nd$_{0.40}$Sr$_x$CuO$_4$ (LNSCO) system [7], which has a low-temperature-tetragonal (LTT) phase below a LTT/LTO transition at a temperature $70$ K $< T'_F < 80$ K in the compositional range 0.12 $\leq x \leq 0.20$. In the LTT phase, the octahedral Cu sites rotate around [100] and [010] axes in alternate CuO$_2$ sheets, and this tilting pins at $x = 1/8$ hole-rich stripes containing one Cu-$(x^2 - y^2)^*$ electron per two Cu centres in every fourth Cu–O–Cu row of a CuO$_2$ sheet [18]. Thermal conductivity shows an abrupt increase in both $\kappa_{ab}(T)$ and $\kappa_c(T)$ on
cooling through \( T'_e \). Phonons are clearly established in the LTT phase. Suppression of both \( \kappa_{ab}(T) \) and \( \kappa_c(T) \) at temperatures \( T > T'_e \) would appear to reflect lattice instabilities associated with fluctuations in the octahedral-site rotations as was shown to occur above the LTO/HTT transition in \( \text{La}_{2-x}\text{Sr}_x\text{CuO}_4 \). However, a larger suppression of \( \kappa_{ab}(T) \) in the LTO phase is consistent with Cu–O bond-length fluctuations associated with a dynamic segregation of hole-rich and hole-poor domains as in the normal state of the \( \text{La}_{2-x}\text{Sr}_x\text{CuO}_4 \) superconductive compositions; the \( x = 0.15 \) and 0.20 LNSCO samples are superconductive with a \( T_c < 15 \) K.

We return finally to the overdoped \( 0.19 \leq x \leq 0.22 \) samples of the \( \text{La}_{2-x}\text{Sr}_x\text{CuO}_4 \) system shown in figure 3(c). In these samples, the electronic component \( \kappa_e(T) \) is less than 30\% of the measured \( \kappa_{ab}(T) \). Here, we find a suppression of both \( \kappa_{ab}(T) \) and \( \kappa_c(T) \) that increases with decreasing temperature in the range \( T_c < T < 200 \) K; it signals the existence of structural fluctuations. Figure 4 shows that \( T_c \) crosses \( T_e \) in this compositional range, so the suppression of both \( \kappa_{ab}(T) \) and \( \kappa_c(T) \) signals the coexistence of fluctuating LTO and HTT phases. These fluctuations apparently mask any suppression in \( \kappa_{ab}(T) \) by fluctuating superconductive and metallic phases as a result of a spinodal phase segregation between the two. However, the recovery of the low-temperature phonon peak below \( T_c \) manifests an ordering of in-plane

Figure 5. Comparison of \( \alpha^2(T) \) (○) and \( 1/\kappa_{ab}(T) \) (+) for (a) \( x = 0.05 \), (b) \( x = 0.06 \), (c) \( x = 0.08 \) and (d) \( x = 0.10 \).
bond-length fluctuations in the superconductive phase within the LTO phase in the $x = 0.19$ and 0.20 samples; the low-temperature $\kappa_c(T)$ peak remains suppressed in the $x = 0.21$ and 0.22 samples and the $\kappa_{ab}(T)$ peak in the $x = 0.22$ sample where the HTT phase is stabilized to lowest temperatures.

5. Conclusions

From these experiments, we draw the following conclusions.

1. Thermal conductivity is a useful indirect probe of lattice instabilities that suppress phonons and can distinguish structural phase fluctuations from electronic phase fluctuation in this quasi-2D electronic system at the crossover from localized to itinerant electronic conduction.

2. Isolated polarons, correlation bags or strong-correlation fluctuations in a matrix that retains percolation do not destroy thermal conductivity by phonons.

3. Octahedral-site rotational fluctuations associated with the coexistence of two structural phases (e.g., LTO + HTT, LTT + LTO, or as shown for short-range charge order in La$_{1.65}$Sr$_{0.33}$NiO$_4$) suppress phonons within and perpendicular to the basal planes of the layered Ruddlesden–Popper A$_2$MO$_4$ structure.

4. Locally cooperative bond-length fluctuations that segregate hole-rich from hole-poor regions within mixed-valent basal planes suppress in-plane phonons where there is a percolating matrix; but long-range ordering, static or dynamic, of the bond-length fluctuations restores in-plane phonons. The $c$-axis phonons are not suppressed by fluctuations of in-plane atomic displacements.

5. In-plane magnons conduct heat. The long-range antiferromagnetic order within CuO$_2$ sheets that is needed for magnons is destroyed sharply in La$_{2-x}$Sr$_x$CuO$_4$ or La$_2$CuO$_4+$δ by the introduction of holes that occupy one-hole correlation bags containing 5–6 Cu centres and by spinodal phase segregation below $T = T_F$.

6. The thermal conductivity data for the La$_{2-x}$Sr$_x$CuO$_4$ system are consistent with the model proposed previously to interpret the properties summarized in figure 4. This model identifies the superconductive phase at $x = 1/6$ as thermodynamically distinguishable below room temperature from the parent phase on the underdoped side and the metallic phase on the overdoped side. Spinodal phase segregations occur between these phases below room temperature and competitive stripe phases may appear at $x = 1/8$ and 5/24. These low-temperature phase segregations are accomplished by locally cooperative atomic displacements within the CuO$_2$ sheets. In the normal state of the superconductive phase, strong electron–phonon interactions produce fluctuations of hole-rich and hole-poor regions within the CuO$_2$ sheets that suppress phonons; restoration of in-plane phonons on cooling through $T_c$ signals long-range ordering of the bond-length fluctuations. A dynamic ordering that restores phonons implies the stabilization of a travelling charge-density wave with a possible hybridization of electrons and phonons in the superconductive phase to give heavy vibrons.

7. Although the driving force for the formation of superconductive pairs is not revealed by these indirect thermal conductivity measurements, they do demonstrate that strong electron–lattice interactions cannot be ignored.
Acknowledgments

JQY thanks Professor T Egami for stimulating discussions. The authors thank the NSF and the Robert A Welch Foundation of Houston (TX, USA) for financial support.

References

[1] Cohn M L and Anderson P W 1972 AIP Conf. Proc. (New York 1972) ed D H Douglas p 17
[2] Pines D 1994 Physica C 235 113
    Pines D 1994 Physica B 199/200 300
[3] Goodenough J B 2003 J. Phys.: Condens. Matter 15 R257
[4] Laughlin R B 1998 Phys. Rev. Lett. 80 5188
    Sachdev S 2000 Science 288 475
[5] Castellani C, Di Castro C and Grilli M 1995 Phys. Rev. Lett. 75 4650
[6] Kulic M L 2000 Phys. Rep. 338 1
[7] Yan J-Q, Zhou J-S and Goodenough J B 2003 Phys. Rev. B 68 104520
[8] Yan J-Q, Zhou J-S and Goodenough J B 2004 Phys. Rev. B 69 134409
[9] Takagi H, Cava R J, Marezio M, Batlogg B, Krajewski J J, Peck W F Jr, Bordet P and Cox D E 1992 Phys. Rev. Lett. 68 3777
[10] Kawamata T, Adachi T, Noji T and Koike Y 2000 Phys. Rev. B 62 R11981
[11] Radaelli P G, Hinks D G, Mitchell A W, Hunter B A, Wagner J L, Dabrowski B, Vandervoort K G, Viswanathan H K and Jorgensen J D 1994 Phys. Rev. B 49 4163
[12] Nakamura Y, Uchida S, Kimura T, Motohira N, Kishio K, Kitazawa K, Arima T and Tokural Y 1991 Physica C 185–189 1409
[13] Kaskel D, Stern E A, Hinks D G, Mitchell A W, Jorgensen J D and Budnick J L 1996 Phys. Rev. Lett. 76 439
[14] Goodenough J B, Zhou J-S and Chan J 1993 Phys. Rev. B 47 5275
[15] Zhou J-S, Bersuker G I and Goodenough J B 1995 J. Supercond. 8 541
[16] Zhou J-S, Chen H and Goodenough J B 1994 Phys. Rev. B 50 4168
[17] Yu R C, Salamon M B, Lu J P and Lee W C 1992 Phys. Rev. Lett. 69 1431
[18] Tranquada J M, Ichikawa N and Uchida S 1999 Phys. Rev. B 59 14712