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Effect of doping concentration on normal state resistivity of $La_{2-x}(Ba, Sr)_xCuO_4$ superconductors

DINESH VARSHNEY and R K SINGH*
School of Physics, Devi Ahilya University, Indore 452001, India
*School of Physics, Barkatullah University, Bhopal 462026, India

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Abstract. Based on free electron layered electron gas model of quasi two dimensional CuO_2 layers in La(Ba/Sr)CuO superconductors a model potential $V(q)$ is developed earlier with the electron–electron and electron–phonon interactions. The model approach facilitates the dielectric functions and the dispersion relations of 2D acoustic phonon and plasmon modes. We have then worked out the coupling strength (γ) linking electrons to the 2D acoustic phonon mode ($\hbar\omega_-$) from the residue at the pole of $V(q)$. Furthermore, the scattering time (τ_{e-ph}) during electron–phonon interaction (EPI) for this simplified system is also estimated. The contribution to the normal state in plane resistivity due to EPI is then evaluated. Finally, the variations of τ and ρ is studied with the doping concentration (x) and temperature (T) and the results obtained by us show reasonably good agreement with the available experimental data.

Keywords. Layered electron gas model; coupling strength; electron–phonon interaction; normal state resistivity; scattering time.

1. Introduction

The alkaline-earth (Ba, Sr and Ca)/alkali-metal (Na and K) doped lanthanum cuprates (La_2CuO_4) represents a fascinating class of superconductors. Because of its simple crystal structure and the high transition temperature (T_c), 214 superconductor receive wide attention and still considerable efforts are in progress to look for the pairing mechanism involved as well as the anomalous physical properties observed (Ginsberg 1990; Kresin and Wolf 1991). The transition temperature of alkaline-earth doped La_2CuO_4 is quite high, 28 K for the barium-doped and 36 K for the strontium-doped compound (Moodenbaugh *et al* 1988; Torrance *et al* 1988). The study of normal state transport properties of these SCR's will give the clues to understand the mechanism of pairing and normal state resistivity (ρ) behaviour with increasing doping concentration (x) as well as temperature (T). Since the discovery of high temperature superconductors (HTSCs) (Bednorz *et al* 1987), many theoretical models have been proposed to understand their normal and superconducting state properties. The Eliashberg theory (1960) among them is understood to describe better the superconducting state properties, while the normal state transport properties are well explained from the Bloch–Boltzmann theory (Ziman 1960) based on Migdal approximation, which is valid if the electron–phonon scattering process dominates i.e. on the existence of small parameter $[N(E_F)\hbar/\tau_{e-ph}]$ and $[N(E_F)\hbar\omega_-]$ where $N(E_F)$ is the density of states at the Fermi level, $\hbar\omega_-(q)$ is a typical phonon energy and $1/\tau_{e-ph}$ is the electron scattering rate. The behaviour of normal state resistivity (ρ) as functions of doping concentration (x) and temperature (T) in 214 SCRs has been investigated by several experimental groups (Bharathi *et al* 1989; Torrance *et al* 1989; Koike *et al* 1991).

Recently, Bharathi *et al* (1989) measured ρ as a function of x in pellet form of La–Sr–CuO using positron annihilation spectroscopy and found that $\rho(300\,K)$ goes

through a minimum for $x = 0{\cdot}16$, T_c is maximum for this Sr doping concentration. Subsequently, Torrance *et al* (1989) and Koike *et al* (1991) have reported the variation of resistivity with T in pellet form of $La_{2-x}(Ba, Sr)_xCuO_4 (0{\cdot}0 \leqslant x \leqslant 0{\cdot}30)$. They observed that ρ is generally linear above critical temperature for high values of x, at $T = T_c$, ρ abruptly falls to zero. Furthermore, the fermion scattering time (τ_{e-ph}) due to electron–phonon interaction has been found to decrease with increasing Sr(x) and is $0{\cdot}52 \times 10^{-14}$ sec ($x = 0{\cdot}04$), $1{\cdot}3 \times 10^{-14}$ sec ($x = 0{\cdot}3$) (Suzuki 1989).

In order to understand the mechanism involved and the superconducting state property i.e. T_c, we have discussed in our earlier paper (Varshney and Singh 1995) in which the Coulomb repulsion, electron–phonon and electron–plasmon interactions are incorporated within the BCS frame work. The results deduced from the FTEP approach on T_c for 214 SCRs are quite consistent with the experimental data. Motivated from the earlier achievements and observations on normal state resistivity (ρ), we have attempted to perform an extensive theoretical analysis of doping (x) and temperature (T) dependence of ρ in the alkaline-earth substituted lanthanum cuprates by taking into account the electron–phonon interaction part of FTEP approach. In this paper, we have calculated (a) the coupling strength (γ), (b) the scattering rate (τ) and used these model parameters to deduce the normal state resistivity. We have also estimated the sound velocity (v_s) and the slope of resistivity ($\alpha = \rho_{e-ph}/T$). The analysis based on electron–phonon interaction strongly supports the phonon mechanism.

The plan of this paper is as follows. In § 2 we briefly introduce the FTEP approach as proposed by Varshney and Singh (1995). In the present theoretical investigation, we have separated the electron–phonon and electron–plasmon interactions and used EPI followed by Bloch–Boltzmann theory to interpret the normal state resistivity of layered cuprate superconductors. We have obtained the expressions for coupling strength between electrons via 2D acoustic phonon modes $\gamma_{e-ph}(q)$, and the scattering rate τ_{e-ph}. From these model parameters, we have deduced ρ and ρ_{e-ph} in § 3. The expressions obtained in §§ 2 and 3 are computed numerically in § 4. The results obtained from the model calculations are analysed and discussed in § 5. Finally, in § 6 we end up with some conclusions.

2. Interaction potential

The single crystal of La based cuprate superconductors can be modelled to an infinite array of two-dimensional (2D) conducting planes of charge carriers. In a unit cell of the crystal one conducting CuO_2 plane is considered and the nonconducting planes between CuO_2 planes will form a uniform dielectric host medium. The polarization of dielectric host medium is taken into account via the high frequency dielectric constant ε_∞. The dimension of the unit cell along the c-axis is taken to be d.

The effective interaction potential between the charge carriers can be given as (Varshney and Singh 1995)

$$V(q, q_z, \omega) = \frac{2\pi e^2 S(q, q_z)}{q\varepsilon_\infty \varepsilon(q, q_z, \omega)}, \tag{1}$$

where $\varepsilon(q, q_z, \omega)$ is dielectric response function and is concerned with the polarization function via

$$\varepsilon(q, q_z, \omega) = 1 + P(q, \omega) S(q, q_z), \tag{2}$$

with

$$P(q,\omega) = -\frac{2\pi e^2}{q\varepsilon_\infty}\Pi(q,\omega), \tag{3}$$

and

$$S(q,q_z) = \frac{\sinh(qd)}{\cosh(qd) - \cos(q_z d)}, \tag{4}$$

where q is the in plane wave vector and q_z is in z direction. The effective interaction potential between the charge carriers confined to a single conducting plane is

$$V(q,\omega) = \frac{d}{2\pi}\int_{-\pi/d}^{+\pi/d} V(q,q_z,\omega)\mathrm{d}q_z. \tag{5a}$$

The integral can thus be evaluated by making use of (1) and (5a) as

$$V(q,\omega) = \frac{2\pi e^2 D(q,\omega)\sinh(qd)}{q\varepsilon_\infty |D(q,\omega)|(|D^2(q,\omega) - 1|)^{1/2}}, \tag{5b}$$

with

$$D(q,\omega) = P(q,\omega)\sinh(qd) + \cosh(qd). \tag{6}$$

The interaction potential $V(q,\omega)$ reduces to the effective interaction potential of an isolated 2D conducting plane when $d \to \infty$, i.e. in the long wavelength limit and is

$$V(q,\omega) = \frac{2\pi e^2}{q\varepsilon_\infty} R^{-1}(q,\omega), \tag{7}$$

with

$$R(q,\omega) = 1 + P(q,\omega). \tag{8}$$

The dynamic polarizability in the lanthanum cuprates is due to the active charge carriers as the substitution of trivalent La^{+3} ions by divalent M^{+2}(M = Ba and Sr) ions introduces free holes in the CuO_2 plane which are the active charge carriers. Besides these active charge carriers, ions will also contribute to the mechanism. The polarization function due to charge carriers (c) and ions (i) in the long wavelength limit ($q \to 0$) is expressed as

$$P(q,\omega) = P_c(q,\omega) + P_i(q,\omega), \tag{9}$$

$$= \frac{\Omega_1^2}{(\$^2 - \omega^2)} - \frac{\Omega_2^2}{(\omega^2)}, \tag{10}$$

where Ω_1 and Ω_2 are the 2D electron and ion plasmon frequency and are expressed as

$$\Omega_1^2 = \frac{2\pi n_e e^2 q}{m^* \varepsilon_\infty}, \tag{11}$$

$$\Omega_2^2 = \frac{2\pi n_i e^2 Z^2 q}{M\varepsilon_\infty}, \tag{12}$$

and

$$\$^2 = q^2 v_F^2/2, \tag{13}$$

where n_c, n_i, m^*, M and v_F are the 2D charge carrier density, ionic areal density, effective mass, mass of the unit cell and Fermi velocity of the charge carriers, respectively.

Zeros of the dielectric function $D(q,\omega)$ will yield the frequencies of coupled 2D acoustic plasmon and phonon modes. The resonant frequencies are described by

$$2\omega_{\pm}^2 = [\Omega_1^2\alpha + \Omega_2^2\alpha + \$^2] \pm [(\Omega_1^2\alpha + \Omega_2^2\alpha + \$^2)^2 - 4\Omega_2^2\alpha\$^2]^{1/2}, \tag{14}$$

with $\alpha = qd$.

The dielectric function in the present study provides two coupled modes i.e. $\hbar\omega_+(q)$ and $\hbar\omega_-(q)$. The energy $\hbar\omega_+(q)$, is basically electron plasmon energy and it screens the electron motion along the a–b plane. The modes $\hbar\omega_-(q)$ are the low frequency acoustic phonon modes and have a pronounced effect on the normal state transport properties. Simplification of (14), yields

$$\omega_+^2 = \alpha\Omega_1^2 + \$^2. \tag{15}$$

Here, we made an approximation that $\Omega_1 \gg \Omega_2$,

$$\omega_-^2 = \alpha\Omega_2^2\{1 + \alpha\Omega_1^2/\$^2\}^{-1} \tag{16}$$

where $\Omega_2^2 \ll \$^2$ and $\Omega_2^2/\$^2 \simeq 1$.These modes follow the dispersion relation for 2D acoustic plasmon and phonon modes as

$$\omega_+ = aq^{1/2} + bq, \tag{17}$$

$$\omega_- = cq, \tag{18}$$

where, a, b and c are constants. In (18) c represents the velocity of sound.

In terms of $\hbar\omega_+(q)$ and $\hbar\omega_-(q)$ modes the Fourier transformed effective interaction potential between the charge carriers is (Varshney and Singh 1995)

$$V(q,\omega) = \frac{2\pi e^2}{q\varepsilon_\infty}\left[1 + \frac{(\omega_+^2)(\omega_+^2 - \$^2)}{(\omega^2 - \omega_+^2)(\omega_+^2 - \omega_-^2)} + \frac{(\omega_-^2)(\omega_-^2 - \$^2)}{(\omega^2 - \omega_-^2)(\omega_-^2 - \omega_+^2)}\right]. \tag{19}$$

Here, the first term represents the Coulomb repulsion and the second term denotes the interaction potential from the exchange of 2D acoustic plasmons $[\hbar\omega_+(q)]$. The third term in FTEP is contributed from the exchange of 2D acoustic phonons $[\hbar\omega_-(q)]$. In the following section we will describe the normal state resistivity behaviour by using the Coulomb repulsion and electron–phonon interaction part of the FTEP.

3. Normal state resistivity

The normal state transport properties are well explained from the Bloch–Boltzmann theory where the electron-phonon interaction is dominant. We have followed this approach and deduced the coupling strength of electrons $[\gamma_-(q)]$ at the frequencies $[\omega_-(q)]$. The residues of Fourier transformed effective potential will yield the coupling strength between the electrons linked with 2D acoustic phonon energy. Here, we have considered only electron–electron and electron–phonon term in FTEP to get

$\gamma_-(q)$ as

$$\gamma_-^2(q) = \frac{\pi e^2 \hbar^2 (\omega_-^2)(\omega_-^2 - \$^2)}{q\varepsilon_\infty (\omega_-^2 - \omega_+^2)}, \tag{20}$$

which on simplification becomes

$$\frac{\gamma_-^2(q)}{\omega_-(q)} = \frac{\pi e^2 \hbar^2}{q\varepsilon_\infty} \frac{(D' - \$^2)}{(D'' - \$^2)}, \tag{21}$$

with

$$D' = (\alpha\Omega_2^2 - \alpha\Omega_1^2), \tag{22}$$

and

$$D'' = (\alpha\Omega_2^2 - 2\alpha\Omega_1^2). \tag{23}$$

The temperature dependent part of the normal state resistivity is included in the scattering time due to the electron–phonon interaction. Therefore, the mean time to absorb or emit a phonon of energy $\hbar\omega_-(q)$ is expressed as

$$\frac{1}{\tau_{e-ph}} = \frac{2\pi}{\hbar} \int_0^{2k_F} \frac{d^2q}{(2\pi)^2} |\gamma_-(q)|^2 (1 - \cos\theta)$$
$$[n_q \delta(\varepsilon_{k+q} - \varepsilon_k - \hbar\omega_-) + (n_{q+1})\delta(\varepsilon_{k+q} - \varepsilon_k + \hbar\omega_-)], \tag{24}$$

where n_q is the phonon occupancy factor, θ is the angle between k and $k+q$. In the high temperature limit n_q is reduced to $K_B T/\hbar\omega_-$. As the typical Debye temperature in these cuprates is of the order of 400 K, for such temperatures we have used the exact form of (24) and substituting the value of $\gamma_-(q)$ from (21) and integrating over q yields

$$\frac{1}{\tau_{e-ph}} = \frac{m^* K_B T}{\pi\hbar^4 k_F^3} \int_0^{2k_F} \frac{|\gamma_-(q)|^2}{\omega_-} \frac{q^2 dq}{\{1 - (q/2k_F)^2\}^{1/2}}, \tag{25}$$

$$= \frac{m^* K_B T}{\pi\hbar^4 k_F^3} \frac{\pi e^2 \hbar}{\varepsilon_\infty} \int_0^{2k_F} \frac{(D' - \$^2)}{(D'' - \$^2)} \frac{q^2 dq}{\{1 - (q/2k_F)^2\}^{1/2}},$$

$$= \frac{m^* K_B T}{\hbar^3 k_F} \frac{2\pi e^2}{\varepsilon_\infty} \int_0^{2k_F} \frac{(V' - \#^2)}{(V'' - \#^2)}, \tag{26}$$

with

$$V' = (\Omega_2^2 - \Omega_1^2)\frac{d}{q}, \tag{27}$$

$$V'' = (\Omega_2^2 - 2\Omega_1^2)\frac{d}{q} \tag{28}$$

$$\#^2 = V_F^2. \tag{29}$$

The normal state resistivity due to EPI in terms of τ_{e-ph} is expressed as

$$\rho_{e-ph} = \frac{m^*}{n_c L e^2} \frac{1}{\tau_{e-ph}}, \tag{30}$$

where L is the number of conducting layers per unit length along 'c' axis. The above describes the temperature dependent part of the normal state resistivity due to the scattering of electrons by 2D acoustic phonons. From (26) and (30) we have

$$\rho_{e-ph} = \frac{m^{*2} K_B T}{\hbar^3 k_F n_c} \frac{2\pi}{L\varepsilon_\infty} \frac{(V' - \#^2)}{(V'' - \#^2)}. \tag{31}$$

The slope of resistivity is defined as

$$\alpha = \frac{\rho_{e-ph}}{T}. \tag{32}$$

Doping of alkaline earths (Ba, Sr and Ca)/alkali-metals (Na and K) in the La_2CuO_4 compound will develop the defects. Hence, besides the scattering of electrons via phonons, scattering via defects also exists. The scattering of electrons due to impurities are temperature independent. The temperature dependent part of resistivity is indebted to electron–phonon scattering process. Hence, one could think of two-component resistivity and

$$\rho_{total} = \rho_0 + \rho_{e-ph}(T), \tag{33}$$

with ρ_0 as the residual resistivity developed due to impurities scattering.

Using the developed expressions following the Bloch–Boltzmann theory we have obtained the normal state resistivity as a function of composition (x) and temperature (T) in $La_{2-x}(Ba, Sr)_xCuO_4$ superconductors. The analysis is presented and discussed below.

4. Method of computation

In order to compute normal state resistivity (ρ) as a function of composition (x) and temperature (T), in 214 SCRs we have used the values of the effective mass m^* equal to 6 m_e and 5 m_e for $La_{2-x}Ba_xCuO_4$ and $La_{2-x}Sr_xCuO_4$, respectively. These are consistent with the specific heat measurements on the pellets (Batlogg *et al* 1987). The high frequency background dielectric constant ε_∞ is 4·5 (Bozovic 1990) from the electron energy loss spectroscopy for cuprate superconductors. The number of conducting layers per unit length along 'c' axis is $L = 7{\cdot}58 \times 10^{16}\,cm^{-1}$ (Zhao *et al* 1989). The charge distribution in the copper oxide plane contain the ionic charge, the free electron charge and the electronic charge that is bound to the ions. The sum of ionic and bound electronic charge lying on the conducting plane is denoted by Ze. The value of ionic charge, Z is based on the valences for ions in CuO_2 plane and is $Z = -2$. M denotes the reduced mass of Cu and O in CuO_2 plane and is 12·77 amu indicating that the vibration considered is the Cu–O vibration.

The ionic areal density n_i is the reciprocal of the unit cell projection perpendicular to the 'c' axis and is $n_i = 1/ab$, where a and b denote the lattice parameters. The 2D electronic density $n_c = n_0 \times 10^{14}\,cm^{-2}$ is obtained from the Hall effect measurements for different values of x on pellets of Sr doped lanthanum cuprate (Ong *et al* 1987). Using these input parameters we have computed the normal state resistivity (ρ) as a function of x and T for $La_{2-x}(Ba, Sr)_xCuO_4$ superconductors and are presented along with discussions in following section.

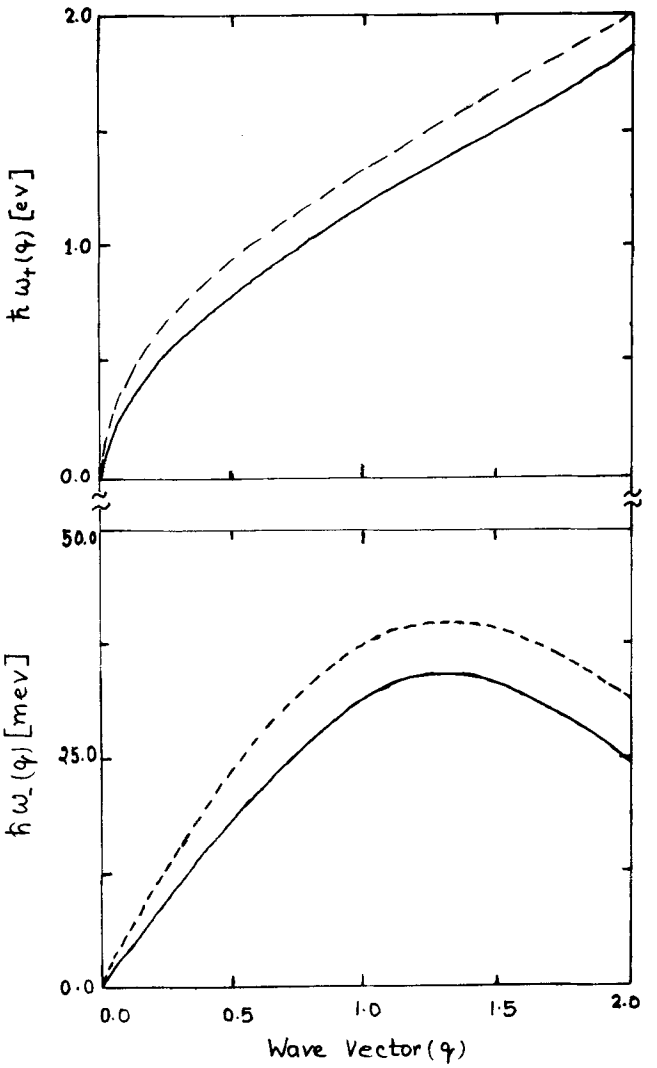

Figure 1. Variation of 2D acoustic phonon (plasmon) energy $\hbar\omega_-(\hbar\omega_+)$ as a function of wave vector (q) [$Å^{-1}$] shown by (———) (– – –) for Ba(Sr) dopings, respectively.

5. Results and discussion

We have plotted the 2D acoustic phonon $\hbar\omega_-$ and plasmon $\hbar\omega_+$ energy from (15) and (16) developed in § 2 as function of wave vector ($q = 2K_F \sin\theta$) in $Å^{-1}$ along the a–b plane for Ba and Sr dopings, respectively in figure 1. It is seen from figure 1 that $\hbar\omega_+(q)$ displays a characteristic square root behaviour for small wave vectors and becomes linear for higher wave vectors. This is in accordance with the 2D acoustic plasmon dispersion relation. The values of $\hbar\omega_+(q)$ are consistent with the experimental data (Bozovic 1990). Furthermore, the plot of $\hbar\omega_-(q)$ as a function of wave vector (q) is characterized by the low acoustic phonon frequencies. The long wavelength sound velocity is expressed as $v_s = (K_B a\omega_-/\hbar\pi)$ with 'a' as the interatomic distance. Taking

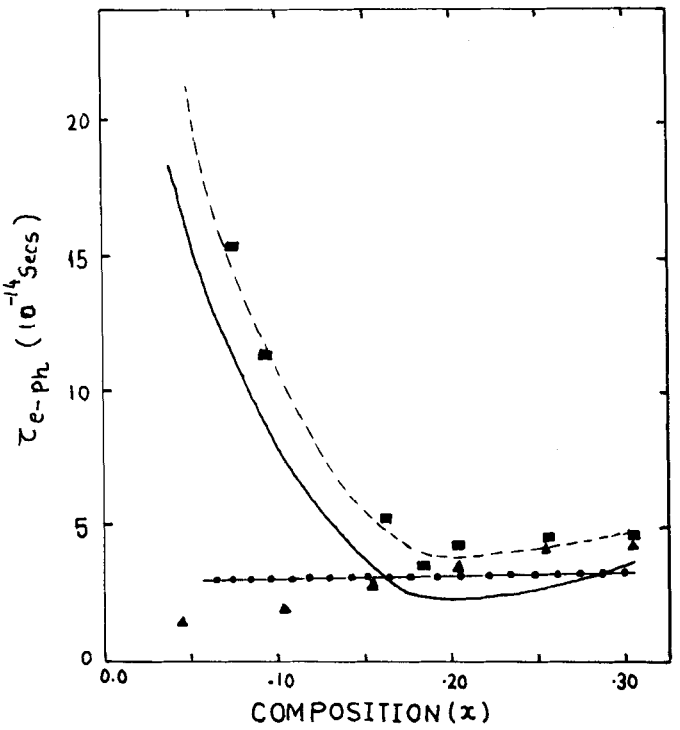

Figure 2. Variation of electron–phonon scattering time (τ_{e-ph}) as a function of composition (x) for Ba(———) and Sr(– – – –) doped superconductors, respectively along with experimental data (▲) Suzuki (1989); (●–●–●–●–●) Orenstein *et al* (1990); and theoretical prediction (■) Kim *et al* (1991).

$a = 3{\cdot}77$ Å (Cava *et al* 1987) for $La_{1{\cdot}85}Sr_{0{\cdot}15}CuO_4$, we have estimated the sound velocity as 4×10^5 cm sec^{-1}. This is in good agreement with the data reported by Bishop *et al* (1987).

The scattering rate between electrons and phonons is inversely proportional to the temperature (T) as is evident from (26). In figure 2 the computed values of τ_{e-ph} with composition (x) have been plotted along with the experimental data (Suzuki 1989; Orenstein *et al* 1990) and theoretical predictions by Kim *et al* (1991) at $T = 100$ K. Earlier Kim *et al* (1991) evaluated the electron–phonon contribution to the transport life time by solving the Boltzmann transport equation variationally and the Eliashberg function $\alpha^2(\omega)$ F(ω). The effects of strong Coulomb correlations on the electron–phonon interactions are calculated using a Frozen-phonon approach and the electron–phonon coupling constant is deduced to see the effects of this on normal state resistivity. The figure inferred that electron phonon scattering time τ_{e-ph} is high at low concentrations and is low at the higher dopings. We notice that τ_{e-ph} is rather insensitive to the carrier concentration in the metallic regime ($x > 0{\cdot}28$) and is approximately 5×10^{-14} seconds. This is consistent with the experimental data (Suzuki 1989; Orenstein *et al* 1990). For low dopings our results do not agree with the experimental data because of the fact that besides electron–phonon scattering other scattering mechanisms are also present. The presence of multiple scattering mechanism lowers the scattering time and concentration independent. We found that scattering time due to

the EPI is rather sensitive to low dopings, i.e. in the insulating phase but is insensitive for higher dopings. Here, in the present context we have estimated τ_{e-ph} by taking proper layered structure and the Frohlich coupling strength linking between neighbouring electrons is evaluated from the residues of the model potential. We find that the results on τ_{e-ph} follow the trend as earlier predicted by Kim *et al* (1991). This is attributed to the strong electron–electron effects on the electron–phonon interactions in our simple model calculations. However, our results on Ba doping follow the same trends as of Sr doped and are of only academic interests, at present, as they could not be compared due to the lack of experimental data on them.

The normal state resistivity (ρ) due to EPI at room temperature have been plotted against the Ba and Sr compositions (x) in figure 3 along with the experimental data (Bharathi *et al* 1989) on pellets. The ρ_{e-ph} includes extrapolated values of residual resistivity ρ_0. It is apparent from the curves that at low composition (x), ρ_{e-ph} is higher and at higher doping concentrations (x), it remains lower. We have noticed that ρ_{e-ph} diverges as the insulator region ($x < 0{\cdot}08$) is approached as the weak coupling in between electrons and phonons exists. Furthermore, the low values of ρ_{e-ph} at higher compositions signify the metallic behaviour. The variations of normal state resistivity ρ_{e-ph} due to EPI together with the extrapolated ρ_0 with the Sr compositions (x) are

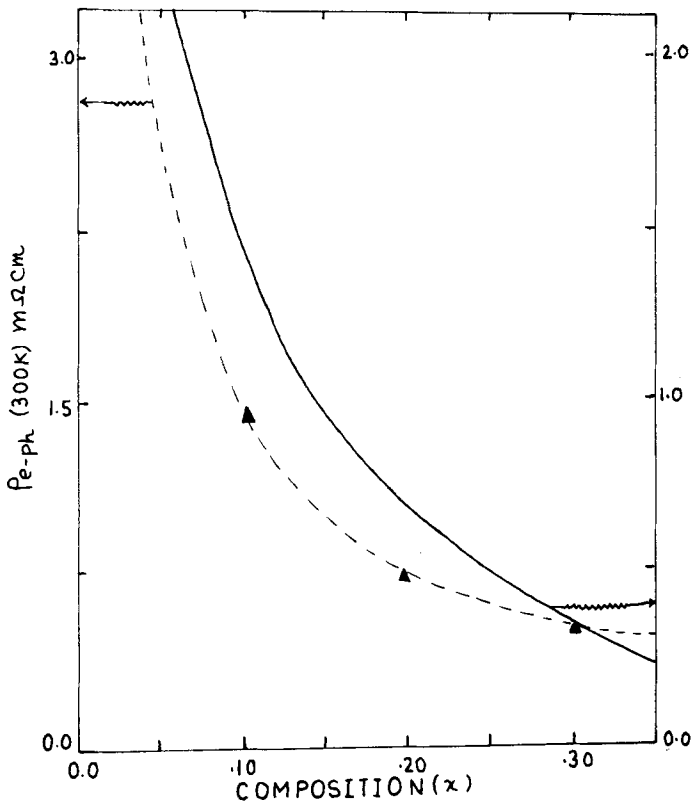

Figure 3. Variations of normal state resistivity (ρ_{e-ph}) with Ba(———) and Sr(– – –) compositions, respectively. The triangles (▲) are the experimental data taken from Bharathi *et al* (1989).

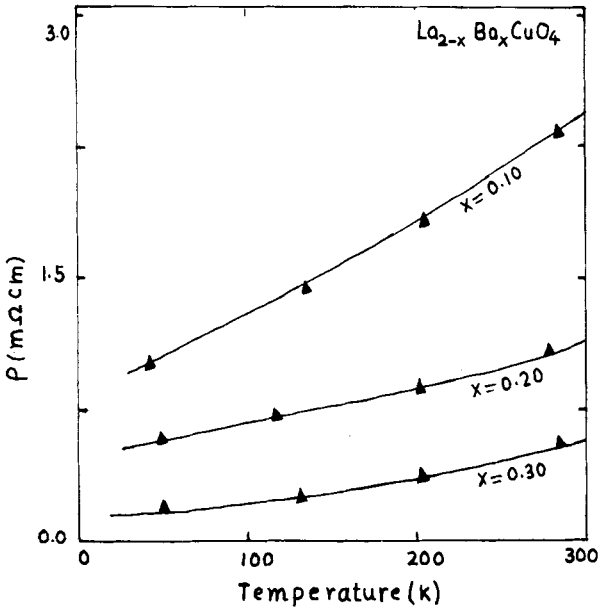

Figure 4. Variation of normal state resistivity (ρ) with temperature (T) for Ba dopings (x). The triangles (▲) are the experimental data taken from Koike *et al* (1991).

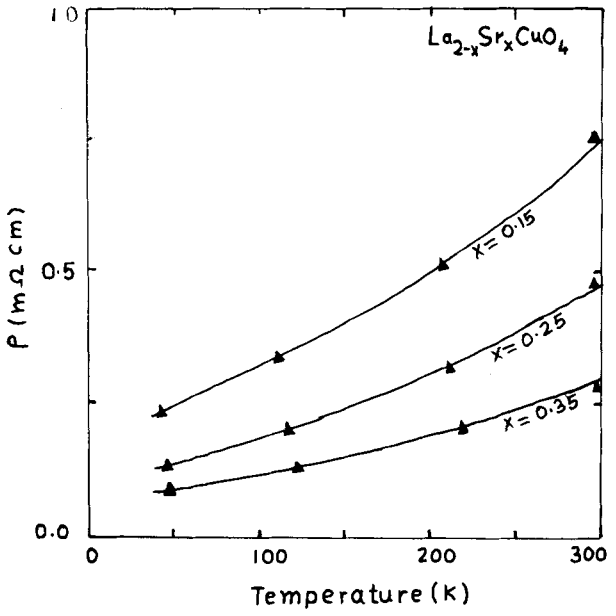

Figure 5. Variation of normal state resistivity (ρ) with temperature (T) for Sr dopings (x). The triangles (▲) are the experimental data taken from Suzuki (1989).

consistent with the experimental data (Bharathi *et al* 1989) on pellets. We have used the data for charge carrier density as depicted from the Hall effect measurements on pellets of $La_{2-x}Sr_xCuO_4$ superconductors (Ong *et al* 1987).

Figures 4 and 5 depict the behaviour of normal state resistivity as a function of temperature (T) for Ba and Sr doping respectively, using (24) and (30) and compared them with the available experimental data (Suzuki 1989; Koike *et al* 1991). It is inferred from the figures that temperature dependence of ρ is almost linear at high temperature. This is in good agreement with the experimental data. Furthermore, the experimental observations show that ρ extrapolates to be zero, but our results lead to a finite intercept. This intercept is interpreted in terms of residual resistivity (ρ_0) that arises from the scattering due to impurities and defects developed by the doping in pure La_2CuO_4 compound. We estimate $\rho_0 = 0{\cdot}6\,(\cdot 15)$ mΩcm at $x = 0{\cdot}20$ for Ba(Sr) dopings, respectively. The roughly linear part of NS ρ is due to the scattering of phonons and is the temperature dependent contribution. We obtained the slope of resistivity $\alpha = 4{\cdot}5\,\mu\Omega$cm/K at 300 K. This is comparable with reported experimental data of $\alpha = 4{\cdot}0\,\mu\Omega$cm/K at room temperature for Sr doped superconductors (Torrance *et al* 1989).

6. Conclusions

We have succeeded in predicting the observed behaviour of normal state resistivity in Ba and Sr doped SCRs from FTEP approach which properly takes into account the electron–electron and electron–phonon interactions. In the present studies the observed behaviour of ρ with x and T are comparable with the experimental data. In fact, we have calculated the model parameters ($\hbar\omega_-$, $\hbar\omega_+$, γ_{e-ph}, τ_{e-ph} and ρ_{e-ph}) by means of a strong coupling theory in which the electron–phonon interaction dominates and proper care is taken of the layered structure to get better results on normal state resistivity which are in agreement with the experimental data. This strongly supports the mechanism proposed earlier by Varshney and Singh (1995). The estimated values of v_s and α are consistent with the experimental observations. The deviations in agreement of τ_{e-ph} strongly supports the presence of some other scattering mechanism, besides electron–phonon. From our point of view, we do think that at the low dopings the electron–electron scattering is rather important.

In conclusion the normal state resistivity (ρ) of $La_{2-x}(Ba, Sr)_x CuO_4$ may be explained by means of mechanism involved in the strong coupling theory of electron–phonon interactions.

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