Superconductivity of underdoped PrFeAs(O,F) investigated via point-contact spectroscopy and nuclear magnetic resonance

Original
Superconductivity of underdoped PrFeAs(O,F) investigated via point-contact spectroscopy and nuclear magnetic resonance / Daghero, Dario; Piatti, Erik; Zhigadlo, Nikolai D.; Ummarino, Giovanni A.; Barbero, Nicolò; Shiroka, Toni. - In: PHYSICAL REVIEW. B. - ISSN 2469-9969. - STAMPA. - 102(2020), p. 104513.

Availability:
This version is available at: 11583/2847020 since: 2020-09-29T14:56:09Z

Publisher:
American Physical Society

Published
DOI:10.1103/PhysRevB.102.104513

Terms of use:
openAccess
This article is made available under terms and conditions as specified in the corresponding bibliographic description in the repository

Publisher copyright
Superconductivity of underdoped PrFeAs(O,F) investigated via point-contact spectroscopy and nuclear magnetic resonance

D. Daghero* and E. Piatti
Department of Applied Science and Technology, Politecnico di Torino, 10129 Torino, Italy

N. D. Zhigadlo†
CrystMat Company, CH-8046 Zurich, Switzerland

G. A. Ummarino
Department of Applied Science and Technology, Politecnico di Torino, 10129 Torino, Italy and National Research Nuclear University MEPhI (Moscow Engineering Physics Institute), Kashira Hwy 31, Moscow 115409, Russia

N. Barbero, T. Shiroka‡
Laboratory for Muon-Spin Spectroscopy, Paul Scherrer Institut, CH-5232 Villigen PSI, Switzerland and Laboratorium für Festkörperphysik, ETH Zürich, CH-8093 Zurich, Switzerland

(Dated: September 29, 2020)

Underdoped PrFeAs(O,F), one of the less known members of the 1111 family of iron-based superconductors, was investigated in detail by means of transport, SQUID magnetometry, nuclear magnetic resonance (NMR) measurements and point-contact Andreev-reflection spectroscopy (PCARS). PCARS measurements on single crystals evidence the multigap nature of PrFeAs(O,F) superconductivity, shown to host at least two isotropic gaps, clearly discernible in the spectra, irrespective of the direction of current injection (i.e., along the $ab$ planes or along the $c$ axis). Additional features at higher energy can be interpreted as signatures of a strong electron-phonon coupling, as demonstrated by a model which combines Andreev reflection with the Eliashberg theory. Magnetic resonance measurements in the normal phase indicate the lack of a magnetic order in underdoped PrFeAs(O,F), while $^{75}$As NMR spin-lattice relaxation results suggest the presence of significant electronic spin fluctuations, peaking above $T_c$ and expected to mediate the superconducting pairing.

I. INTRODUCTION

The discovery of superconductivity in LnFeAsO oxypnictides ($Ln_{1111}$, $Ln$: lanthanide) generated a widespread interest among the condensed matter physicists. These compounds, which belong to the 1111 family of Fe-based superconductors, exhibit a ZrCuSiAs-type structure, composed of alternating stacks of LnO and FeAs layers. They become superconductors either through chemical substitution at the different atomic sites, or through the application of external pressure. Consequently, the resulting electronic phase diagrams depend sensitively on the particular doping element. In either case, the original antiferromagnetic state is partially or fully suppressed. In particular, it has been shown that the LnFeAsO parent compounds can be doped with holes by partially replacing the $Ln^{3+}$ ions with divalent ions, such as $Sr^{2+}$, as e.g. in La$_{1-x}$Sr$_x$FeAsO$^5$ or Pr$_{1-x}$Sr$_x$FeAsO$^6$. By converse, $n$-type doping can be achieved by substituting $Ln^{3+}$ with tetravalent ions, such as $Th^{4+}$ (Sm$_{1-x}$Th$_x$FeAsO)$^{7,8}$, or by partially replacing $O^{2-}$ with $F^-$ or $H^-$ ($Ln$FeAs$_{1-x}$O$_x$F$_x$, LnFeAsO$_{1-x}$H$_x$)$^{9,10}$. In addition, in case of isovalent doping of the $Ln_{1111}$ parent compound, as e.g., in the $As_{1-x}P_x$ case, one can tune the magnetic interactions without changing the carrier concentration$^{9,10}$. Until now, the electron-doped $Ln_{1111}$-type oxypnictides ($O_{1-x}F_x$ and $O_{1-x}H_x$) seem to exhibit the highest $T_c$'s. The control of $T_c$ through carrier concentration is, therefore, a versatile and powerful mean of elucidating the intrinsic nature of superconductivity.

In most $Ln_{1111}$ families, an increase in doping level shifts the system from an antiferromagnetically (AF) ordered state towards a purely superconducting (SC) state, via a region where the AF and SC phases coexist$^{11-13}$. By contrast, in PrFeAsO$_{1-x}$F$_x$, the Neél order (and the tetragonal-to-orthorhombic structural transition) appear to vanish rather rapidly, possibly in a first-order-type transition, as the fluorine concentration approaches the critical value $x \sim 0.08$14. This behavior has been observed to occur also in the LaFeAsO$_{1-x}$F$_x$ family15. On the other hand, it differs significantly from the structurally related families (where the $Ln$ ion is, e.g., Sm, Nd, Ce, etc.), whose AF- to-SC transitions are much more extended. To investigate this in further detail, homogeneously underdoped samples, preferentially in a single-crystalline form, are required.

To date, despite extensive evidence that superconductivity in Fe-based materials is mediated by spin fluctuations, a conclusive experimental confirmation is still missing. In particular, the interplay between the AF fluctuations and superconductivity in the underdoped regime remains unclear, mainly reflecting the difficulties associated with the preparation of high-quality underdoped $Ln_{1111}$ samples. The first step toward the elucidation of the nature of superconductivity is the growth of high-quality crystals.

Here, we report on advanced point-contact Andreev reflection spectroscopy (PCARS) and nuclear magnetic resonance (NMR) studies of superconductivity in underdoped PrFeAs$_{1-x}$F$_x$ crystals with a $T_c$ of $\sim 24$ K. Through an exhaustive set of measurements, we directly assess the multi-
gap nature of superconductivity in this compound and determine the amplitudes of the gaps, that appear to be isotropic in-plane and out-of plane, with no evidence of extended node lines. We bring evidence of a strong coupling between electrons and a bosonic mode, whose characteristic energy agrees well with that of spin fluctuations. Finally, we show that the magnetic order, typical of the parent compounds, is completely suppressed in these underdoped crystals, while sizable spin fluctuations persist, as indicated by NMR. Altogether, these results strongly point towards a spin-fluctuation mediated multiband superconductivity in PrFeAsO$_{1-x}$F$_x$.

II. CRYSTAL GROWTH AND EXPERIMENTAL DETAILS

The PrFeAs(O,F) crystals were grown by using a cubic-anvil high-pressure high-temperature technique. The details of the setup can be found in Ref. 16 and 17. Starting powders of PrAs, FeF$_2$, Fe$_2$O$_3$, and Fe of high purity ($\geq 99.95\%$) were weighed according to the stoichiometric ratio, thoroughly grounded in a mortar and then mixed with NaAs flux. For one growth batch we used 0.45 g of PrFeAsO$_{0.60}$F$_{0.35}$ and 0.2 g of NaAs. The crystal growth process was performed by heating the mixture up to $\sim 1500$°C in 2 h. The mixture was kept there for 5 h, cooled to 1250°C in 6 h, held at this temperature for 3 h, and finally cooled down to room temperature. The crystalline products were separated by dissolving the flux in distilled water. Further details on the crystal growth of PrFeAs(O,F) can be found in Ref. 17.

The x-ray analysis confirmed that the obtained crystals belong to the 1111-type structure, with the refined model being consistent with that from our previous x-ray diffraction studies (see Table 1 in Ref. 18). Compositional analysis via energy-dispersive x-ray (EDX) measurements confirmed that the ratio of praseodymium, iron, and arsenic is close to 1:1:1. Light elements such as oxygen and fluorine cannot be measured accurately via EDX. Therefore, we could not determine the exact doping level of the PrFeAs(O,F) crystals. Nevertheless, by a comparison of our transition temperatures with those of polycrystalline samples (see below), we estimate an F doping of $\sim 0.1$ in our case.

The details of the PCARS technique are given in App. A. As for the NMR study, this consisted in $^{75}$As lineshape- and spin-lattice relaxation measurements, performed at 7.057 T over a temperature range from 4 to 295 K. The NMR signals were detected by means of standard spin-echo sequences, consisting in $\pi/2$ and $\pi$ pulses of 3 and 6 $\mu$s, with recycling delays ranging from 0.01 to 1 s. The lineshapes were obtained via fast Fourier transform (FFT) of the echo signal. Spin-lattice relaxation times $T_1$ were measured via inversion recovery, by using a $\pi-\pi/2-\pi$ pulse sequence.

III. SUPERCONDUCTING PROPERTIES

A. Preliminary characterization of superconductivity

The dependence of magnetic susceptibility vs. temperature in a single PrFeAsO$_{1-x}$F$_x$ crystal, measured in a magnetic field of 0.2 mT parallel to the c axis, is shown in Fig. 1(a). Here, the effective superconducting transition temperature $T_c$,eff is defined as the crossing of the linear extrapolations from the two regions of the high-temperature normal state and low-temperature superconducting state. In the underdoped case the transition is relatively sharp, indicative of a good sample quality. The resistance was measured by using a standard four-probe technique, with the current flowing in the ab plane. Upon lowering the temperature, as shown in Fig. 1(b), the resistance first decreases linearly, to reach a broad minimum around 70 K, and then it increases again. Similar features were also observed in the other investigated crystals. A closeup of the superconducting transition region is shown in the inset. After a saturation around 30 K the resistance starts dropping and reaches its zero value at 23.5 K, fully consistent with the onset of the magnetic transition as measured via SQUID magnetometry. This behavior (and the

FIG. 1. (a) Temperature dependence of magnetic susceptibility of a single PrFeAs(O,F) crystal. The zero-field-cooled (ZFC) curve was obtained on heating in a magnetic field of 0.2 mT applied along the c direction. (b) Resistance as a function of temperature. The inset shows a closeup of the superconducting transition.
A very tiny kink, while a shows the conductance curves of a 74-Ohm, 20 2 and 21). It is worth noting also that the position of the shows typical conductance curves, measured at 2.7 K, for contacts made either on the top surface (c-axis contacts) or on the side (ab plane contacts). The labels refer to the direction of current injection, i.e., perpendicular or parallel to the FeAs planes, respectively. Despite the different shapes of the curves, the different directions of current injection, and the different resistance of the contacts, it is clear that they all show structures at approximately the same energies. In particular, the position of the low-energy maxima (±3.5 meV, solid vertical lines) is very robust. Additional features, that can take the form of maxima, shoulders, or slope changes are present at about ±7.5 mV (dashed vertical lines). These two values are particularly interesting because a full SC gap with an amplitude 1.6k_BTc = 3.5 meV was observed in underdoped PrFeAs(O,F) single crystals by microwave penetration depth and by quasiparticle conductivity measurements19, while a gap of 3.5k_BTc ≈ 7.5 meV was detected by optical conductivity measurements20. Other structures whose occurrence is apparently less systematic can be observed at higher energies (arrows). From the spectroscopic point of view, the fact that there are structures whose position does not depend on the resistance of the contacts means that: i) all the contacts are spectroscopic, at least at low temperature; ii) these structures are intrinsic, i.e., unrelated to the contact, but instead directly connected to the properties of the material. In particular, they are suggestive of the presence of multiple (at least two) superconducting energy gaps. This is a rather common feature of Fe-based systems (including here materials of the same family, such as La-111121 and Sm-111122). It is worth noting also that the position of the spectral features does not depend on the direction of current injection, which suggests that the system does not show a clear in-plane/out-of-plane anisotropy, at least for the gap amplitudes. Clearly, one cannot exclude small anisotropies (i.e., k dependence of the SC gaps), undetectable by our technique.

Finally, the shape of the curves and, in particular, the absence of zero-bias maxima indicates the absence of significant contributions of low-energy quasiparticles to the conductance, thus suggesting a fully-gapped SC and the absence of node lines (as also demonstrated by microwave penetration depth measurements in underdoped single crystals19). A similar situation was observed in other 1111 compounds like F-doped Sm-1111 and La-111122. By contrast, the PCARS spectra of some 122 systems, featuring accidental node lines, show zero-bias maxima at least in one of the directions of current injection; a typical example being the Ca-122 system23,24. Based on the spectra shown in Fig. 2, and consistently with the results of penetration depth19, critical field19, and infrared spectroscopy20 measurements, as well as with electronic band-structure calculations for 1111 compounds, from now on we will assume PrFeAs(O,F) to be a multiple-gap, s±-wave superconductor.

Figure 3 shows the conductance curves of a 74-Ohm, ab-plane contact on a PrFeAs(O,F) single crystal as a function of temperature. The lowest-temperature spectrum was already depicted in the second-last panel of Fig. 2. The temperature dependence is crucial in identifying the critical

---

**FIG. 2.** Low-temperature (2.7 K) conductance curves for various contacts made on the same single crystal. Despite the different shapes, all the curves display structures (maxima, shoulders, slope changes) at approximately the same positions, as indicated by the vertical lines at ±3.5 mV and ±7.5 mV. Arrows indicate structures related to the strong electron-boson coupling.
temperature $T_A$ where the Andreev signal disappears, thus allowing us to determine the normal-state conductance of the contact. As a matter of fact, the progressive decrease in amplitude of the spectra on increasing temperature is due to the decrease in amplitude of the superconducting gap(s). Hence, $T_A$ can be identified with the temperature at which the curves recorded at increasing temperatures start to overlap and the Andreev-reflection features disappear. In our case, the gap vanishes between 24 K and 25 K, since the 25-K conductance curve is superimposed to those recorded at 26 K and 27 K. Thus, the curve recorded at 25 K represents the normal-state spectrum of the contact and we will assume $T_A = 24.5 \pm 0.5$ K. Note that this value is consistent with the $T_c$ obtained from magnetometry and transport data (see Fig. 1). This, again, is an indirect, yet quite convincing proof of the spectroscopic nature of the contact. Indeed, if the conduction through the contact were diffusive, Joule heating would occur within the contact and the Andreev signal would disappear at a lower bath temperature.

To extract the gap amplitudes more accurately, the conductance curves must first be normalized and then fitted to a suitable model. The normalization is obtained by dividing the differential conductance recorded at a given $T < T_c$ by the normal-state conductance of the same contact. The curve measured just above $T_c$ can be used, under the reasonable assumption that the normal-state properties do not change much between $T$ and $T_c$. The lowest-temperature curve in Fig. 3, once divided by the normal-state curve recorded at 27 K (and symmetrized, to better highlight the intrinsic structures and to suppress noise fluctuations) is shown in Fig. 4(a) (circles). On top of it, we plot some theoretical curves, obtained through an automatic fitting procedure based on the minimization of the sum of squared residuals. The model used to fit the experimental data is a two-band, 2D version of the Blonder-Tinkham-Klapwijk (BTK) model. This model contains as free parameters the amplitude of the gaps $\Delta_1$ and $\Delta_2$, the broadening parameters $\Gamma_1$ and $\Gamma_2$, the barrier parameters $Z_1$ and $Z_2$ and the relative weight of band 1 in the conductance, $w_1$. These parameters are not completely free, because the values of the gaps reflect the position of the maxima and shoulders, and the barrier parameters determine the percentage of tunnel vs. Andreev-reflection conduction through the junction and, in practice, are related to the depth of the zero-bias minimum and to the shape of the curve between $\Delta_1$ and $\Delta_2$. Details of the fitting procedure can be found elsewhere. The blue curve in Fig. 4(a) is an attempt to fit the conductance data across the whole voltage range. The overall fit seems fairly good, yet the fitting function completely
fails to reproduce the features at $\pm 7.5 \text{ mV}$, most likely reflecting a superconducting gap (see inset). Moreover, the amplitude of the large gap $\Delta_2 = 15.6 \text{ meV}$ is far too big for a system with $T_c = 25 \text{ K}$, since the gap ratio $2\Delta_2/k_B T_c$ would be 14.5. This value is completely unreasonable even though, in other compounds of the 1111 family, the (larger) gap ratio can be as high as 8.21,22. Finally, the values of the $\Gamma$ parameters are too high, and comparable to the gap values themselves, which should not happen in a spectroscopic contact. The fit is thus unsatisfactory and meaningless. The reason is that, as already demonstrated in the case of Ba(Fe,Co)$_2$As$_2$,28 SmFeAs(O,F),29 and Fe(Fe,Se)$_2$,30 in Fe-based superconductors the relatively strong coupling between the electrons and bosons that mediates the Cooper pairing gives rise to additional structures (shoulders) in the tunnel- and PCARS spectra, better seen as peaks in the second derivative $-d^2I/dV^2$, that do not occur at the gap edge, but at a higher energy $E_p$. As discussed in Ref.29, in case of multiple gaps this energy is $E_p \approx \Delta_{\text{max}} + \Omega_0$, where $\Omega_0$ is the characteristic boson energy and $\Delta_{\text{max}}$ is the largest gap. The electron-boson interaction does not affect the spectra in the energy region where the gap features are observed, but it gives rise to shoulders that can extend to rather high energies and can enormously enhance the apparent width of the conductance curve. These structures cannot be fitted by the Blonder-Tinkham-Klapwijk model, even in its various extended versions, if energy-independent gaps are used as in the BCS, weak-coupling theory. To include the effects of strong coupling in the theory, a much more complicated procedure has to be used, which involves the solution of the Eliashberg equations (see Appendix B). As demonstrated in Appendix B, since the BCS theory represents the low-energy limit of the Eliashberg theory, the low-bias region of the spectrum is completely and uniquely determined by the Andreev reflection. Hence, in this region, the BTK model with constant gaps can be safely used to extract the gap values.31 In Fe-based superconductors, the superconductivity is thought to be mediated by spin fluctuations. Indeed, the position of the electron-boson structures we observe in the aforementioned materials agrees well with a characteristic boson energy $\Omega_0$ that obeys the empirical law $\Omega_0 \approx 2T_c/5$, where $T_c$ is in Kelvin and $\Omega_0$ in meV.32 In our samples, $\Omega_0 \approx 10 \text{ meV}$ and the structures are expected to fall at energies larger than the maximum gap amplitude, that we will call $\Delta_2$.

We have thus to abandon the idea of fitting the whole curve, and focus instead on the low-energy region that hosts the structures related to the gaps, i.e., on the region $|V| \leq V_{\text{max}}$. The choice of $V_{\text{max}}$ is somewhat arbitrary and can (slightly) values of the energy gaps. The red curve in Fig. 4(a) was obtained by setting $V_{\text{max}} = 10 \text{ mV}$ which implies much more reasonable values for the parameters (reported in the labels). In particular, the value of the small gap $\Delta_1 = 3.53 \text{ meV}$ is perfectly compatible with the results of penetration depth and theryarticle conductivity measurements.19

The fact that the high-energy tails of the unnormalized curves (Fig. 3) are affected by the electron-boson structures, and the fact that these structures depend on the energy gap and disappear only at $T_c$,28,29, means that also the normalization is somewhat arbitrary. The usual criterion, i.e., that the high-voltage tails ($V > 3\Delta_0$) of the conductance curves must fall on top of the normal-state conductance curve without any shift. The blue line is a single-gap fit, while the red curve is a two-gap fit, both made with $|V_{\text{max}}| = 10 \text{ mV}$. The inset shows a closeup of the low-bias region.

![Normalized differential conductance at 2.7 K (circles) and relevant fit (red line) for a c-axis contact with a resistance of 8.7 $\Omega$. (b) Open circles: low-temperature (2.7 K) conductance spectrum of a 179-$\Omega$, ab-plane contact, divided by the normal-state conductance curve without any shift. The blue line is a single-gap fit, while the red curve is a two-gap fit, both made with $|V_{\text{max}}| = 10 \text{ mV}$. The inset shows a closeup of the low-bias region.](image_url)
value exhibits a certain variability, depending on the normalization. In the spectrum shown in Fig. 5(a), recorded at 2.7 K in a c-axis contact with resistance 8.7 Ω, the features associated to the large gap are rather clear and well separated from the electron-boson structures. Here, the best fit gives $\Delta_1 = 3.78$ meV and $\Delta_2 = 7.63$ meV.

Figure 5(b) shows the low-temperature conductance curve of an ab-plane point contact on the same crystal. Here, the features related to the small gap are dominant and no clear structures associated to the large gap can be detected by eye. With the normalization shown in the figure (obtained without any shift of the normal-state conductance), the two-gap fit (red line) is superior to the single-gap one (blue line), because it can reproduce both the position of the maxima (see inset) and the width of the curve. This fit gives $\Delta_1 = 3.68$ meV and $\Delta_2 = 7.13$ meV. The gap value obtained by the single-band fit is $\Delta = 5.03$ meV and can be seen as a sort of an average of $\Delta_1$ and $\Delta_2$, as usually happens when multiple gaps are insufficiently resolved in the spectra. Actually, by choosing different normalizations, the single-band fit can become almost indistinguishable from the two-band one (in particular, if the amplitude of the maxima (see inset) and the width of the curve.

IV. 75As NMR RESULTS: ROLE OF SPIN FLUCTUATIONS

While PCARS provides detailed information on the nature and value of the superconducting gaps, NMR can be used to investigate also the normal-state properties of PrFeAs(O,F). Here, we employ mostly 75As-NMR measurements at 7 T to determine the static (line widths and shifts), as well as the dynamic (spin-lattice relaxation) electronic properties of PrFeAs(O,F). The very small size of the single crystals implied a rather poor S/N ratio. Therefore, the NMR measurements had to be performed on powder samples (obtained by crushing the available single crystals).

Since 75As has a nuclear spin $I = 3/2$ with a moderately large quadrupole moment ($Q = 31.4$ fm$^2$), the observed NMR line consists of the central Zeeman $+1/2$ to $-1/2$ transition broadened by a second-order quadrupole perturbation, while the two satellites are much too weak and far apart. Both the two-peak lineshape (see Fig. 7) and its variation with temperature (at least down to 40 K) are similar to those of the 75As NMR lines observed in lightly F-doped LaFeAsO$^{33,34}$ or in pure ThFeAsO$^{35}$.

A. NMR lineshapes and lack of magnetic order

As shown in Fig. 7, down to $T_c$ (~ 15 K at 7 T), the position of the NMR lines does not change significantly with temperature, hence suggesting the absence of a magnetic order. This is in agreement with the magnetometry results which, above $T_c$, also show a weak and almost flat response. Below $T_c$, instead, the NMR lines do not exhibit the expected drop in frequency in the superconducting
As for the NMR linewidths, these too are similar to those of analogous compounds (see references above), mostly in terms of the FWHM value (here, typically around 1.30 MHz), but less in terms of its temperature dependence. In the PrFeAs(O,F) case, too, we observe a moderate increase of FWHM with decreasing temperature (indicative of enhanced magnetic spin fluctuations). However, unlike the generic case, here the two-peak 75As central transition shows a progressive broadening of the peaks, which, at the same time, become closer, until they merge below 15 K. While these contrasting trends leave the global FWHM practically unchanged, the clear change in line shape indicates a progressive enhancement of the Pr$^{3+}$ magnetic effects at low temperature, reflected also in the $1/T_1$ relaxation rates (see below). Since the observed increase of FWHM is smooth, this is in sharp contrast with the abrupt changes expected in case of a magnetic phase transition. The lack of appreciable variations of FWHM vs. $T$ strongly suggests that PrFeAs(O,F) does not exhibit any AF order but, at the same time, it may sustain AF fluctuations, as we show in detail below. Finally, note that, in our case, the linewidth broadening may also arise from quadrupole effects, mostly reflecting disorder or defects intrinsic to doped samples. However, comparisons of pure-NQR with NMR spectra have shown that the former is of secondary importance and does not lead to the observed temperature dependence. In particular, the quadrupole broadening is quantitatively less pronounced for the central $+1/2\to-1/2$ transition, affected only to second order by the quadrupole effects.

B. NMR relaxation rates and spin fluctuations

57Fe Mössbauer spectroscopy studies on the PrFeAsO parent compound found an itinerant 3d magnetic order of Fe$^{2+}$ ions, with an onset at about 165 K, accompanied by an orthorhombic distortion of the unit cell. Upon lowering the temperature, this evolves into a complete longitudinal incommensurate spin-density-wave (SDW) order below 139 K. At much lower temperatures (12.8 K) also the localized Pr$^{3+}$ magnetic ions order. Although the critical temperature of this second magnetic system depends on the type of rare-earth, such phenomenology is common to many 1111 parent compounds. Upon F doping, the SDW magnetism is expected to vanish, yet the spin fluctuations to survive. Indeed, as we show below, in the slightly doped PrFeAs(O,F) case, the $^{75}$As NMR spin-lattice relaxation shows clear signatures of strong magnetic fluctuations.

The $^{75}$As spin-lattice relaxation times $T_1$ were evaluated from the magnetization recovery curves recorded on the leftmost peak (at ca. 50.97 MHz) at different temperatures. One of such curves, for $T = 100$ K, is shown in the inset of Fig. 8(a). The $T_1$ value for the central transition of a spin-3/2 $^{75}$As nucleus is obtained via:

$$M_{x}(t) = M_{x}^0 \left[1 - f \left(0.9 e^{-\frac{\Delta E}{kT}} + 0.1 e^{-\frac{\gamma \Delta E}{9T}}\right)\right].$$

Here $M_{x}^0$ is the magnetization value at thermal equilibrium,
Here, the particularly conspicuous increase of \( 1/T_1 \) at lower temperature is clear evidence of the increasing importance of electronic spin fluctuations at low \( T \). Such fluctuations are nothing but a residual of the original SDW magnetic order present in the undoped parent compound\(^{45}\). Indeed, it has been postulated that optimally doped samples strike a balance between the competing AF phase at lower dopings and the too weak fluctuations to sustain superconductivity at higher dopings. Our results confirm such a scenario also for the Pr-1111 case.

Let us now discuss the NMR relaxation data in more detail. As the temperature is lowered from room temperature, first we observe a gradual increase of fluctuations, peaking at a cusp-like maximum above \( T_\text{c} \), followed by an abrupt decrease at low temperatures. If we compare our results with those obtained in similar Ln1111 compounds with different dopings\(^{37,43}\), it emerges that parent – or very low doped – compounds show a diverging behavior at \( T_{\text{SDW}} \), while overdoped samples show a very weak peak at \( T_\text{c} \). Our case, closer to optimal doping (given the relatively high \( T_\text{c} \)), indicates that in PrFeAs(O,F) spin fluctuations are still dominant down to \( T_\text{c} \). They would continue growing as the suppressed SDW is shifted towards 0 K, but a change in the physics of the system close to \( T_\text{c} \) clearly changes also their behavior. While in the undoped case, such event would be the opening of a SDW gap, in our case, the possibilities seem restricted to the interving SC phase. However, since the \( 1/T_1 \) peak occurs at 25 K, i.e., 10 K above \( T_\text{c} \) (~ 15 K at 7 T), the most likely explanation for its occurrence might be given by the Bloembergen-Purcell-Pound model\(^{45}\), as observed also in other under- or optimally doped La-1111 compounds\(^{46}\). Such model describes the behavior of the spin-lattice relaxation rate, \( 1/T_1 \), under the influence of local fluctuating magnetic fields \( h(t) \) and indeed predicts a peak in \( 1/T_1 \) at the temperature where the effective correlation time of the spin fluctuations \( \tau_c \) equals the inverse of the Larmor frequency \( \omega_L \). Considering the similarity of PrFeAs(O,F) with other Ln1111 compounds, we expect the BPP model to apply also in our case. Here, the particularly sharp cusp in \( 1/T_1 \) might reflect the joint effect of the Re\(^{2+}\) and Pr\(^{3+}\) spin fluctuations.

Finally, we consider the \( 1/T_1 \) behavior below \( T_\text{c} \). Given the high quality of data, normally one could use them to study the superconducting gap and pairing. Unfortunately, close to 0 K, the \( 1/T_1 \) data converge at 0.3 s\(^{-1}\) and not at zero, as expected for a superconductor. In fact, deep in the SC phase, all electrons are bound into Cooper pairs, making the hyperfine interactions with the nuclei a very inefficient relaxation mechanism and driving the relaxation rate to zero. In our case, the finite value of \( 1/T_1(0) \) indicates that, at very low temperatures, other relaxation mechanisms are at play. This excludes a possible use of the data collected in the SC phase and explains why here we limited our NMR study to the normal phase. Among the alternative relax-
ation mechanisms one could think of disorder-related relaxation channels (intrinsic to doping). However, a comparison with La-1111 results excludes it, since compounds with widely different dopings still exhibit a $1/T^3$ behavior at low $T$. On the other hand, the presence of a magnetic ion, such as Pr$^{3+}$, could well justify our results. Indeed, an almost identical dependence of the $^{75}$As NMR relaxation rate vs. temperature is also found in the Ce-1111 case. We recall that Ce, Pr, and Nd have similar magnetic moments (free-ion values of 2.54, 3.58, and 3.62 $\mu_B$, respectively), whose strong coupling with Fe spin fluctuations in the FeAs layer could explain our results, as well as the different low-$T$ behavior of relaxation compared with the nonmagnetic La-1111 case.

V. CONCLUSIONS

By combining point-contact Andreev spectroscopy with nuclear magnetic resonance methods we investigated in detail the normal- and superconducting state properties of underdoped PrFeAs(O,F), a member of the Ln1111 family. Point-contact Andreev spectroscopy performed on single crystals provides evidence of the multiband/multigap nature of the PrFeAs(O,F) superconductivity. No indications of low-energy quasiparticles were found in the point-contact spectra, suggesting a fully gapped superconductor with no nodes. A small $\Delta_1 \approx 3.5 \pm 0.5$ meV gap was found not only to be very robust, but also to agree well with the results of microwave penetration depth and quasiparticle conductivity measurements. Additional structures, in the form of conductance maxima or shoulders, could be interpreted as due to one or possibly two larger gaps, whose amplitudes lie in the energy ranges $\approx 6.0$–$7.5$ meV (in agreement with optical conductivity measurements) and $\approx 8$–$10$ meV. The latter values are quite large and would correspond to gap ratios $2\Delta/k_BT_c$ of the order of 9. Moreover, we have shown that additional high-energy structures, ubiquitous in the conductance curves and not predicted by any BCS-based theory, are the hallmark of strong coupling between the electrons and spin fluctuations and can only be accounted for in the framework of a strong-coupling theory of superconductivity. Finally, magnetic resonance results in the normal phase provide clear evidence about the lack of any magnetic order in PrFeAs(O,F). Further, spin-lattice relaxation data suggest that this compound, similarly to other members of the 1111 family, hosts substantial electronic spin fluctuations (here enhanced by the presence of Pr$^{3+}$ ions), which are expected to mediate the superconducting pairing.

ACKNOWLEDGMENTS

G.A.U. acknowledges support from the MEPhI Academic Excellence Project (Contract no. 02.a03.21.0005). This work was partly supported by the Swiss National Science Foundation (SNSF) through Grant no. 200021-169455.

Appendix A: Basics of the PCARS technique

To ensure proper PCARS measurement conditions, the contact must be smaller than the electronic mean free path, so that the conduction is ballistic, no Joule dissipation occurs in the contact region, and the resistance of the contact largely exceeds the resistance of the normal bank. In these conditions, the voltage drop at the N/S interface practically coincides with the total potential difference between the electrodes, $V$, and the excess energy with which electrons are injected in the S side of the junction is just $eV$. Provided that there is no insulating layer at the sample surface, the conduction is dominated by the Andreev reflection, even though the probability of quasiparticle tunnelling is not zero. The raw $dI/dV$-vs-$V$ spectrum already contains qualitative information on the number, amplitude, and symmetry of the gap(s). However, a more quantitative analysis can be made by fitting it with suitable models for the Andreev reflection at the N/S interface.

To fabricate the contacts, we used the so-called “soft” technique, widely described elsewhere. In a few words, we stretch a thin Au wire over the crystal, until it touches the surface in a single point. Typically, ballistic contacts have resistances of a few tens of Ohms (even though the actual value depends on the properties of the sample, namely on its normal-state resistivity). This type of contacts can be mechanically unstable, especially during cooling/heating, because of the different thermal coefficients of the materials. Thus, in some cases, we used a drop of conducting Ag glue to improve the stability. Independent of the presence of Ag glue, the actual contact must be thought of as a parallel of nanoscopic contacts between a normal metal and a superconductor.

Appendix B: A strong-coupling model for the superconductivity in Pr-1111

As already mentioned, one can account for the presence of electron-boson structures only by using a strong-coupling extension of the BCS theory, i.e., the Eliashberg theory. Based on the similarity with other electron-doped Fe-based superconductors, we assume Pr-1111 to be described by an effective $s\pm$-wave three-band model, with one hole-like band centered at $\Gamma$ and two electron-like Fermi surface sheets at the corners of the Brillouin zone. This model is described in detail elsewhere.

To calculate the SC gaps and the critical temperature in this model, one has to solve six coupled equations for the complex order parameters $\Delta_1(i\omega_n)$ and the renormalization functions $Z_1(i\omega_n)$, where $i=1,2,3$ is the band index and $\omega_n$ are the Matsubara frequencies. The frequency (energy) dependence of the order parameters, normally ignored in the BCS theory, is here the key factor that accounts for the presence of the electron-boson coupling features. There are many input parameters, including: i) nine electron-phonon spectral functions, $a^2_i \sigma^{ph}(\Omega)$; ii) nine electron-boson (spin fluctuations) spectral functions, $\alpha^2_i F^{ij}(\Omega)$; iii) nine ele-
ments of the Coulomb pseudopotential matrix, \( \mu_{ij}^{\text{ph}}(\omega) \). To a first approximation, we neglect the disorder, thus assuming that all the scattering rates (from either magnetic or nonmagnetic impurities) are zero. To further simplify the problem, some additional assumptions can be made, shown to be valid in the case of iron pnictides.\(^{53-55}\) In particular, we know that phonons do not contribute significantly to the (dominant) interband coupling\(^{49}\) (i.e., \( \lambda_{ij}^{\text{ph}} \approx 0 \)) and that the total electron-phonon coupling constant is small.\(^{56}\) Hence, we can simply neglect the phonon contribution and assume \( \lambda_{ii}^{\text{ph}} = 0 \). Spin fluctuations, instead, are known to provide mostly the interband coupling between the hole- and electron bands, so we can assume \( \lambda_{ii}^{\text{sf}} = 0 \). Finally, following Ref.\(^{57}\), we will assume that the Coulomb pseudopotential matrix is identically zero, i.e., \( \mu_{ii}^{\text{ph}}(\omega_c) = \mu_{ij}^{\text{ph}}(\omega_c) = 0 \). With these approximations, the electron-phonon coupling constant matrix \( \lambda_{ij} \) becomes,\(^{49,53-55}\):

\[
\lambda_{ij} = \begin{pmatrix}
0 & \lambda_{12}^{\text{sf}} & \lambda_{13}^{\text{sf}} \\
\lambda_{12}^{\text{sf}} & 0 & \lambda_{23}^{\text{sf}} \\
\lambda_{13}^{\text{sf}} & \lambda_{23}^{\text{sf}} & 0
\end{pmatrix},
\]

(B1)

where \( \nu_{ij} = N_i(0)/N_e(0) \), with \( N_i(0) \) the normal density of states at the Fermi level for the \( i \)-th band. The electron-boson coupling constants are defined through the Eliashberg functions \( \alpha_{ij}^{\text{sf}} F_{ij}^{\text{sf}}(\Omega) \). Following Refs.\(^{53-55}\) we choose these functions to have a Lorentzian shape, with a maximum at the energy \( \Omega_0 \) and half-width \( \Omega_0/2 \). \( \Omega_0 \) is the characteristic energy of the mediating boson, which corresponds to the characteristic energy of the spin resonance\(^{52}\) and is related to the critical temperature by the empirical law \( \Omega_0 = 2T_c/5 \) that has been demonstrated to hold, at least approximately, for many iron pnictides.\(^{58}\)

The factors \( \nu_{ij} \) that enter the definition of \( \lambda_{ij} \) can normally be determined from the band-structure calculations, unfortunately not available for the Pr-1111 case. However, since our aim is just to show that the wide shoulders observed in our datasets (we will refer in particular to the data in Fig. 4b) are due to electron-boson coupling, we may use the values employed in Co-doped Ba-122\(^{55}\) (because this also is an electron-doped Fe-based superconductor with a similar \( T_c \)), i.e., \( \nu_{12} = 1.12 \) and \( \nu_{13} = 4.50 \). Hence, only two free parameters remain, \( \lambda_{12} \) and \( \lambda_{13} \), that need to be fixed in order to reproduce the experimental \( T_c \) and the experimental gaps. We find \( \lambda_{12} = 0.7 \) and \( \lambda_{13} = 1.8 \), giving a total coupling constant \( \lambda_1 = 2.374 \).

Owing to the \( s \pm \) symmetry, the order parameter has opposite signs on the hole-like and the electron-like Fermi surface sheets. The low-temperature values of the gaps turn out to be \( \Delta_h = 6.33 \text{ meV} \), \( \Delta^1 = 3.35 \text{ meV} \) and \( \Delta^2 = 8.56 \text{ meV} \) where the subscripts \( h \) and \( e \) refer to the hole-like and electron-like Fermi surfaces, respectively. The gap values agree very well with the gap distribution shown in Fig. 6. The calculated critical temperature is \( T_c = 28.6 \text{ K} \). This value is slightly larger than the onset of the superconducting transition (see Fig. 1). However, since the coupling is of electronic origin, there is a feedback effect of the SC condensate on the spin-fluctuation spectrum.\(^{55,52}\) Taking this effect into account, as we already did in Ref.\(^{55}\), the critical temperature turns out to be \( T_c = 22.86 \text{ K} \). Once the order parameters as a function of energy are known, they can be inserted into the equations for the Andreev reflection (i.e., in the three-band version of the BTK model) and the conductance curve can be calculated. The BTK model contains, in addition to the gap amplitudes, the relative weights of the bands, the barrier parameters and the broadening parameters. We chose the values of these parameters in order to obtain a curve similar to that in Fig. 4b. In particular, we took \( Z_h = Z_e = Z_{c2} = 0.33 \), \( T_h = 2.95 \text{ meV} \), \( R_{h1} = 1.45 \text{ meV} \), \( R_{c2} = 4.05 \text{ meV} \), \( w_h = 0.20 \), \( w_{e1} = 0.25 \), \( w_{c2} = 0.55 \). The resulting curve is shown in Fig. 9 with a blue line. Clearly, accounting for the energy dependence of the order parameters gives rise to very wide (in energy) and very high (in amplitude) shoulders that resemble very closely those of the experimental data (actually, a proper normalization could easily reproduce an experimental spectrum with the same shape), but does not affect the low-energy part of the spectrum, where the gap-related features show up. Indeed, the fit of the theoretical curve with the same two-band BTK model we used to fit the experimental data would have given again \( \Delta_1 = 3.3 \text{ meV} \) and \( \Delta_2 = 8.5 \text{ meV} \).


