ARPES Spectra of Bi2212 give the Coulomb Coupling $\lambda^C \approx 1$ and the Electron-Phonon Coupling $\lambda^{EP} = 2 - 3$

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We show that the double kink-structure in the electronic self-energy of Bi2212 near the nodal point at low energy $\omega_1 \approx 50 - 70 \text{ meV}$ and at high energy at $\omega_2 \approx 350 \text{ meV}$, observed recently in the ARPES measurements by Valla et al.¹, gives that the electron-phonon (EPI) coupling constant $\lambda^{EP}$ in the normal part of the self-energy $\Sigma(\omega)$ is twice larger than the Coulomb coupling $\lambda^C$. The experimental data for $\text{Re} \Sigma(\omega)$ up to energies $\sim 350 \text{ meV}$ can be satisfactorily explained by $\lambda^{EP} \approx 2.1$ and $\lambda^C \approx 1.1$. Additionally the low energy slope of the ARPES $\text{Re} \Sigma(\omega)$ at $\omega < 20 \text{ meV}$ gives a hint that the low energy phonons might contribute significantly to the EPI coupling, i.e. $\lambda^{low,EP} > 1$, thus giving the total EPI coupling constant $\lambda_{tot} = \lambda^{EP} + \lambda^{low,EP} > 3$. In order to test the role of low frequency phonons by ARPES measurements a much better momentum resolution is needed than that reported in [1]. Possible pairing scenarios based on ARPES, tunnelling and magnetic neutron scattering measurements are discussed.

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In the last several years ARPES measurements in high temperature superconductors (HTSC), with much better momentum and energy resolution, made a breakthrough in determining the energy and momentum dependence of the quasiparticle self-energy in HTSC. Important measurements by the Shen group [2], [3] gave evidence on the low-energy kink in the quasiparticle spectrum around the phonon energy between 40-70 meV, in both nodal and antinodal points. Since the low-energy kink is pronounced practically in all HTSC materials, above and below $T_c$, it is most probable of the phononic origin, as it was recently confirmed by the ARPES isotope effect in $\text{Re} \Sigma(\omega)$ [4]. These results gave a new impetus in the HTSC physics by renewing the importance of the electron-phonon interaction (EPI) in the quasiparticle scattering and in the pairing mechanism of HTSC. In spite of the very convincing evidence in ARPES, tunnelling, STM and optic measurements that EPI in HTSC is rather strong [5], [6], [7], the proponents of the spin-fluctuation (SF) mechanism assume a phenomenological form for the spectral function $\alpha^2 F_1^SF(\omega,\mathbf{k}) \sim g^2 SF \text{Im} \chi(\mathbf{k},\omega)$, where $\chi(\mathbf{k},\omega)$ is the spin susceptibility. In systems where $\chi(\mathbf{k},\omega)$ is strongly peaked around some $\mathbf{k}$-points, like in under-doped HTSC where $\mathbf{k} = Q_{AF}(\pi/a,\pi/a)$, this approximation for $\alpha^2 F_1^SF(\omega,\mathbf{k})$ is unwarranted, since the vertex correction (terms beyond the Migdal-Eliashberg approximation) can significantly influence the self-energy which is most probably suppressed [18], [19].

It is worth of mentioning, that in order to fit the quasiparticle self-energy by the SF theory it is necessary to make a radical assumption that the SF coupling constant is (unrealistically) large, i.e. $g_{SF} \sim 1 \text{ eV}$ [5]. Such a large value of $g_{SF}$ is difficult to justify, both theoretically and experimentally. In Ref. [16] it was shown that the SF phenomenology fails to give large $T_c$ even for $g_{SF} \gg 1 \text{ eV}$, if the spectral function $\text{Im} \chi(\mathbf{k},\omega)$ is taken from the neutron scattering data and not from the low frequency NMR spectra as it was done in [5]. The situation is even worse if one fits the slope $d\rho(T)/dT$ of the resistivity $\rho(T)$, which gives $g_{SF} \sim 0.3 \text{ eV}$ and $T_c < 7 K$ - see discussions in [5], [6], [7] and references therein. The most impressive evidence for ineffectiveness of SF to explain (solely) the self-energy effects came from the magnetic neutron scattering measurements by the Bourges’s group [11]. They show, that by changing doping from slightly underdoped to optimally doped systems there is a dramatic change in the magnitude and $\omega$-shape of the magnetic spectral function $\text{Im} \chi(\mathbf{k} \approx Q,\omega)$ in the normal state at $T > T_c$. These experiments show [11] that $AT > T_c$ in the slightly underdoped $YBa_2Cu_3O_{6.92}$ (with $T_c = 91 K$)
Im $\chi(Q, \omega)$ is peaked (with very large value) at $\omega \approx 35$ meV, while in the optimally doped YBa$_2$Cu$_3$O$_{6.97}$ (with $T_c = 92.5$ K) Im $\chi(Q, \omega)$ is drastically suppressed and practically negligible. This result was confirmed quite recently in Ref. [12] in the optimally doped systems where Im $\chi(k, \omega)$ is practically negligible at all $k$. Moreover, in spite of the large suppression of Im $\chi(k, \omega)$ in the optimally doped HTSC [11] there is a very small difference in $T_c$ (91 K vs 92.5 K) of slightly underdoped and optimally doped HTSC. This means that the large reconstruction and suppression of Im $\chi(k, \omega)$ by doping does not affect superconductivity at all. This result means also that $g_{SF} \ll 1$ eV in the SF phenomenology and that the latter should be abandoned as the leading pairing mechanism in HTSC. It is also worth of mentioning that the recent numerical calculations on Hubbard and t-J model render valuable evidence that there is no high temperature superconductivity (SC) in these models [13]. If SC exists at all in these models its $T_c$ must be rather low. Therefore, other interactions, such as EPI and the direct Coulomb interaction, should be taken into account.

In Refs. [14, 15] it was argued that the low-energy (below 1 eV) behavior of the ARPES $\Sigma(k, \omega)$, obtained recently in [2, 3, 4], can be qualitatively and semi-quantitatively explained by the combined effect of the EPI and Coulomb interaction. The recent ARPES results [2, 3, 4] for the effective real part of the self-energy at energies $\omega < \omega_{ph}^{\text{max}} \approx 80$ meV ($\omega_{ph}^{\text{max}}$ is the maximal phonon frequency) give evidence that the effective EPI coupling [3] is rather strong $\sim 1$. However, in [2, 3, 4] the EPI self-energy was obtained by subtracting the high energy slope of the quasiparticle spectrum $\omega(\xi_k)$ at $\omega \sim 0.3$ eV. The latter is due to the Coulomb interaction. Although the position of the low-energy kink is not affected by this procedure (if $\omega_{ph} \ll \omega_{C}$), the above (subtraction) procedure gives in fact an effective EPI self-energy $\Sigma_{\text{eff}}^{EP}(k, \omega)$ and coupling constant $\lambda_{z,\text{eff}}^{EP}(k)$ only. Let us briefly demonstrate that $\lambda_{z}^{EP}(k)$ is smaller than the real EPI coupling constant $\lambda_{z}^{EP}(k)$. (The total self-energy is $\Sigma(k, \omega) = \Sigma^{EP}(k, \omega) + \Sigma^{C}(k, \omega)$ where $\Sigma^{C}$ is the contribution due to the Coulomb interaction. At very low energies $\omega \ll \omega_{C}$ one has usually $\Sigma^{C}(k, \omega) = -\lambda_{z}^{C}(k)$, where $\omega_{C}(\sim 1$ eV) is the characteristic Coulomb energies and $\lambda_{z}^{C}$ the Coulomb coupling constant. The quasiparticle spectrum $\omega(k)$ is determined from the condition

$$\omega - \xi(k) - \text{Re}[\Sigma^{EP}(k, \omega) + \Sigma^{C}(k, \omega)] = 0,$$

where $\xi(k)$ is the bare band structure energy. At low energies $\omega \ll \omega_{ph}^{\text{max}} \ll \omega_{C}$ Eq. (2) can be rewritten in the form

$$\omega - \xi^{\text{ren}}(k) - \text{Re} \Sigma_{\text{eff}}^{EP}(k, \omega) = 0,$$

where $\xi^{\text{ren}}(k) = [1 + \lambda_{z}^{C}(k)]^{-1}\xi(k)$ and $\text{Re} \Sigma_{\text{eff}}^{EP}(k, \omega) = [1 + \lambda_{z}^{C}(k)]^{-1}\text{Re} \Sigma^{EP}(k, \omega)$. Since at very low energies $\omega \ll \omega_{ph}^{\text{max}}$ one has $\text{Re} \Sigma^{EP}(k, \omega) = -\lambda_{z}^{EP}(k)\omega$ and $\xi^{\text{ren}}(k) - \lambda_{z}^{C}(k)$ the real coupling constant is related to the effective one by

$$\lambda_{z}^{EP}(k) = [1+\lambda_{z}^{C}(k)]\lambda_{z,\text{eff}}^{EP}(k) > \lambda_{z,\text{eff}}^{EP}(k).$$

At higher energies $\omega_{ph}^{\text{max}} < \omega < \omega_{C}$, which are less important for pairing and where the EPI effects are suppressed, one has $\text{Re} \Sigma(k, \omega) \approx -\lambda_{z}^{C}(k)\omega$. The conclusion is that in order to obtain the correct values for the EPI and Coulomb coupling constant the self-energy $\text{Re} \Sigma(k, \omega)$ should be measured in a broad energy interval. This was done in the recent ARPES measurements on Bi2212 and LSCO by Valla et al. [1] where $\text{Re} \Sigma(k, \omega)$ was experimentally determined in the broad energy interval. The measured $\text{Re} \Sigma^{\text{exp}}(k, \omega)$ at $T = 10$ K near and slightly away from the nodal point in the optimally doped Bi2212 with $T_c = 91$ K [1] is shown in Fig. 1. It is seen in Fig. 1 that $\text{Re} \Sigma^{\text{exp}}(k, \omega)$ has two kinks—the first one at low energy $\omega_{1} \approx \omega_{phil} \approx 50 - 70$ meV which is most probably of the phononic origin [2, 3, 4], while the second kink at higher energy $\omega_{2} \approx \omega_{C} \approx 350$ meV is evidently due to the Coulomb interaction (by including spin fluctuations too). However, one of the most important results in Ref. [1] is that the slopes of $\text{Re} \Sigma^{\text{exp}}(k, \omega)$ at low $(\omega < \omega_{phil}^{\text{high}})$ and high energies $(\omega_{phil}^{\text{high}} < \omega < \omega_{C})$ are very different. The low-energy and high-energy slope near the nodal point are depicted and shown in Fig. 1 semantically (thick lines). From Fig. 1 it is obvious that EPI prevails at low energies $\omega < \omega_{phil}^{\text{high}}$. More precisely digitalization of $\text{Re} \Sigma^{\text{exp}}(k, \omega)$ in the interval $\omega_{phil}^{\text{high}} < \omega < 0.4$ eV gives the Coulomb interaction.
coincidence \( \lambda_z^C \)

\[(\lambda_z^{SF} < \lambda_z^C) \approx 1.1 \quad (4)\]

while the same procedure at 20 meV \( \approx \omega_{ph}^{low} < \omega < \omega_{ph}^{high} \approx 50 \text{–} 70 \text{meV} \) gives the total coupling constant \( \lambda_z^{EP} \approx 3.2 \) and the EPI coupling constant \( \lambda_z^{EP} \approx \lambda_z^{high,EP} \).

\[\lambda_z^{EP} \approx 2.1. \quad (5)\]

This estimation tells us that at and near the nodal point the EPI interaction dominates in the quasiparticle scattering at low energies since \( \lambda_z^{EP} \approx 2 \lambda_z^{C} > 2 \lambda_z^{SF} \), while at large energies (compared to \( \omega_C \)) the Coulomb interaction with \( \lambda_z^{C} \approx 1.1 \) dominates. We point out that EPI near the anti-nodal point can be even larger than in the nodal point, mostly due to the higher density of states at the anti-nodal point.

It is well known that the low energy quasiparticle scattering dominates the transport and thermodynamic properties of metallic systems such as HTSC. Comparing the Valla et al. results \( \text{[1]} \) with the previous ARPES measurements \( \text{[2, 3]} \) it is apparent that the real EPI coupling constant \( \lambda_z^{EP}(k) \) - obtained from experiments in Ref. \( \text{[1]} \), are at least twice larger than the effective one \( \lambda_z^{eff}(k) \) - from Refs. \( \text{[2, 3, 4]} \), i.e. \( \lambda_z^{EP}(k) \approx 2 \lambda_z^{EP} \).

Moreover, the results for \( \Re \Sigma^{EP}(k, \omega)(= -\lambda_z \omega) \) at very low energies \( \omega < \omega_{ph}^{low} \approx 20 \text{meV} \), shown in Fig. \( \text{[4]} \) give a hint to an even larger slope which gives the low-energy coupling \( \lambda_1 = \lambda_z^{low,EP} + \lambda_z^{EP} + \lambda_z^C \). This slope is larger than those in the (higher energy) interval 20 meV \( < \omega < 80 \text{meV} \) \( \text{[1]} \) which means that the EPI coupling due to low energy phonons may be rather strong, i.e. \( \lambda_z^{low,EP} > 0.3 \text{–} 1.3 \), while the total EPI coupling is \( \lambda_z^{EP} |_{z,tot} = \lambda_z^{low,EP} + \lambda_z^{EP} > 2.4 \text{–} 3.4 \) - we call this the L-scenario.

It is worth of pointing out that if the L-scenario turns out to be correct and if the high value of \( \lambda_z^{low,EP} > 1.3 \) is realized, then the vibrations of heavier atoms (than oxygen) may contribute significantly to pairing in HTSC.

One of the consequence of this result would be a reduction (from the canonical value 0.5) of the oxygen isotope effect in optimally doped HTSC materials, as it was observed experimentally - see \( \text{[8]} \) and references therein. Furthermore, the value \( \lambda_z^{low,EP} > 1.3 \) is compatible with the earlier tunnelling measurements \( Bi_2Sr_2CaCu_2O_8 \) (with \( T_c \approx 70 \text{K} \) \( \text{[10]} \) which give the total EPI coupling constant \( \lambda_{\text{tunn}}^{EP} \approx 3.5 \), while \( \lambda_{\text{tunn}}^{low,EP} \approx 2.1 \) for \( \omega = 20 \text{meV} \), as well as with the EPI experiments from Ref. \( \text{[22]} \) were analyzed. However, to elucidate the role of low energy phonons by ARPES measurements a much better momentum resolution (\( \delta k \)) is needed than that reported in \( \text{[3]} \), where \( \delta k \approx 0.04 \text{Å}^{-1} \) and the self-energy resolution is \( \delta \Re \Sigma > 20 \text{meV} \). This resolution is insufficient for a definitive conclusion on the contribution of low energy phonons.

Finally, although it is still premature for giving a definitive pairing scenario in HTSC, the experimental evidence for d-wave pairing and for the strong EPI in HTSC imply a necessary condition for EPI in HTSC: EPI in HTSC must be peaked at small transfer momenta, i.e. there is forward scattering peak in EPI. Otherwise, if the rather large EPI would be weakly momentum independent it would inevitably destroy d-wave pairing in HTSC - see more in \( \text{[4, 5, 6]} \). In that respect several pairing scenarios are imaginable depending on the strength of the EPI coupling in the d-wave channel \( \lambda_{d,\text{tot}}^{EP} \). Since the SF coupling is small and \( \lambda_z^{SF} \ll \lambda_z^{C} \approx 1 \), then if \( \lambda_{d,\text{tot}}^{EP} \) is sufficiently strong \( \lambda_{d,\text{tot}}^{EP} > 1 \) than d-wave pairing is dominated by EPI which in conjunction with the Coulomb interaction (by including both the short- and long-range parts) may give high \( T_c \). In that case the s-wave part \( \lambda_z^{C} \) of the Coulomb interaction suppresses s-wave pairing while the d-wave part \( \lambda_{d}^{EP} (\ll \lambda_{d}^{C} \approx 1 \) 0 affects d-wave pairing weakly. If the Coulomb interaction is attractive in the d-channel \( \text{[17]} \), i.e. \( \lambda_{d}^{C} (\approx 0) \), it strengthens d-wave pairing additionally. Very interesting situation arises if both EPI and the (total) Coulomb interaction give appreciable attraction in the d-channel with \( \lambda_{d}^{EP} \approx | \lambda_{d}^{C} | \). The reasons for the large EPI coupling constant in the d-wave channel have been discussed in \( \text{[6, 7]} \), where it was argued that strong correlations in conjunction with the weakly screened EPI Madelung coupling in the ionic-metallic structure of HTSC can produce the forward scattering peak in the EPI coupling and other charge scattering processes \( \text{[20, 6]} \). In such a case the EPI coupling in the d-wave channel \( \lambda_{d}^{EP} \) may be appreciable, while on the other hand the transport coupling due to EPI \( (\rho(T) \sim \lambda_{d} T) \) is suppressed, i.e. \( \lambda_{d}^{EP} < \lambda_{d}^{EP}/3 \text{[20, 6]} \). The latter result resolves the long-standing experimental puzzle in HTSC that the experimental value of \( \lambda_{d} \) is too small to give high \( T_c \). Contrary to low temperature superconductors, where in most materials \( \lambda_{d} \approx \lambda_{d} \approx \lambda_{d}^{EP} \), in HTSC materials one has \( \lambda_{d} < \lambda_{d} \approx \lambda_{d}^{EP} \approx \lambda_{d}^{C} \), then SF and EPI would interfere constructively giving high \( T_c \). However, as we discussed above this scenario contradicts the magnetic neutron scattering measurements of the Bourges’s group \( \text{[11, 12]} \), which imply that \( \lambda_{d}^{EP} \ll \lambda_{d}^{C} \). Therefore, the SF approach should be abandoned.

In conclusion, we have argued that the recent ARPES measurements of the self-energy in Bi2212 at and near the nodal point by Valla et al. \( \text{[1]} \) give evidence that at energies \( \omega < 70 \text{meV} \) the EPI interaction is at least twice stronger than the Coulomb interaction (which includes spin fluctuations too). It turns out that \( \lambda_z^{EP} > 2.1 \) and \( \lambda_z^{SF} \ll \lambda_z^{C} \approx 1.1 \). These ARPES measurements give also a hint that the low-energy phonons can give an appreciable EPI coupling \( \lambda_z^{EP} > 1 \) and \( \lambda_z^{EP} > 3 \). In order to clear up the role of low frequency phonons in EPI ARPES measurements need much better resolution
in momentum space than that reported in [1]. If confirmed by other experiments the ARPES results by Valla et al. [1] will render undisputable evidence for the leading role of the electron-phonon interaction in the scattering and pairing mechanism of HTSC materials.

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