The ARPES spectra in high-$T_c$ superconductors (HTSC) show four distinctive features in the quasiparticle self-energy $\Sigma(k,\omega)$. All of them can be explained consistently by the theory in which the electron-phonon interaction (EPI) with the forward scattering peak (FSP) dominates over the Coulomb scattering. In particular, this theory explains why there is no shift of the nodal kink at 70 meV in the superconducting state, contrary to the clear shift of the anti-nodal singularity at 40 meV. The theory predicts a “knee”-like structure of $|\ Im \Sigma(\omega)| = |\ Im \Sigma_{\text{ph}}(\omega) + \ Im \Sigma^G(\omega) |$, which is phonon dominated. $|\ Im \Sigma(\omega)| \approx |\ Im \Sigma_{\text{ph}}(\omega_{\text{ph}})| \sim \pi \lambda_{\text{ph}} \omega_{\text{ph}}/2$, for $\omega \approx \omega_{\text{ph}}^{(70)}$, and shows linear behavior $|\ Im \Sigma(\omega)| \approx |\ Im \Sigma_{\text{ph}}(\omega_{\text{ph}})| + \pi \lambda_{C,\omega}/2$, for $\omega > \omega_{\text{ph}}^{(70)}$ due to the Coulomb scattering.

ARPES spectra give $\lambda_{\text{ph}} > 1$ - which is obtained from $Re \Sigma$, and $\lambda_C < 0.4$ - obtained from $Im \Sigma$, i.e. $\lambda_{\text{ph}} \gg \lambda_C$. The dip-hump structure in the spectral function $A(k_F,\omega)$ comes out naturally from the proposed theory.

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Introduction - The pairing mechanism in high-temperature superconductors (HTSC) is under intensive debate [1, 2]. In that respect ARPES experiments play a central role for theory, since they give information on the quasiparticle spectrum, life-time effects and indirectly the pairing potential. Recent ARPES experiments on various HTSC families, such as $La_{2-x}Sr_xCuO_4$ and BISCO [3, 4, 5, 6], show four distinctive features in the quasiparticle self-energy $\Sigma(k,\omega)$: (I) There is a kink in the normal state quasiparticle spectrum, $\omega(\xi_k)$, in the nodal direction $(0,0) - (\pi,\pi)$ at the energy $\omega_{\text{kink}} \lesssim 70$ meV, which is a characteristic oxygen vibration energy $\omega_{\text{ph}}^{(70)}$. However, the kink is not shifted in the superconducting state, contrary to the prediction of the standard Eliashberg theory [3]. The latter contains integration over the whole Fermi surface and over the energy giving that singularities in $\omega(\xi_k)$ (along all directions) must be shifted in the superconducting state by the maximal gap value $\Delta_0$: (II) In the anti-nodal region, near $(\pi,0)$ or $(0,\pi)$, there is a singularity in $\omega(\xi_k)$ in the normal state at $\omega_{\text{sing}}^{(40)} \approx 40$ meV - which is also a characteristic oxygen vibration energy $\omega_{\text{ph}}^{(40)}$. This singularity is shifted in the superconducting state at $T << T_c$ to $\omega \approx 60$ meV($= \omega_{\text{ph}}^{(40)} + \Delta_0$), where $\Delta_0(\approx 20$ meV) is the maximal superconducting gap at the anti-nodal point. The experimental slopes of $Re \Sigma(k,\omega)$ at the kink (and singularity) give the EPI coupling constant $\lambda_{\text{ph}} > 1$. The different shifts of $\omega_{\text{kink}}^{(70)}$ and $\omega_{\text{sing}}^{(40)}$ in the superconducting state we call the ARPES non-shift puzzle; (III) There is a a “knee”-like structure of $|\ Im \Sigma(\omega)| = |\ Im \Sigma_{\text{ph}}(\omega) + \ Im \Sigma^G(\omega) |$, which is phonon dominated $|\ Im \Sigma(\omega)| \approx |\ Im \Sigma_{\text{ph}}(\omega_{\text{ph}})| \sim \pi \lambda_{\text{ph}} \omega_{\text{ph}}/2$ with $\lambda_{\text{ph}} > 1$ (obtained from $Re \Sigma$) for $\omega \approx \omega_{\text{ph}}^{(70)}$, and for $\omega > \omega_{\text{ph}}^{(70)}$ there is a pronounced linear behavior of $|\ Im \Sigma(\omega)| \approx |\ Im \Sigma_{\text{ph}}(\omega_{\text{ph}})| + \pi \lambda_{C,\omega}/2$, and $\lambda_C < 0.4$ (obtained from $Im \Sigma$) - due to the Coulomb scattering. It turns out that $\lambda_{\text{ph}} \gg \lambda_C$; (IV) There is a dip-hump structure in the spectral function $A(k_F,\omega)$ with the quasiparticle peak sharpening in the superconducting state near the anti-nodal point.

Before explaining these distinctive features by the phonon-type theory we stress, that the ARPES spectra especially the non-shift puzzle and “knee”-like structure, can not be explained by the spin-fluctuation interaction (SFI) due to the following reasons: (i) the intensity of the SFI spectrum $(\approx Im \chi(Q,\omega)$ - the spin susceptibility at $Q=(\pi,\pi))$, although pronounced in slightly underdoped materials, is strongly suppressed (even below the experimental resolution) in the normal state of the optimally doped HTSC oxides [7], although their critical temperatures differ only slightly $(\delta T_c \sim 1 K)$. Such a huge reconstruction of the SFI spectrum around the optimal doping but with small effect on $T_c$ gives strong evidence for the ineffectiveness of the SFI in pairing. (ii) The SFI theory assumes unrealistically large coupling $g_{sf} \approx 0.65$ eV, (with the coupling constant $\lambda_{sf} \approx 2.5$), while the ARPES [8, 9, 10], resistivity [11] and magnetic [12] measurements give much smaller $g_{sf} \lesssim 0.1$ eV, i.e. $\lambda_{sf} < 0.2 < \lambda_C \lesssim 0.4$. Such a small $\lambda_{sf}$ gives small $T_c$; (iii) if the kink at 70 meV in the normal state would be due to the magnetic spectrum, then it would be strongly rearranged in the superconducting state, contrary to the ARPES results. On the other hand, the phonon energies are only slightly $(\lesssim 5\%)$ changed in the superconducting state.; (iv) the magnetic resonance mode at 41 meV, which appears only in the superconducting state [8], can not cause the kink, since the latter is present in $La_{2-x}Sr_xCuO_4$, where there is no magnetic resonance mode at all [8]; (v) The non-shift puzzle and “knee”-like structure of the ARPES spectra are related
to phonon features, which definitely disqualify the SFI and favor the electron-phonon interaction (EPI) as the pairing mechanism. The four distinctive features in the ARPES spectra can be explained by the theory in which the electron-phonon interaction (EPI) with the forward scattering peak (FSP) dominates over the Coulomb interaction - the EPI-FSP model.

The EPI-FSP model - The central question for the EPI theory is - why is the anti-nodal singularity $\omega_{\text{sing}}^{(0)}$ shifted in the superconducting state, but the nodal kink $\omega_{\text{kink}}^{(70)}$ is not? We show that in order to solve the ARPES non-shift puzzle one should go, as it is said before, beyond the standard Eliashberg theory for the EPI.

To remain the reader, the standard Eliashberg theory implies that $\omega_{\text{sing}}^{(40)}$ and $\omega_{\text{kink}}^{(70)}$ should be shifted in the superconducting state to $\omega_{\text{sing}}^{(40)} \to \omega_{\text{ph}}^{(40)} + \Delta_0$ and $\omega_{\text{kink}}^{(70)} \to \omega_{\text{ph}}^{(70)} + \Delta_0$, respectively. Here we show that the ARPES non-shift puzzle can be explained by the EPI-FSP model which contains the following basic ingredients: (1) The EPI is dominant in HTSC and its spectral function $\alpha^2 F(k, k', \Omega)$, which enters the Eliashberg equations below, has a pronounced FSP at $k - k' = 0$, due to strong correlations. Its width is very narrow $|k - k'| \ll k_F$ even for overdoped systems [11, 14]. Near the Fermi surface one expects that $\alpha^2 F(k, k', \Omega) \approx \alpha^2_{ph} F(\varphi', \varphi, \Omega)$ and in strongly correlated systems one has $\alpha^2 F(\varphi, \varphi', \Omega) \sim \gamma^2(\varphi - \varphi')$, where the charge vertex $\gamma(\varphi - \varphi')$ is strongly peaked at $\varphi - \varphi' = 0$ with the width $\delta \varphi_{\text{W}} \ll \pi$ [11, 14]. Thereby, one can put in leading order $\alpha^2_{ph} F(\varphi, \varphi', \Omega) \approx \alpha^2_{ph} F(\varphi, \Omega) \delta(\varphi - \varphi')$, which picks up the main physics whenever $\delta \varphi_{\text{W}} \ll \pi$. The EPI-FSP model, which results from the $t-J$ model with the electron-phonon interaction [11, 14] predicts the following important results: (a) the strength of pairing is due to the EPI, while the residual Coulomb interaction (including spin fluctuations) triggers the pairing to d-wave one; (b) the transport coupling constant $\lambda_{tr}$ entering the resistivity, $\rho \sim \lambda_{tr} T$ is much smaller than the pairing one $\lambda_{ph}$, i.e. $\lambda_{tr} < \lambda_{ph}/3$. We stress that the FSP in the EPI of strongly correlated systems is a general effect by affecting electronic coupling to all phonons. This is an important result, since for some phonons (for instance the half-breathing modes of O ions) the bare coupling constant $g_0^2(q)$ is peaked at large $q \sim 2 k_F$ and therefore detrimental for d-wave pairing, while the one renormalized by strong correlations $g_{2\text{en}}^2(q) = g_0^2(q) \gamma_2^2(q)$ is peaked at much smaller $q$, thus contributing constructively to d-wave pairing. The Monte Carlo calculations on the Hubbard model with the EPI and finite repulsion [12] confirm the existence of the FSP in the EPI - previously found analytically in [11]; (2) the dynamical part (beyond the Hartree-Fock) of the Coulomb interaction is characterized by the spectral function $S_C(k, k', \Omega)$. The ARPES non-shift puzzle implies that $S_C$ is either peaked at small transfer momenta $|k - k'| \ll k_F$, or it is so small that the shift is weakly affected and below the experimental resolution of ARPES. Since the ARPES data give also that the Coulomb coupling constant $\lambda_C < 0.4$ is much smaller than $\lambda_{ph} > 1$, then the kink is practically insensitive to the $k$-dependence of $S_C$. Due to simplicity we assume the former case - see also discussion after Eq.(4); (3) The scattering potential due to non-magnetic impurities has pronounced forward scattering peak - due to strong correlations [11, 14], thus making d-wave pairing robust in the presence of impurities - see more below.

Eliashberg equations for the EPI-FSP model - The Matsubara Green’s function is defined by $(k = (k, \omega_n))$

$$G_k = \frac{1}{i\omega_{\text{k}} - \xi_k - \Sigma_k(\omega)} = -\frac{i\tilde{\omega}_k + \xi_k}{\tilde{\omega}_k^2 + \xi_k^2 + \Delta_k^2}$$

(1)

where $\xi_k = \omega_k + \Delta_k$ are the bare quasiparticle energy, renormalized frequency and gap, respectively [13]. The 2D Fermi surface of HTSC is parameterized by $k = (k_F + k_{\perp}, k_F \varphi)$, where $k_F(\varphi)$ is the Fermi momentum and $k_F \varphi$ is the tangent on the Fermi surface [13]. In that case $\xi_k \approx v_{\text{ph}} k_{\perp}$ and $\int d^2k \approx \int d\varphi \int d^2k_{\perp} v_{\text{ph}}(\varphi) = \int \int N_{\varphi, c} d\xi d\varphi$. After the $\xi$-integration the Eliashberg equations in the FSP model read

$$\tilde{\omega}_{n, \varphi} = \omega_n + \pi T \sum_{m} \frac{\lambda_{n, \varphi}(n - m)\tilde{\omega}_{m, \varphi} + \Sigma_{n, \varphi}}{\sqrt{\tilde{\omega}_{m, \varphi}^2 + \Delta_{m, \varphi}^2}} + \Sigma_{n, \varphi}$$

(2)

$$\Delta_{n, \varphi} = \pi T \sum_{m} \frac{\lambda_{n, \varphi}(n - m)\Delta_{m, \varphi} + \Sigma_{n, \varphi}}{\sqrt{\tilde{\omega}_{m, \varphi}^2 + \Delta_{m, \varphi}^2}}$$

(3)

where $\lambda_1(\varphi)(n - m) = \lambda_{ph, \varphi}(n - m) + \delta_{\text{nn}} \gamma_{1(2), \varphi}$ with the electron-phonon coupling function $\lambda_{ph, \varphi}(n)$

$$\lambda_{ph, \varphi}(n) = 2 \int_0^{\infty} d\Omega \frac{\alpha^2_{ph, \varphi} F_{\varphi}(\Omega)\Omega}{\Omega^2 + \omega_n^2}$$

(4)

Note, that Eqs.(2-3) have a local form as a function of the angle $\varphi$, i.e. the energies at different points on the Fermi surface are decoupled. Just this (decoupling) property of the Eliashberg equations in the EPI-FSP model, is crucial for solving the ARPES non-shift puzzle. The term $\Sigma_{n, \varphi}$ is due to the dynamical Coulomb effects and its calculation is the most difficult part of the problem. $\Sigma^C$ is proportional to the charge vertex $\gamma(\varphi - \varphi')$ and, as we said in (2), we assume that it is also almost "local" on the Fermi surface, although this assumption is not crucial at all, since $\lambda_C \ll \lambda_{ph}$. After the $\xi$-integration it reaches the same form as the second term in $\Sigma_{n, \varphi}$, where $\lambda_{1, \varphi}(n - m)$ is replaced by the Coulomb coupling function $\lambda_{C, \varphi}(n - m)$.

The latter has the same form as Eq.(4) but $\alpha^2_{ph, \varphi} F_{\varphi}(\Omega)$ is replaced by $S_C(\varphi, \Omega)$. ARPES spectra give evidence that $\text{Im}\Sigma^C_{\varphi}(\omega) \approx -\pi \lambda_C(\omega)/2$ at $T < \omega < \Omega_c$ which we reproduce by taking $S_C(\varphi, \Omega) = A_{\varphi} \Theta(\Omega_c - |\omega|)$, where $A_{\varphi}$ is normalized to obtain $\lambda_{C, \varphi} \approx 0.4$. The contribution of the Coulomb
interaction to the gap, \( \Delta_{n,\varphi} \), in Eq.(2) includes the following effects: (i) of the Hartree-Fock pseudopotential - which maximizes \( T_c \) when \( \Delta_{n,\varphi} = \phi = 0 \) and favors unconventional (d-wave) pairing; (ii) of the dynamical part of the Coulomb interaction which is unknown and an approximation for \( \Delta_C \) is needed. The SPI approach assumes that \( \Delta_C(k,\omega_n) \) depends on the dynamical spin susceptibility \( \chi_\omega \). Since \( Im\chi_\omega(q,\omega) \) is peaked at \( Q = (\pi,\pi) \) this term is repulsive and favors d-wave pairing. Although \( \Delta_C \) contributes little to \( \Delta_{n,\varphi} \), it is important to trigger superconductivity from s- to d-wave pairing [1][1].

In Eqs.(1 - 2) non-magnetic impurities are included. Strong correlations induce the FSP in the impurity scattering matrix, being \( t(\varphi,\varphi,\omega) = \gamma_1^2(\varphi - \varphi') \). In leading order one has \( t(\varphi,\varphi,\omega) = \delta(\varphi - \varphi') \), thereby not affecting any pairing. In reality impurities are pair-breaking for d-wave pairing and the next to leading term is necessary. This term is controlled by two scattering rates, \( \gamma_1,\varphi \) and \( \gamma_2,\varphi \), where \( \gamma_1,\varphi > \gamma_2,\varphi \geq 0 \). The case \( \gamma_1,\varphi = \gamma_2,\varphi \) leads to the extreme forward scattering - not affecting \( T_c \), while \( \gamma_2,\varphi = 0 \) means an isotropic and strong pair-breaking scattering [1].

Quasiparticle renormalization - The quasiparticle energy \( \omega(\xi_k) \) is the pole of the retarded Green’s function. For numerical calculations we take for simplicity the Lorentzian shape for \( \alpha^2 F(\omega) \) centered at \( \omega_{ph} \). Since our aim is a qualitative explanation of the ARPES non-shift puzzle, we perform calculations only for moderate coupling constants \( \lambda_{ph,\varphi} \approx \lambda_{ph} = 1, \lambda_C = 0.3 \) in both, the nodal and anti-nodal direction. In fact they can take larger values, i.e. \( \lambda_{ph} < \pi/2 \) especially in the anti-nodal region. It is apparent from Eqs.(1 - 2) that the quasiparticle renormalization is local (angle-decoupled) on the Fermi surface. This behavior is expected to be realized in a more realistic model with the finite width \( \delta\varphi \) when \( \delta\varphi \ll \pi/2 \).

(I) Kink in the spectrum in the nodal direction - The kink at \( \omega_{kink} \approx 70 \) meV in \( \omega(\xi_k) \) means that the quasiparticles moving along the nodal direction \( (\varphi = \pi/4) \) interact with phonons with frequencies up to \( 70 \) meV [3], i.e. \( \alpha^2 F(\omega) \neq 0 \) for \( 0 < \omega \leq 70 \) meV. Since \( \Delta_{\pi/4}(\omega) = 0 \) them the "local" form of Eq.(2) implies that the spectrum \( \omega(\xi_k) \) is not shifted in the superconducting state. Numerical calculations in Fig.1a confirm this analytical result what is in agreement with ARPES results [3]. It is expected that for a realistic phonon spectrum the theoretical singularity in \( \omega(\xi_k) \) (shown in Fig.1a) will be smeared having also an additional structure due to other phonons which contribute to \( \alpha^2 F(\omega) \).

(II) Singularity in the anti-nodal direction - The singularity (not the kink) in \( \omega(\xi_k) \) at \( \omega_{sing}^{(40)} \) in the anti-nodal direction \( (\varphi \approx \pi/2) \) is observed in ARPES in the normal and superconducting state of \( Lo_2-xSr_xCuO_4 \) and BISCO [4]. This means that the quasiparticles moving in the anti-nodal direction interact with a narrower phonon spectrum centered around \( \omega_{ph}^{(40)} \approx 40 \) meV.

Since \( \Delta_{\pi/2}(\omega) \leq \Delta_0 \) then Eq.(1) gives that in the normal state \( \omega(\xi_k) \) is singular at \( \omega_{sing}^{(40)} \), while in the superconducting state the singularity is shifted to \( \omega_{sing}^{(40)} = \pm \omega_{ph}^{(40)} + \Delta_0 ) \). This is confirmed by numerical calculations in Fig.2a for \( \omega(\xi_k) \), and in Fig.2b for \( Im\Sigma(\omega,\varphi) \), for \( \lambda_{ph} = 1 \) and \( \lambda_C = 0.3 \). ARPES spectra give \( \Delta_{ph} > 1 \) in the anti-nodal region.

Note, that the theoretical singularity in Fig.1a is stronger than in Fig.2a, because the calculations are performed for the same temperature, and since \( \omega_{ph}^{(70)} > \omega_{ph}^{(40)} \) the latter singularity is smeared by temperature effects more than the former. The real shape of these singularities depends on microscopic details, such as for instance the presence of the van Hove singularity slightly below the Fermi surface in the anti-nodal region, etc. This will be studied elsewhere.

(III) The “knee”-like shape of \( Im\Sigma(\xi = 0,\omega) \) - The “knee” is shown in Fig.1b for the nodal kink (at \( \omega_{ph} = 70 \) meV) and in Fig.2b for the antinodal singularity (at \( \omega_{ph} = 40 \) meV). In both cases there is a clear “knee”-like structure for \( \omega \) near \( \omega_{ph} \), what is in accordance with the recent ARPES results in various HTSC families [2][2][2][2][2]. From Fig.1b it is seen that for energies \( \omega_{ph}^{(70)} < \omega < \Omega_C \) the linear term is discernible in \( Im\Sigma | = | Im\Sigma_{ph} + Im\Sigma_C | \sim | Im\Sigma_{ph}(\omega_{ph}) | + \pi\lambda_C/\omega/2 \), while for \( \omega \approx \omega_{ph}^{(70)} \) the slope of \( | Im\Sigma(\xi = 0,\omega) | \)
imaginary self-energy  is steeper, since for \( \lambda_{ph} = 1 \gg \lambda_C = 0.3 \) the term \( \text{Im} \Sigma_{ph}(\omega_{ph}) \) dominates. The "knee"-like shape of \( \text{Im} \Sigma(\xi = 0, \omega) \) is seen in Fig. 3a and is also supported by the pronounced long-range Madelung EPI, which is caused by the ionic-metallic character of layered HTSC oxides. \( \lambda_{ph} = 1 \) in the anti-nodal region may influence \( \Sigma_{ph}(\xi = 0, \omega) \) significantly and change its shape too. This will be studied elsewhere.

IV) ARPES dip-hump structure - The EPI-FSP model explains qualitatively the dip-hump structure in \( A(\varphi, \omega) = -\text{Im} G(\varphi, \omega)/\pi \) which was observed recently in ARPES. In Fig. 3a it is seen that the dip-hump structure is realized in the normal state (also in the presence of impurities) already for a moderate coupling constant \( \lambda_{ph} = 1 \). The dip is more pronounced in the superconducting state where the peak in \( A(\omega) \) is appreciable narrowed, what is in accordance with ARPES experiments. Contrary to expectations, the dip-energy does not coincide with the (shifted) phonon energy at \( \omega_{ph} = 40 \) meV. However, the positions of the maxima of \(-dA/d\omega\) appear near the energies \((-\Delta_0 - n\omega_{ph})\) as it is seen in Fig. 3b. The calculations give also a dip in both, the anti-nodal and nodal density of states \( N(\omega) \) (not shown) already for \( \lambda_{ph} = 1 \), which is more pronounced for larger \( \lambda_{ph} (> 1) \).

Discussion and conclusions - In obtaining Eqs. (1 - 4) in the EPI-FSP model the Migdal vertex corrections due to the electron-phonon interaction are neglected. In [10] it is shown that these corrections may increase \( T_c \) significantly, at the same time decreasing the isotope effect, even for \( \lambda_{ph} < 1 \). However, these results cannot change the qualitative picture obtained by the present theory.

In conclusion, the four distinctive features in the quasiparticle self-energy \( \Sigma(\kappa, \omega) \), obtained from the ARPES spectra in HTSC materials, are explained consistently by the theory in which the electron phonon interaction (EPI) with the forward scattering peak (FSP) dominates over the Coulomb scattering. In particular, this theory explains why there is no shift of the nodal kink at 70 meV in the superconducting state, contrary to the clear shift of the anti-nodal singularity at 40 meV. The "non-shift" effect is a direct consequence of the existence of the FSP in the EPI, i.e. due to the long-range character of the electron-phonon interaction in HTSC oxides. This is also supported by the pronounced long-range Madelung EPI, which is caused by the ionic-metallic character of layered HTSC oxides. However, for the quantitative theory the EPI-FSP model must be refined by including the realistic phonon and band structure of HTSC oxides.

Finally, based solely on the ARPES and tunnelling spectra, as well as on dynamical conductivity measurements, one can make a reliable phenomenological theory for the pairing in HTSC oxides. Its basic ingredient is the
electron-phonon interaction (which provides the strength for pairing) with the pronounced forward scattering peak, whatsoever is its cause, while the Coulomb scattering, in spite of its weakness compared to the EPI, triggers superconductivity to d-wave like \[1, 11, 17\]. This phenomenology is described by Eqs.(1-4), or their generalization to a realistic forward scattering peak with small but finite width \(\delta \phi_w \ll \pi\).

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