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Robustness of superconductivity to structural disorder in $\text{Sr}_{0.3}(\text{NH}_2)_3(\text{NH}_3)_{1-x}\text{Fe}_2\text{Se}_2$

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The superconducting properties of a recently discovered high-$T_c$ superconductor, Sr/ammonia-intercalated FeSe, have been measured using pulsed magnetic fields down to 4.2 K and muon spin spectroscopy down to 1.5 K. This compound exhibits intrinsic disorder resulting from random stacking of the FeSe layers along the $c$ axis that is not present in other intercalates of the same family. This arises because the coordination requirements of the intercalated Sr and ammonia moieties imply that the interlayer stacking (along $c$) involves a translation of either $a/2$ or $b/2$ that locally breaks tetragonal symmetry. The result of this stacking arrangement is that the Fe ions in this compound describe a body-centered-tetragonal lattice in contrast to the primitive arrangement of Fe ions described in all other Fe-based superconductors. In pulsed magnetic fields, the upper critical field $H_c2$ was found to increase on cooling with an upward curvature that is commonly seen in type-II superconductors of a multiband nature. Fitting the data to a two-band model and extrapolation to absolute zero gave a maximum upper critical field $\mu_0H_c2(0)$ of 33(2) T. A clear superconducting transition with a diamagnetic shift was also observed in transverse-field muon measurements at $T_c \approx 36.3(2)$ K. These results demonstrate that robust superconductivity in these intercalated FeSe systems does not rely on perfect structural coherence along the $c$ axis.

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I. INTRODUCTION

The discovery of superconductivity in iron-based materials [1] has led to a new family of systems with substantial structural variations, but all compounds are composed of FeB layers (where $B =$ As, Se, Te, P or some mixture) and the variations arise from the way these layers are assembled and which atoms are included between them [2–5]. Increasing the interlayer separation in FeSe has been found to give rise to a dramatic effect on the superconducting transition temperature $T_c$. For example, pure FeSe ($T_c = 8.5$ K [6]) can be intercalated with alkali-metal ions and ammonia [7–9] or other organic molecules [7,10] to produce new superconductors with transition temperatures of up to around 45 K. Similar high-$T_c$ behavior can be induced using metal hydroxides as the spacer layer [11–14]. This trend of increase in $T_c$ with increasing layer separation does not continue indefinitely [10], and this effect can be rationalized through first-principles calculations [15]. However, the effect of controlled structural disorder on $T_c$ has not been so closely examined. We have identified an intercalated FeSe compound in which random stacking of well-defined layers results in a paracrystalline structure. In this paper, we demonstrate that the superconducting state is nevertheless robust.

The compound $\text{Sr}_x(\text{NH}_2)_3(\text{NH}_3)_{1-x}\text{Fe}_2\text{Se}_2$ ($x = 0.3$, $0.2 \leq y \leq 0.6$) belongs to a family of layered intercalates $\text{A}_x(\text{NH}_2)_3(\text{NH}_3)_{1-y}\text{Fe}_2\text{Se}_2$ ($A = \text{Li, Na, K, Rb, Cs, Ca, Sr, Ba, Eu, and Yb}$). In these materials, bulk superconductivity occurs in anti-PbO-type FeSe layers composed of edge-sharing FeSe4 tetrahedra that are separated by metal ions, amide ions, and ammonia molecules [8,9,16,17]. Figure 1 shows the various structures and the conventional unit cell with $a = b \approx 3.8$ Å and $c = 16.5–17.4$ Å. These compounds remain tetragonal down to low temperatures and do not exhibit an orthorhombic distortion. Neutron scattering and x-ray-diffraction measurements have revealed some structural differences related to the size of the cation used; when it is small ($A = \text{Li}$), Fe ions in adjacent layers occupy the same primitive tetragonal sublattice [8], as shown in Fig. 1(a). For larger cations ($A = \text{K, Rb}$), the arrangement of Fe and Se ions is unchanged, but the larger electronegative cations share the same sites as the amide and ammonia moieties in the body-centered position of the primitive FeSe sublattice [see Fig. 1(b)]. As we describe in detail in Sec. III, the Sr case presents a scenario in which the coordination requirements of the Sr cation and the amide and ammonia moieties are best satisfied by an arrangement of adjacent FeSe layers that results in a random stacking of these layers along the $c$ axis, which may be described as paracrystalline. We will demonstrate through SQUID magnetometry, pulsed magnetic fields, and muon spin rotation ($\mu$SR) measurements that, despite this unusual structural disorder, superconductivity remains robust with a $T_c \approx 36$ K that is significantly higher than that of a parent compound.

II. SYNTHESIS

In an argon-filled glove box, finely ground FeSe powder (1.970 g, 14.611 mmol) synthesized from the elements as described elsewhere [8] and pieces of strontium metal (0.640 g,
The crystal structures of the products were analyzed using synchrotron X-ray powder diffraction on beamline I11 at the Diamond Light Source, or beamline ID31 at the ESRF. Samples were sealed under an argon atmosphere within 0.5-mm-diam borosilicate glass capillaries. The diffractograms revealed a series of extremely sharp reflections that could all be indexed on a body-centered-tetragonal unit cell with lattice parameters $a_{\text{sub}} = 2.7$ Å and $c = 17.4$ Å [see Fig. 2(a)]. The intensities of these reflections were accounted for by the model shown in Figs. 2(b) and 2(c) in which the familiar square planar Fe nets found in all the iron selenide superconductors with Fe-Fe = $a_{\text{sub}} \sim 2.7$ Å are related by the body-centering translation.

Since in this model the basal lattice parameter $a_{\text{sub}}$ was equal to the Fe-Fe distance, the selenide ions were inevitably modeled as disordered over two sites located above and below the centers of each square of Fe ions. Further sites were located in the interlamellar space, and they were occupied by N atoms from ammonia or amide moieties (N:Fe ratio 1:2) and Sr ions (Sr:Fe ratio 0.3:2). This model was also obtained using a charge flipping algorithm implemented within TOPAS ACADEMIC [18], and it was consistent for all the samples investigated. The interatomic distances and the coordination environments using this model were chemically realistic providing that local ordering of occupied and unoccupied Se, N, and Sr sites was imposed. This required an $a = b = \sqrt{2} a_{\text{sub}}$ basal expansion of the tetragonal unit cell to achieve a chemically realistic Fe-Se layer and stacking disorder along $c$ to account for the apparent smaller cell.

Weak structured diffuse scattering was evident in the diffractograms as a result of the stacking disorder. This diffuse scattering was accounted for in a semiquantitative manner by constructing a superstructure in which layers were stacked along the $c$ axis in a way that respected the coordination environments for the intercalate species shown in Figs. 1(a) and 1(b). NH$_3$ and NH$_2^+$ moieties were sixfold-coordinated by selenide ions [a square of four in one layer and a pair in the layer above, producing an isosceles triangular prism; see Fig. 2(g)] such that N-H···Se distances were 3.7 Å, similar to those found in the analogous Li/NH$_3$ intercalates. Sr$^{2+}$ ions were coordinated by a triangle of selenide ions about 3.2 Å apart and by N atoms from the amide or ammonia species.
When using a model in which chemically realistic layers [boxed region in diagram (e)] with basal lattice parameters $a_{\text{sub}} / b_{\text{sub}} \times c_{\text{sub}}$ are related by the body-centered-tetragonal sublattice. Figures 1(c)–1(f) demonstrate how the nature of these translations, represented by dotted lines, results in a paracrystalline structure as follows. Consider starting from an identical base layer (the region below the first dotted line), in which the Fe, Se, and intercalated species occupy the same crystallographic sites as those found in the Li case. To construct the next layer, each atom must be translated by either $a/2$ or $b/2$. In each case, the Fe sites are the same but the positions of the Se, ammonia, and metal species change depending on the chosen direction. As the translation occurs along a randomly chosen direction with the addition of each layer, it follows that for $n$ layers there would be $2^n$ possible configurations. Our interpretation of the structural data is that despite the random stacking, the well-separated FeSe layers themselves remain well ordered, and this helps to explain our findings, described below, that robust superconductivity is preserved in these systems.

IV. SQUID MAGNETOMETRY

Magnetic susceptibility measurements were recorded using a Quantum Design MPMS-XL superconducting quantum interference device (SQUID) magnetometer. Samples of the Sr/NH$_3$ intercalate (30.2 and 29.8 mg) and the Sr/NH$_3$ intercalate (34.0 and 30.2 mg) were filled and immobilized in gelatin capsules. Measurements were conducted in dc fields of 5 mT in the temperature range 2–150 K after cooling in zero applied field (ZFC) and in the measuring field (FC).

The magnetometry measurements performed directly after the synthesis (Fig. 3) indicate bulk superconductivity with a noticeable sharp drop to negative susceptibility values at $T_c = 36(1)$ K. There is no effect of H/D substitution on $T_c$. Although the superconducting volume fraction is larger in the case of D, we believe this is consistent with the natural variability of volume fractions achieved in different syntheses.
Even after nine months, the superconducting volume fractions and the sharp drop at $T_c$ were found to be unchanged. A small reduction in $T_c$ (of 2–3 K) was detected, although this was not found to be correlated with any significant change in structure. For the remainder of the paper, we will present magnetometry and $\mu$SR data on the Sr/NH$_2$ intercalate.

V. PULSED-FIELD MAGNETOMETRY

Powder samples were measured at the Nicholas Kurti Magnetic Field Laboratory, Oxford using a proximity detector oscillator (PDO) dynamic susceptometer [20,21]. The sample is mounted in an ampoule under an argon atmosphere and placed inside a small sensor coil that is inductively coupled to the PDO circuit. This is essentially an LCR circuit with a resonant frequency $f$ that is measured as a function of field and temperature. In metals (superconductors) $f$ is highly dependent on changes in the skin (penetration depth), and in insulators the signal is dominated by the magnetic permeability. Thus when sweeping in field or temperature, the superconducting phase transition manifests as a large change of $f$ at $H_{c2}$ or $T_c$, which corresponds to the difference in skin and penetration depth of the two states.

Data were taken for temperatures in the range 4.2–40 K in fields of up to 38 T. Below $T_c$, the onset of superconductivity is marked by a sudden rise in frequency and a deviation from the normal-state signal [see Fig. 4(a)]. Additional heating can occur due to eddy currents that are generated in the mixed or normal state by the applied field. In some cases, this produces hysteresis in the data and a shift in $H_{c2}$, with the effect being more pronounced with a larger $dB/dt$ (a shorter pulse length for the same maximum field). We find our samples to be insensitive to these effects and observe no difference in $f$ at varying pulse lengths between 7 and 14 ms [see Fig. 4(c) for field profiles of the pulses].

We note that due to the broadness of the phase transition, the value of $H_{c2}$ is particularly difficult to determine at low $T$. This may be due to the critical field anisotropy between the in-plane and out-of-plane directions increasing as the system is cooled, and so for a powdered sample there is a more gradual change in frequency. Thus, we define two different methods for extracting $H_{c2}$. First, we take $H_{c2}$ as the peak in $d^2f/dH^2$ and its uncertainty as the half-width at half-maximum [see Fig. 4(d)]. Second, one may perform a linear extrapolation of the curve on either side of the transition and define $H_{c2}$ as the point of intersection of these two slopes. In this case, another criterion is the point at which the line extrapolated from the curve below the critical field crosses the normal-state signal. We take the uncertainty as the difference between these two points of intersection [see Fig. 4(b)]. We find that both methods give similar results, with $H_{c2}$ increasing steadily as the sample is cooled (see Fig. 8). We note that the curve exhibits a concave form that has been observed in multiband high-$T_c$ superconductors [22].

VI. TRANSVERSE FIELD $\mu$SR MEASUREMENTS

To probe the internal field distribution in the vortex state, transverse field $\mu$SR ($\mu$SR) measurements were carried out using the GPS instrument at the Swiss Muon Source (PSI), Switzerland and the MuSR spectrometer at ISIS, UK. In these measurements, spin-polarized muons are implanted into the material with an external field $B_{app}$ applied perpendicular to the initial muon spin direction. They will then rotate at the Larmor frequency $\omega = \gamma_B B_{loc}$, where $\gamma_B/2\pi = 135.5$ MHz T$^{-1}$ is the muon gyromagnetic ratio and $B_{loc}$ is the local field, which may include contributions from $B_{app}$, magnetic ions, and nuclear dipoles. This is measured directly by the time-dependent positron decay asymmetry $A(t)$ [23].

Muons implanted in a type-II superconductor will settle in particular crystallographic sites and experience a magnetic field due to the applied field as well as any internal variation resulting from the formation of a vortex lattice. However, as the vortex lattice is, in general, incommensurate with the crystal lattice, the full range of the field distribution within the vortex lattice is sampled, although it is broadened slightly by the field from nuclear spins. In TF measurements, the measured asymmetry is proportional to the spin polarization as measured by individual detector banks. This experimental situation can be modeled using a fit function,

$$A(t) = A_{bg} \cos(\gamma B_{app} t + \phi) \exp[-(\sigma_{bg} t)^2/2] + A_{VL} \cos(\gamma B_{pk} t + \phi) \exp[-(\sigma_{VL} t)^2/2],$$

(1)
where the phase $\phi$ results from the detector geometry, and $A_{VL}$ and $A_{bg}$ are the relaxing asymmetry due to vortex lattice and background contributions, respectively (the latter originates from muons stopping in the nonsuperconducting fraction of the sample or silver sample holder). This model assumes a symmetric Gaussian distribution of local fields with a peak value $B_{pk}$ ($B_{app}$) and standard deviation, or damping factor, $\sigma_{VL}$ ($\sigma_{bg}$) corresponding to the vortex lattice (background) contributions. The width of the field distribution is given by $B_{rms} = \sigma_{VL}/\gamma_\mu$ and its temperature dependence is shown in Fig. 5. We observe a broadening of the field distribution (increase in $B_{rms}$) when the sample is cooled through $T_c$ in an applied field of 10 mT that can be attributed to the formation of the vortex lattice with superconductivity setting in at around 36.3(2) K. As the material enters the superconducting state, the applied field is partially screened causing a diamagnetic shift in $B_{pk}$ for $T < T_c$. For powdered samples of anisotropic superconductors such as these, the in-plane penetration depth $\lambda_{ab}$ is related to the field width via

$$B_{0} = \frac{\Phi_0}{(3\pi\lambda_{ab})^2},$$

where $\Phi_0$ is a magnetic flux quantum and $B_0$ is the vortex lattice contribution to $B_{rms}$ [24]. For these data, it is assumed that the only other contribution to $B_{rms}$ is from nuclear dipole fields, which are temperature-independent and add in quadrature. We fit our data to the phenomenological functions

$$B_{rms}(T) = \left[B_0^2 \left(1 - \frac{T}{T_c}\right)^{\alpha - \beta} + B_{dip}^2\right]^{\frac{1}{2}},$$

$$B_{pk}(T) = B_{app} - B_{dia} \left[1 - \left(\frac{T}{T_c}\right)^{\alpha'}\right],$$

where $B_{dia}$ is the maximum diamagnetic shift of the peak field, and $B_0$ and $B_{dip}$ are the widths corresponding to the vortex lattice (at $T = 0$) and nuclear dipole contributions, respectively. For the fits shown in Fig. 5, the fitted parameters are $\beta = 0.33(3)$ and $\beta' = 0.45(3)$ with fixed $\alpha = \alpha' = 1$. Using a weighted average between the extracted value of $B_0$ of these data and additional measurements on a second batch of sample, we calculate a penetration depth of $\lambda_{ab}(T = 0) = 292(3)$ nm.

VII. ZERO-FIELD AND LONGITUDINAL-FIELD $\mu$SR MEASUREMENTS

Zero-field measurements (ZF $\mu$SR) were used to further probe the intrinsic magnetism of the system. No spontaneous oscillations were observed in the forward-backward asymmetry across the whole temperature range, nor was there any discontinuous change in amplitudes or recovery of the baseline...
asymmetry at low temperatures [see Fig. 6(a)]. Together these make the presence of any long-range magnetic order unlikely. Given that the spectra do not follow a Kubo-Toyabe relaxation, it is also unlikely that the relaxation is caused solely by nuclear moments. The lack of recovery of the baseline asymmetry at late times suggests dynamic fluctuations, and we therefore attribute the ZF signal to relaxation caused by disordered, fluctuating electronic moments. The data were fitted to the function

$$A(t) = A_{\text{rel}} \exp(-\lambda t) + A_{\text{base}},$$

where $A_{\text{rel}}$ is the relaxing asymmetry with relaxation rate $\lambda$, and $A_{\text{base}}$ is a nonrelaxing background. $A_{\text{rel}}$ can be considered as the lower bound for the superconducting volume fraction, which for this sample was around 1/3. Exponential relaxation often describes dense spin systems that are dynamically fluctuating, where the relaxation rate $\lambda$ is proportional to the variance of the local magnetic field distribution and also to the fluctuation time, or a dilute spin systems in the presence of dynamic fluctuations. Upon cooling, the relaxation rate $\lambda$ [Fig. 6(b)] is seen to increase slowly with decreasing temperature. However, below superconducting $T_c$ it increases far more rapidly, suggesting that the relaxation rate is in some way coupled to the superconducting order parameter.

Longitudinal field measurements (LF $\mu$SR) were made at 5 and 100 K in which a field was applied in the initial direction of the muon polarization (see Fig. 7). For data measured at 5 K [Fig. 7(a)], the sample was cooled below $T_c$ in zero applied field. At both temperatures, the relaxation is decoupled at relatively low fields ($\approx 5$ mT at 100 K and $\approx 20$ mT at 5 K), suggesting that the relaxation is due to fairly dilute and/or static spins with residual dynamic fluctuations. One scenario is that there is a sizable contribution to the ZF relaxation from static nuclear moments, with dilute, fluctuating electronic moments providing additional relaxation. The lack of any peak in $\lambda$ suggests the absence of any freezing of the dynamics of the moments (as one would expect in a spin glass due to magnetic interactions), supporting the interpretation that these are dilute and not strongly interacting with each other. The fact that the relaxation rate is coupled to the superconducting order parameter implies that the moments are embedded in the superconductor, rather than in phase-separated pockets. At high temperatures compared to $T_c$, the moment size and/or fluctuation rate increase slowly upon cooling. Upon cooling below $T_c$, the muon spins likely experience a combination of slower fluctuation times, larger moments, and an increased width of the field distribution. The latter could arise in a manner analogous to the increase in $B_{\text{rms}}$ in the TF measurements. Specifically, if the moments are locally in a normal (rather than superconducting) region and sufficiently numerous that there is a degree of overlap of their magnetic fields in the superconducting regions, then the decreasing penetration depth upon cooling would lead muons, which decorate all of the sample, to experience a broader distribution of local fields. Taken together, the ZF measurements therefore suggest the presence of a small concentration of dilute magnetic moments dispersed in the superconducting volume, but whose presence does not seem to adversely affect the presence of superconductivity.
The model parameters are the band diffusivities unlike the WHH model, it does not account for spin-orbit from nonmagnetic impurities, orbital pair breaking, strong band BCS superconductor \[29,30\]. This includes scattering developed by Gurevich that is based on a weakly coupled two-band \[28\] in the dirty limit did not converge. Instead, we use a model Hohenberg (WHH) model \[28\] for a one-band superconductor impurity scattering does not affect \(\rho\) and weakly dependent on the coupling constants, which contribute significantly only in the mK temperature range. Consequently, fits to the data converge only when one set of coupling constants (either interband or intraband) are fixed. Fits are shown in Fig. 9 for various values of \(\lambda_{11}, \lambda_{22}, \) and \(w = \lambda_{11}\lambda_{22} - \lambda_{12}\lambda_{21}\). The Gurevich model was successful in modeling the upward curvature of \(H_{c2}\), which stems from the difference in diffusivities of the two bands, i.e., \(d_{2} < 1\) \[30\]. However, we found that it was possible to fit for both a strong (when \(w < 0\)) and weak (\(w > 0\)) interband coupling, and fitted parameters varied considerably depending on the values of the chosen coupling constants and the method for finding \(H_{c2}\). This is particularly noticeable when extrapolating \(H_{c2}(T)\) to absolute zero. When using \(H_{c2}\) values extracted by the slope method as described earlier, we find \(\mu_{0}H_{c2}(0) \approx 33(2)\) T to be fairly consistent across various different scenarios (see Fig. 9). The derivative method for calculating \(H_{c2}\) could not be extended to the lowest-temperature datum and hence it produced a poorly constrained estimate, thus this method was deemed less reliable.

The Uemura relation \[31\] is a scaling relation between \(T_{c}\) and the superfluid stiffness \(\rho_{s}\), which seems to hold well for many exotic superconductors \[32\]. It is known to break down for overdoped cuprates \[33\] and may be an oversimplification \[34\], and other scaling relation behaviors have been explored \[35–38\]. Nevertheless, the Uemura plot of \(T_{c}\) against \(\rho_{s}\) provides a convenient means of exploring the energy scale to break up pairs as a function of the strength of the order parameter, and our TF \(\mu\)SR data allow us to extract an estimate of \(\rho_{s} = c^2/\lambda_{ab}\). As shown in Fig. 10, we find that the Sr intercalated compound is close to the main scaling line on the Uemura plot, as is the Li intercalated material \[8\], and this behavior correlates with underdoped cuprates and many other iron-based superconductors. Note that there is also another, lower line in this plot, which is common to electron-doped cuprates \[39,40\] and LiFe\(_{1-x}\)Co\(_{x}\)As \[41,42\]. For those materials, it is found that although the superconducting state is reasonably robust (the superfluid is stiff), the strength of the pairing is significantly suppressed, but these considerations do not seem to apply to our intercalated compounds.

VIII. DISCUSSION

Pulsed-field measurements indicate an upward curvature in \(H_{c2}\) versus temperature (see Fig. 8), which may be indicative of a multiband nature. This type of behavior has been observed in other layered two-band superconductors such as the FeAs-based systems \[25,26\] and the parent compound FeSe using resistivity measurements \[27\]. Fits of the Werthamer-Helfand-Hohenberg (WHH) model \[28\] for a one-band superconductor in the dirty limit did not converge. Instead, we use a model developed by Gurevich that is based on a weakly coupled two-band BCS superconductor \[29,30\]. This includes scattering from nonmagnetic impurities, orbital pair breaking, strong electron-phonon coupling, and spin paramagnetism (note that unlike the WHH model, it does not account for spin-orbit effects). The model parameters are the band diffusivities \(D_{1}\) and \(D_{2}\) and the intraband \(\lambda_{11}\) and interband \(\lambda_{12}\) and \(\lambda_{21}\) coupling constants. Assuming that the nonmagnetic impurity scattering does not affect \(T_{c}\), the upper critical field is given in reduced natural units \(h = H_{c2}D_{1}/2\phi_{0}T\) and \(t = T/T_{c}\) by

\[
a_{0}\ln t + U(h)[\ln t + U(\eta h)] + a_{2}[\ln t + U(\eta h)] + a_{1}\ln t + U(h) = 0,
\]

where

\[
U(x) = \psi \left(1 + x + \frac{\mu H}{2\pi T}\right) - \psi(x),
\]

in which \(\psi(x)\) is the digamma function, \(\eta = D_{2}/D_{1}\) is the ratio of band diffusivities, and the constants \(a_{0}, a_{1},\) and \(a_{2}\) are functions of the coupling constants.

In any case, the model is difficult to fit due to overparametrization; at high temperatures, \(H_{c2}\) is strongly dependent on \(D_{2}/D_{1}\) and weakly dependent on the coupling constants, which contribute significantly only in the mK temperature range. Consequently, fits to the data converge only when one set of coupling constants (either interband or intraband) are fixed. Fits are shown in Fig. 9 for various values of \(\lambda_{11}, \lambda_{22},\) and \(w = \lambda_{11}\lambda_{22} - \lambda_{12}\lambda_{21}\). The Gurevich model was successful in modeling the upward curvature of \(H_{c2}\), which stems from the difference in diffusivities of the two bands, i.e., \(d_{2} < 1\) \[30\]. However, we found that it was possible to fit for both a strong (when \(w < 0\)) and weak (\(w > 0\)) interband coupling, and fitted parameters varied considerably depending on the values of the chosen coupling constants and the method for finding \(H_{c2}\). This is particularly noticeable when extrapolating \(H_{c2}(T)\) to absolute zero. When using \(H_{c2}\) values extracted by the slope method as described earlier, we find \(\mu_{0}H_{c2}(0) \approx 33(2)\) T to be fairly consistent across various different scenarios (see Fig. 9). The derivative method for calculating \(H_{c2}\) could not be extended to the lowest-temperature datum and hence it produced a poorly constrained estimate, thus this method was deemed less reliable.

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IX. CONCLUSION

We have probed the superconducting properties of a recently discovered superconductor \( \text{Sr}_{0.3}(\text{NH}_3)_2(\text{NH}_3)_1-x \text{Fe}_2\text{Se}_2 \) using pulsed-field magnetometry and \( \mu \text{SR} \) techniques. The upper critical field was shown to increase upon cooling and exhibited a concave form that is reminiscent of other multigap high-temperature superconductors, which when extrapolated to absolute zero give a maximum upper critical field of \( \mu_0 H_{\text{c2}}(0) \approx 33(2) \) T. TF \( \mu \text{SR} \) measurements show a clear diamagnetic shift and a broadening of the field width that is highly reproducible between different batches of sample, with the onset of superconductivity at 36.3(2) K. ZF \( \mu \text{SR} \) did not reveal any long-range magnetic order but dilute electronic moments with some residual dynamics and whose behavior is coupled to the superconducting order parameter. We find that intercalation of Sr atoms together with amide and ammonia introduces intrinsic stacking disorder that results in a paracrystalline state. The system nevertheless retains complete structural order of the Fe sublattice, and thus these results demonstrate that robust superconductivity does not rely on perfect structural coherence along the \( c \) axis.

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