Self-energy effects and electron–phonon coupling in Fe–As superconductors

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Abstract
Doping and temperature dependent studies of optical phonon modes in Fe-122 pnictides are performed using Raman scattering experiments and compared with model calculations to elucidate the role of electron–phonon and spin–phonon interaction in this family of compounds. The frequency and linewidth of the $B_{1g}$ mode at around $210 \text{ cm}^{-1}$ is highlighted as appreciable anomalies at the superconducting and spin density wave transitions are observed that strongly depend on composition. We give estimates of the electron–phonon coupling related to this renormalization and calculate the phonon self-energy on the basis of a four-band model comparing different symmetries of the order parameters. In addition, we observe a pronounced quasi-elastic Raman response for the undoped compound, suggesting persisting magnetic fluctuations in the spin density wave state.

(Some figures in this article are in colour only in the electronic version)

The recent discovery of superconductivity in iron arsenide compounds containing Fe$_2$As$_2$ layers with $T_c$ higher than 50 K has stimulated enormous research activity [1–5]. This is largely due to the tantalizing possibility of finding a unified framework and description of a diverse range of unconventional superconducting materials.

Like for the high $T_c$ cuprates, there is evidence for the importance of many-body effects. Inelastic neutron scattering (INS) measurements have uncovered a magnetic resonance peak for superconducting Ba$_{1-x}$K$_x$Fe$_2$As$_2$ [6]. The electron–phonon coupling is estimated to be a factor of five too weak to reproduce the observed transition temperature, even taking account of multiband effects [7]. Besides, the spin density wave (SDW) and the phase diagram with doping indicate an intimate relation between magnetism and superconductivity.

A combination of the specific band structure of the Fe pnictides [8] with antiferromagnetic correlations has led to an $s^\pm$ wave pairing state ($\Delta_{k} = \Delta_{k}(\cos k_x + \cos k_y)$). The $s^\pm$-gap symmetry assumes a change of sign of the superconducting gap between electron and hole pockets and the gap’s magnitudes differing slightly on the hole and electron pockets. There is also a fictitious line of nodes between pockets which, however, does not cross the Fermi surface. This scenario is further supported by several theoretical works [9–11]. Although some spectroscopic experiments are consistent with $s^\pm$ wave symmetry [3, 6, 12], there also exist data [13, 14] which cannot be easily fitted with a nodeless gap symmetry [15]. A nonuniversal gap symmetry can be resolved assuming a relatively large intraband Coulomb repulsion within each of the pockets. Depending on the model parameters, an extended $s$ wave gap with a $\cos k_x/2\cos k_y/2$ dependence and nodal lines crossing electron pockets centered around the M point [16, 17] or different symmetries like the $d_{x^2−y^2}$ wave (cos $k_x − \cos k_y$) symmetries [18, 19] can be obtained.

The prevailing trust in a magnetically mediated pairing mechanism with moderate electron–phonon coupling has been further justified by the observation of an inverse isotope effect on iron with $\alpha \sim −0.18$ in Ba$_{1-x}$K$_x$Fe$_2$As$_2$ [20, 21]. At the same time, Eschrig [22] asserted that an in-plane $B_{1g}$ mode...
phonon coupling are needed. In this paper, we provide further combined experimental/theoretical studies of electron–temperature below 130 K (see figure 1(c)). The full width at using the excitation line refrigerator. Raman scattering experiments were performed of the single crystal. The scattered spectra were collected by a DILOR-XY triple spectrometer and a nitrogen cooled charge-coupled device detector.

Figure 1(a) shows the Raman response Imχ of the undoped SrFe2As2 for (xx) and (xy) polarizations at 290 K, which is corrected by the Bose thermal factor [1 + n(ω)] = [1 − exp(−ω/kB T)]−1 from the measured Raman scattering intensity. At room temperature we observe a single peak at 203 cm−1 in the (xx) polarization. This is part of four symmetry-allowed modes: ΓRaman = A1g(x2 + y2, z2) + B1g(xz, yz) + 2Eg(xz, yz, z2) [26, 27]. The 203 cm−1 mode is assigned to a B1g mode and corresponds to a displacement of Fe atoms along the c axis. The room temperature Raman response exhibits a structured background, which might be composed of a phonon density of states and weak electronic Raman scattering. To remove extrinsic effects, like defect scattering and contributions from the cryostat windows, Imχ(T = 290 K) is subtracted from Imχ(T). The resulting Raman response is plotted in figure 1(b) as a function of temperature. We observe a quasi-elastic scattering maximum, which is well described by a Lorentzian profile, Imχ ∝ AΓ/(ω2 + Γ2), where A is the scattering intensity and Γ is the full width at half-maximum. With decreasing temperature the scattering intensity grows steeply through Tc of the SDW and then shows a saturation for temperature below 130 K (see figure 1(c)). The full width at half-maximum tends to go to zero quasi-linearly upon cooling (see figure 1(d)). A central maximum can arise from the decay of a soft mode into acoustic modes or phonon density fluctuations in the presence of the structural phase transition. In our case, however, this does not give a dominant contribution because the structural phase transition is of first order [28]. Actually, the intensity of the elastic maximum does not diverge at Tc. Instead, it looks similar to the temperature dependence of the elastic neutron scattering intensity at the AF superlattice reflection [32]. Thus, it is ascribed to light scattering by low energy magnetic fluctuations. It should be noted that the Bose-corrected Raman scattering intensity, i.e. the imaginary part of the Raman response function, should vanish in the limit Δω → 0 to fulfill causality [23]. This means that for energies below our window of observability (Δω > 30 cm−1) additional low energy intensity develops with decreasing temperatures. Similar Lorentzian-lineshaped fluctuations are observed, e.g. in low dimensional quantum spin systems due to pronounced energy density fluctuations [33]. Therefore it is tempting to attribute the quasi-elastic scattering and the linear dependence of its linewidth to strong magnetic fluctuations at low energies, i.e. to the proximity of the SDW state to a quantum phase transition [34, 35].

In the following we will focus on Raman scattering on optical phonons. In order to understand the evolution of the SDW state the temperature dependence of the 203 cm−1 B1g mode was analyzed by using a Lorentzian profile. In figure 2 the results are summarized together with those for BaFe2As2 (Tc = 138 K) and CaFe2As2 (Tc = 173 K). We note that this mode is susceptible to any change of the Fe d states around the Fermi level. The phonon frequency and linewidth show characteristic anomalies in the temperature dependence [36]. For all samples
we observe a jump of the phonon frequency and a narrowing of the linewidth below $T_s$. Similar phonon anomalies are observed in the $\Lambda_{1g}$ mode [30], while the $e_g$ phonon modes show a splitting related to its coupling to the structural distortions of the SDW [37]. Thus, we attribute the narrowing of the linewidth below $T_s$ to a longer phonon lifetime in the SDW state and the depletion of electronic states on the Fermi surface. In contrast, we do not find a change of the scattering background that can be partially attributed to electronic Raman scattering. This suggests that the SDW gap is not fully opened on the Fermi surface.

This is consistent with the ARPES experiments [38] on the undoped pnictides where significant deviations from a complete nesting of the electron and hole Fermi surfaces have been found. Our results indicate that some of the non-nested Fermi surface portions will survive below $T_s$, producing a finite linewidth below $T_s$. For comparison, results for the diamagnetic isostructural compound FeAs$_2$ are presented as well. Its temperature dependence of the frequency and linewidth of the $B_{1g}$ mode is consistent with the modeled anharmonic phonon decay processes [42]:

$$\omega_p(T) = \omega_0 + C \left[1 + 2/\exp(h\omega_0/2k_B T) - 1\right]^{-1}$$

where $\omega_0$ is the zero-temperature frequency of the $B_{1g}$ mode and $C$ corresponds to the anharmonic phonon decay processes. Also Ba$_{0.72}$K$_{0.28}$Fe$_2$As$_2$ follows largely the modeled anharmonic behavior. For Sr$_{0.85}$K$_{0.15}$Fe$_2$As$_2$ with a higher SDW transition temperature, however, we can find sizable deviations from the anharmonic profile. The frequency undergoes a small hardening by 1 cm$^{-1}$ at the superconducting $T_c$. This means that although the overall temperature dependence of the frequency is determined by the lattice anharmonicity, there are nonnegligible superconductivity-induced self-energy contributions. Interesting to note is that the SDW lead to an onset of anomalies in the phonon linewidth below 180 K for the Sr-based pnictide. For the Ba pnictide these anomalies for $T < 150$ K are much weaker. The overall decrease of linewidth with temperature is very pronounced and of similar magnitude to that for the undoped samples.

We now turn to the superconducting samples Sr$_{0.85}$K$_{0.15}$Fe$_2$As$_2$ and Ba$_{0.72}$K$_{0.28}$Fe$_2$As$_2$. The respective temperature dependence of the frequency and linewidth of the $B_{1g}$ mode is shown in figure 3. In order to identify possible electron–phonon contributions, data for the diamagnetic isostructural compound FeAs$_2$ are presented as well. Its temperature dependence of the phonon frequency is well described in terms of simple phonon–phonon decay processes [42]:

$$\omega_p(T) = \omega_0 + C \left[1 + 2/\exp(h\omega_0/2k_B T) - 1\right]^{-1}$$

Figure 2. (a) Temperature dependence of the normalized linewidth and (b) peak position of the 203 cm$^{-1}$ mode on a reduced temperature scale, $T/T_s$. For comparison, SrFe$_2$As$_2$ ($T_s = 202$ K) is presented together with BaFe$_2$As$_2$ ($T_s = 138$ K) and CaFe$_2$As$_2$ ($T_s = 173$ K). The dashed, dotted, and solid lines are guides to the eyes.

Figure 3. Temperature dependence of the peak position (upper panel) and linewidth (lower panel) of the $B_{1g}$ mode for the Sr$_{0.85}$K$_{0.15}$Fe$_2$As$_2$ ($T_c = 28$ K) and Ba$_{0.72}$K$_{0.28}$Fe$_2$As$_2$ ($T_c = 32$ K) single crystals. For comparison, results for the diamagnetic isostructural compound FeAs$_2$ are added. The solid lines are a fit to an anharmonic phonon decay process. The striped bar highlights the onset of the superconducting transition regime.

To analyze the effect of superconductivity on the renormalization of the $B_{1g}$ phonons we employ the four-band model proposed previously [11] for the folded Brillouin zone. The unit cell contains two Fe and two As atoms and the band structure predicts a Fermi surface consisting of two hole ($\alpha$)
pockets centered around the Γ point and two electron (b) pockets centered around the M point, respectively. A similar analysis has been previously performed for the buckling B1g mode in cuprates within a single-band model [24].

Without taking into account vertex corrections, the renormalization of the optical phonons is determined by the Dyson equation:

\[ D^{-1}(\mathbf{q}, \omega) = D_0^{-1}(\omega) - g_q^2 \Pi(\mathbf{q}, \omega), \]

where \( D_0(\omega) = \frac{2\omega}{\omega^2 - \omega_{\mathbf{q}}^2} \) is the momentum independent bare phonon propagator and \( g_q \) is the corresponding electron–phonon coupling constant. The polarization operator is given by

\[ \Pi(\mathbf{q}, \omega) = -i \int \text{Tr}[\tau_3 G(\mathbf{k} + \mathbf{q}, \Omega + \omega) \tau_3 G(\mathbf{k}, \Omega)] \frac{d^2 k d\Omega}{(2\pi)^3}, \]

where \( G(\mathbf{k}, \omega) = \frac{\omega^2 - \mathbf{k}^2 - \delta_\mathbf{k}}{\omega^2 - E_\mathbf{k}^2} \) is the propagator, \( E_\mathbf{k}^2 = \epsilon_\mathbf{k}^2 + \Delta_\mathbf{k}^2 \) is the energy-binding dispersion in the superconducting state, \( \epsilon_\mathbf{k} \) are the tight-binding energies for the α and β bands [11], and \( \Delta_\mathbf{k} \) is a (momentum dependent) superconducting gap.

The various symmetries of the superconducting gaps will renormalize the polarization operator in the superconducting state differently. The main effect of superconductivity on the phonon self-energy at \( q = 0 \) results from the change of the polarization operator of the two α bands. The latter, which are centered close to the Γ point, consequently couple rather strongly to the phonon dispersions around the center of the BZ. In figure 4(a) we show the changes of the real and imaginary parts of the polarization operators of the α bands for the extended s wave (s^+) and d_x^2−y^2 wave symmetries of the superconducting gap with respect to the normal state values. One can clearly see that the largest effect occurs for the extended s wave symmetry. The real part of the polarization operator in the superconducting state is larger for the extended s wave symmetry. The respective results are shown in figure 4(b).

The inset shows the experimental data for Sr_{0.15}K_{0.15}Fe_2As_2 extracted from figure 3.

![Figure 4](image)

**Figure 4.** (a) Calculated difference of the polarization operator for the two α bands, \( \Pi(\mathbf{q} \to 0, \omega) \) between the superconducting and normal states. The solid and dashed curves refer to the real and imaginary components, respectively. (b) Calculated normalized temperature dependent B1g frequencies for temperatures below \( T_c \). The inset shows the experimental data for Sr_{0.15}K_{0.15}Fe_2As_2 extracted from figure 3.

Evidence for coexistence, phase competition or separation on the superconducting transition temperature assuming an extended s± wave symmetry. The respective results are shown in figure 4(b). Using the previous estimates for the electron–phonon coupling strength of \( g_{\mathbf{q}=0} \approx 24.8 \) meV and the bare \( \hbar \omega_{\mathbf{B}_{1g}} \approx 26.2 \) meV in iron pnictides [44], we calculate the renormalization of the phonon frequency and find its hardening in the normal state at \( T = 100 \) K to 26.3 meV which corresponds to 210.4 cm\(^{-1}\). Below \( T_c \), Re \( \Pi(\omega) \) further decreases at energies larger than \( 2\Delta_0 \approx 13.5 \) meV which agrees well with the experimental data shown in the inset of figure 4(b). At the same time, we also find a broadening of the phonon mode below \( T_c \) (not shown) though experimentally our data show a quick switch of the broadening tendency into a sizable narrowing. Such change may occur only if \( \hbar \omega_{\mathbf{B}_{1g}} \sim 2\Delta_0 \) or slightly smaller, which would yield anomalously large \( 2\Delta_0/\hbar \omega_{\mathbf{B}_{1g}} \) ratios. One of the possible explanations could be microscopic coexistence of SDW and superconductivity or interband scattering effects which we leave for further studies. Evidence for coexistence, phase competition or separation on
different length scales has indeed been claimed [45]. These effects seem to depend in a critical way on the stoichiometry and therefore deserve further investigations.

Now we estimate the electron–phonon coupling strength. The linewidth of the isostructural compound FeAs$_2$ is given as 3–4 cm$^{-1}$ between 4 K and room temperature. Taking into account the linewidth of 3–7 cm$^{-1}$ and the superconductivity related narrowing of 1–2 cm$^{-1}$ in Ba$_{0.72}$K$_{0.28}$Fe$_2$As$_2$ and Sr$_{0.85}$K$_{0.15}$Fe$_2$As$_2$, the electron–phonon contribution does not exceed 2 cm$^{-1}$. The Allen equation provides a relation between the phonon linewidth, $\Gamma$, due to electron–phonon coupling and the phonon coupling constant [46]: $\Gamma = 2\pi \lambda_B N(0) \omega^2$ where $\lambda_B$ is the strength of the electron–B$_{1g}$ coupling and $N(0)$ is the density of states (DOS) at the Fermi surface. For BaFe$_2$As$_2$, the total DOS at $E_F$ is taken as $N(0) = 3.06$ eV$^{-1}$/f.u. [47]. If we assume that $N(0)$ remains constant for a small doping level, the electron–phonon constant is estimated to be $\lambda_B \approx 0.2$.

We can also obtain the electron–phonon constant at the Brillouin zone center using the superconductivity-induced renormalization constant $\kappa = (\omega^2/\omega^N) - 1 \approx 0.5\%$ [43]; $\lambda_B^{\text{point}} = -\kappa \Re \{\sin(u/\pi)\} \approx 0.01$, where $u = \pi + 2i \cosh^{-1}(\omega N/2\Delta)$. Both values are much smaller than the theoretically estimated total average value of $\lambda \approx 0.2$ [7]. Our results indicate that conventional electron–phonon coupling is weak for the B$_{1g}$ phonon mode. As indicated by the inverse isotope coefficient, however, we cannot rule out the significance of unconventional electron–phonon coupling due to multiband effects.

To conclude, we have presented a Raman scattering study of undoped SrFe$_2$As$_2$ and superconducting Ba$_{0.72}$K$_{0.28}$Fe$_2$As$_2$ and Sr$_{0.85}$K$_{0.15}$Fe$_2$As$_2$ as a function of temperature. We observe a superconductivity-induced self-energy effect of the B$_{1g}$ phonon mode and estimate the corresponding electron–phonon strength of $\lambda_B^{\text{point}} \approx 0.02$.

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