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Nodeless multigap superconductivity in organic ion intercalated (tetrabutyl ammonium)_{0.3}FeSe

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We probe the superconducting order parameter of the organic-ion-intercalated FeSe-based superconductor (tetrabutyl ammonium)_{0.3}FeSe [(TBA)_{0.3}FeSe] using muon-spin relaxation/rotation (μ SR). Zero-field μ SR measurements show only a weak temperature dependence with no evidence for magnetic ordering or broken time-reversal symmetry in the superconducting state. The temperature dependence of the superfluid density is deduced from transverse-field μ SR measurements with fields applied both parallel and perpendicular to the c axis, and can be well described by a nodeless two-gap $s + s$ model. These properties are reminiscent of those of (Li_{1-x}Fe_x)OHFe_{1-y}Se, which also has a comparably enhanced T_c , suggesting that such a gap structure is a common feature of quasi-two-dimensional intercalated FeSe-based superconductors.

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I. INTRODUCTION

The discovery of superconductivity in LaFeAsO_{1-x}F_x with $T_c = 26$ K [1] revealed the iron-pnictides to be the second family of high-temperature superconductors. Shortly after, superconductivity was found in iron-chalcogenide systems, the simplest example being FeSe with $T_c = 8$ K [2]. While an enhanced T_c can be realized in FeSe by applying pressure [3–5] or by producing it in monolayer films (with a T_c above 100 K) [6–8], the enhancement of T_c in bulk systems can also be achieved by the intercalation of spacer layers between the layers of FeSe. The first such examples were AFe₂Se₂ (A = Na, K, Rb, Cs, and Tl) with T_c up to 32 K [9,10], but studies of their intrinsic superconducting properties were hindered by inhomogeneities due to iron vacancies, leading to a phase separation between antiferromagnetic and superconducting regions [11–13]. Subsequently, higher transition temperatures were achieved via the intercalation of organic ions in Li_{0.56}(NH₂)_{0.53}(NH₃)_{1.19}Fe₂Se₂ with $T_c = 39$ K [14] and Li_{0.6}(NH₂)_{0.2}(NH₃)_{0.8}Fe₂Se₂ with $T_c = 44$ K [15]. In LiOH-intercalated (Li_{1-x}Fe_x)OHFe_{1-y}Se, bulk superconductivity with T_c over 40 K coexists with antiferromagnetic order [16,17], and its stability in air, together with a lack of phase separation, made it a promising candidate for examining the intrinsic properties of intercalated FeSe superconductors.

There has been considerable debate as to whether iron-chalcogenide superconductors have an analogous pairing state to the sign-changing s_{\pm} state generally attributed to the iron pnictides [18–20]. In bulk FeSe, some thermodynamic experiments suggest a nodal gap [21–23], however other measurements support a fully gapped behavior, but with a significant gap anisotropy and possible sample dependence [24–28]. The nature of the pairing state has been of particular interest for single-layer FeSe/SrTiO₃ and intercalated FeSe superconductors, since their Fermi surface lacks the hole pocket at the zone center present in bulk FeSe and iron pnictide systems [29–32]. In single-layer FeSe/SrTiO₃, fully gapped superconductivity was revealed [6,33], with the largest gap being of the order of 20 meV. Meanwhile two-gap nodeless superconductivity was found in (Li_{1-x}Fe_x)OHFe_{1-y}Se from various probes [34–37], but there are conflicting conclusions from quasiparticle interference measurements as to whether these correspond to a sign-changing superconducting pairing state [38], or sign preserving s -wave superconductivity [35]. Consequently, it is important to probe the pairing states of other families of intercalated FeSe-based superconductors.

Single crystals of FeSe-based superconductors intercalated with tetrabutyl ammonium [(TBA)_{0.3}FeSe] [39] and cetyltrimethyl ammonium [(CTA)_{0.3}FeSe] [40] ions have also been successfully synthesized, with high onset T_c values of 50 and 45 K, respectively. Such an intercalation leads to both a charge transfer from the sizable organic ions to the FeSe layers, and a greatly enhanced separation between FeSe layers of around 15.5 Å. The latter gives rise to a two-dimensional

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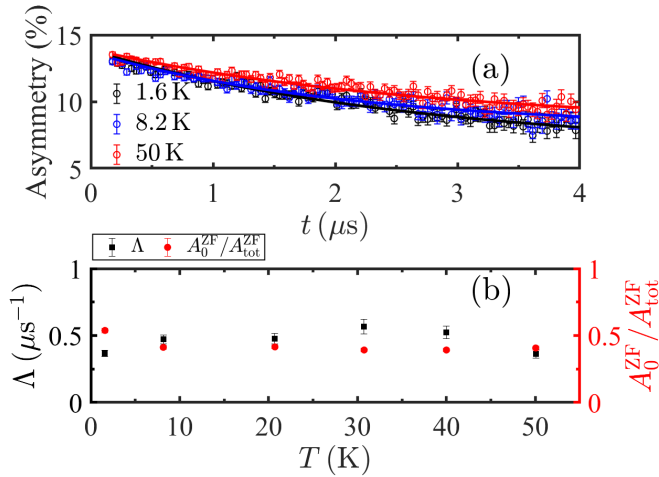


FIG. 1. (a) Zero-field μ SR spectra of $(\text{TBA})_{0.3}\text{FeSe}$ at 1.6 (black), 8.2 (blue), and 50 K (red), with the solid lines showing the fitting results using Eq. (1). (b) Temperature dependence of the fitted Lorentzian relaxation rate Λ and the ratio $A_0^{\text{ZF}}/A_{\text{tot}}^{\text{ZF}}$.

superconductivity, which is evidenced by the observation of a pseudogap phase above T_c in nuclear magnetic resonance (NMR) and in Nernst effect measurements [41], which suggest the presence of preformed Cooper pairs that lack long range phase coherence. However, detailed information about the superconducting gap structure of these superconductors is still lacking.

In this paper, we examine the superconducting pairing state of single crystalline $(\text{TBA})_{0.3}\text{FeSe}$ using zero-field (ZF) and transverse-field (TF) muon-spin relaxation/rotation (μ SR) measurements. From ZF- μ SR measurements, there is no evidence of time-reversal symmetry breaking in the superconducting state, and no signatures of magnetic ordering are observed. The temperature dependence of the magnetic penetration depth is deduced from TF- μ SR for $H \parallel c$ and $H \perp c$, which can be well described by a nodeless two-gap $s + s$ -wave model along both directions.

II. EXPERIMENTAL DETAILS

Single crystals of $(\text{TBA})_{0.3}\text{FeSe}$ were prepared via an electrochemical intercalation method, as described in Ref. [39]. ZF- and TF- μ SR experiments were performed using the general-purpose surface-muon (GPS) instrument at the Swiss Muon Source ($S\mu S$) of the Paul Scherrer Institut (PSI), Switzerland. Spin-polarized positive muons are implanted into the sample, and decay positrons are preferentially emitted along the muon-spin direction. Therefore, the asymmetry of the emitted positrons depends on the dynamics and distribution of the local magnetic fields at the muon stopping site(s). Platelike single crystals were mounted on a silver sample holder so that the c axis was perpendicular to the plate, and parallel to the muon beam. Figure 2(a) shows the configuration used for the TF- μ SR measurements. These were performed with an applied field of 80 mT along the c axis, for which the spin-polarization direction of the muon beam was rotated by 45° , and an applied field of 10 mT in the ab plane, for which the muon spins were not rotated. Note that for the spin rotated case, the maximum asymmetry is approximately

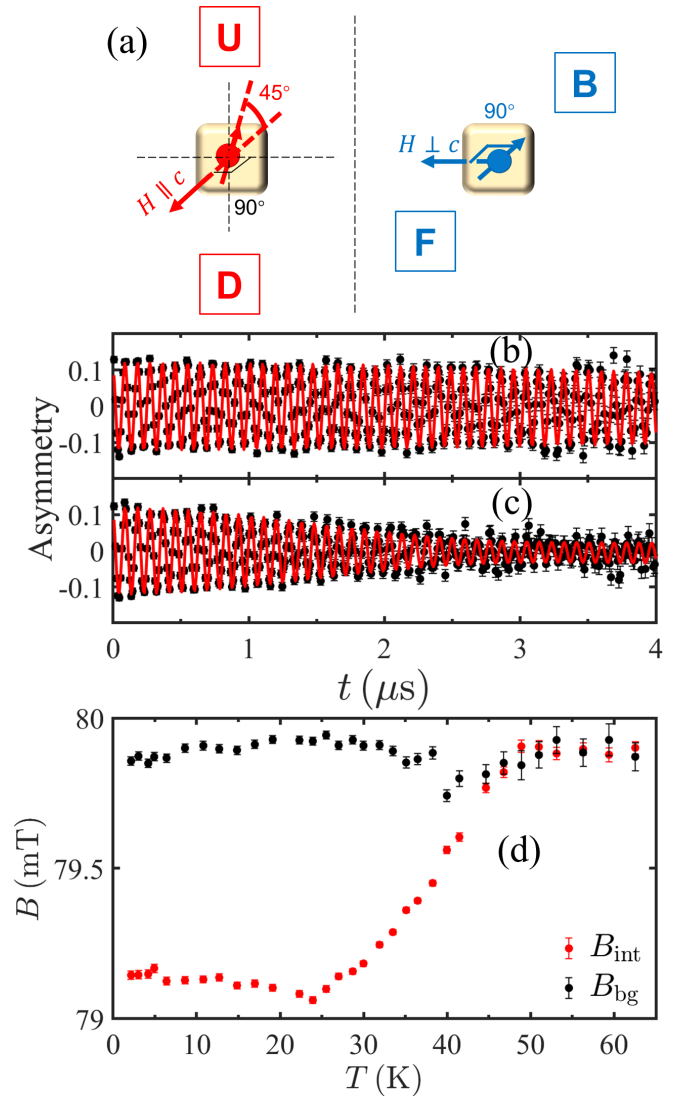


FIG. 2. (a) Illustration of the geometry for TF- μ SR measurements, where for $H \parallel c$ the muon spins are rotated upwards by 45° , and the up(U)-down(D) detector pair is used to determine the asymmetry, while for $H \perp c$ the muon spins are unrotated and the asymmetry is obtained from the forward(F)-backward (B) pair. Transverse-field μ SR spectra are displayed, collected in a transverse field of 80 mT applied parallel to the c axis at (b) 62.5 and (c) 2.2 K. (d) Temperature dependence of the internal magnetic field inside the sample B_{int} and the silver holder (background) B_{bg} as obtained from the fits.

70% of the unrotated muon spin case [i.e., $A(0) \sin(45^\circ)$]. The TF was applied above the superconducting transition temperature T_c , after which the samples were cooled to the base temperature of 1.6 K, and measurements were performed upon warming up the sample. All the data were analyzed using the data analysis program WIMDA [42].

III. RESULTS

A. Zero-field μ SR

In order to probe the magnetic properties and examine whether there is time-reversal symmetry breaking in the

superconducting state, ZF- μ SR measurements were performed on $(\text{TBA})_{0.3}\text{FeSe}$. Figure 1(a) displays ZF- μ SR spectra collected at three temperatures, 1.6, 8.2, and 50 K (below and above $T_c = 48$ K), with the initial spin-polarization of the muon beam rotated by 45° . Since the asymmetry decays relatively quickly at low times, the data were analyzed by an exponential decay function,

$$A(t) = A_0^{\text{ZF}} \exp(-\Lambda t) + A_{\text{bg}}^{\text{ZF}}. \quad (1)$$

Here A_0^{ZF} is the initial asymmetry, corresponding to muons stopping in the sample, $A_{\text{bg}}^{\text{ZF}}$ is a time-independent contribution from muons stopping in the silver sample holder, and Λ is the Lorentzian relaxation rate. Such an exponential decay can correspond to either magnetic fields fluctuating much faster than the muon precession rate, or to dilute static fields with a Lorentzian distribution. The temperature dependence of $\Lambda(T)$ and the ratio $A_0^{\text{ZF}}/A_{\text{tot}}^{\text{ZF}}$ are shown in Fig. 1(b), where $A_{\text{tot}}^{\text{ZF}}$ is the total initial asymmetry. $\Lambda(T)$ exhibits a relatively weak temperature dependence without a pronounced increase upon crossing T_c . Clearly, there is no evidence for either time-reversal symmetry breaking or significant slowing down of spin-fluctuations. In addition, the data points between 20 and 40 K were measured with the spin-polarization of the muon beam parallel to the c axis (unrotated). The very similar behaviors to the measurements with the muon spins rotated also suggests a lack of significant anisotropy of the local fields.

B. Transverse-field μ SR with $H \parallel c$

In order to probe the superconducting gap structure of $(\text{TBA})_{0.3}\text{FeSe}$, TF- μ SR measurements were performed with a TF applied both parallel and perpendicular to the c axis, with the configurations shown in Fig. 2(a). These were performed upon field-cooling the sample below T_c , in order to induce a regular flux line, from which the magnetic penetration depth $\lambda(T)$ can be extracted.

Figures 2(b) and 2(c) display the μ SR time spectra with an applied TF of $\mu_0 H = 80$ mT parallel to the c axis above (62.5 K) and below (2.2 K) T_c , respectively. At 62.5 K, $(\text{TBA})_{0.3}\text{FeSe}$ is in the normal state and the muons precess at a single frequency with a small depolarization due to quasistatic nuclear moments. While at 2.2 K, the more rapid damping arises from the nonuniform field distribution due to the flux line lattice. The data were analyzed using

$$A(t) = A_0 e^{-(\sigma t)^2/2} \cos(\gamma_\mu B_{\text{int}} t + \phi) + A_{\text{bg}} \cos(\gamma_\mu B_{\text{bg}} t + \phi), \quad (2)$$

where σ is the Gaussian relaxation rate corresponding to the sample, A_0 and A_{bg} are the initial asymmetries corresponding to the sample and background components, B_{int} and B_{bg} are the corresponding internal and background magnetic fields, ϕ is the initial phase of the precession signal, and $\gamma_\mu/2\pi = 135.5$ MHz/T is the muon gyromagnetic ratio. The total initial asymmetry $A_{\text{tot}} = A_0 + A_{\text{bg}}$ and ϕ were fixed from fitting the data at the highest temperatures (above T_c) while A_0/A_{bg} was fixed to the value obtained from fitting the data at the lowest temperature, where the background is better visible at long times. It can be seen from the red solid lines in Fig. 2 that Eq. (2) well describes the TF data above and below T_c . The temperature dependence of B_{int} and B_{bg} is shown in Fig. 2(d).

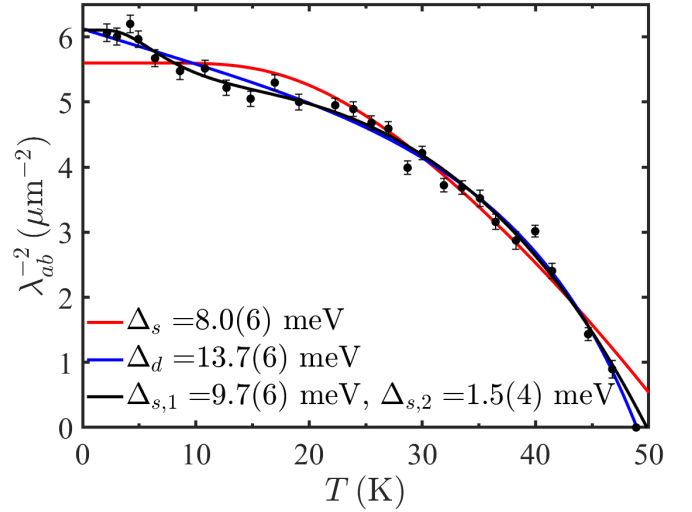


FIG. 3. Temperature dependence of $\lambda_{ab}^{-2}(T)$ of $(\text{TBA})_{0.3}\text{FeSe}$ obtained from a 80 mT transverse field parallel to the c axis. The solid lines are the fitted curves corresponding to a single gap s -wave model (red), a single gap d -wave model (blue) and a two-gap $s + s$ wave model (black).

At temperatures above T_c , B_{int} and B_{bg} show similar values, slightly smaller than the applied field, indicating a weak diamagnetism that is consistent with muons stopping in the organic-ion layers. Below T_c there is a decrease of B_{int} while B_{bg} remains almost unchanged, indicating a diamagnetic response due to the formation of a flux line lattice, as expected for a type-II superconductor.

The superconducting contribution to the Gaussian relaxation rate is calculated using $\sigma_{sc}^2 = \sigma^2 - \sigma_n^2$, where σ_n^2 is the nuclear dipolar contribution, estimated from fitting the data above T_c , where $\sigma_n = 0.200(7)$ μs^{-1} for $H \parallel c$. If the applied field is much less than the upper critical field ($H \ll H_{c2}$), σ_{sc} is proportional to the inverse square of the penetration depth [43]. When the TF is applied parallel to the c -axis, $\sigma_{sc}^{\parallel c}$ probes the in-plane penetration depth λ_{ab} as [36]

$$\lambda_{ab}^{-2} = 9.32(\mu\text{m}^{-2}/\mu\text{s}^{-1}) \times \sigma_{sc}^{\parallel c}(\mu\text{s}^{-1}). \quad (3)$$

The temperature dependence of the inverse square of the penetration depth $\lambda_{ab}^{-2}(T)$ is shown in Fig. 3. At low temperatures, $\lambda_{ab}^{-2}(T)$ saturates below ~ 4.9 K, which is consistent with a nodeless gap structure, since thermal excitations are unable to deplete the superconducting condensate in a fully gapped superconductor at sufficiently low temperatures. If there are nodes in the gap, λ_{ab}^{-2} would continuously increase upon lowering the temperature due to the presence of low energy excitations. Furthermore, at intermediate temperatures, around 12 K, $\lambda_{ab}^{-2}(T)$ exhibits an inflection point, indicative of multigap superconductivity.

The normalized superfluid density can be expressed as $\tilde{n}(T) = \lambda^{-2}(T)/\lambda^{-2}(0)$ and can be analyzed utilizing the London approach [44]:

$$\tilde{n}(T) = 1 + \frac{1}{\pi} \int_0^{2\pi} \int_{\Delta(T, \phi)}^{\infty} \frac{\partial f}{\partial E} \frac{E dE d\phi}{\sqrt{E^2 - \Delta^2(T, \phi)}}. \quad (4)$$

TABLE I. Summary of the fit parameters of the superfluid density for $(\text{TBA})_{0.3}\text{FeSe}$ with $H \parallel c$ (λ_{ab}^{-2}) and $H \perp c$ ($\lambda_{ab,c}^{-2}$), including the goodness of fit χ^2 .

	Model	Gap value(s) (meV)	Gap value(s) ($k_B T_c$)	T_c (K)	$\lambda(0)$ (nm)	x	χ^2
λ_{ab}^{-2}	s -wave gap	8.0(6)	1.8(1)	52(2)	423(4)	—	6.17
	d -wave gap	13.7(6)	3.2(1)	49.0(6)	404(2)	—	2.38
	$s + s$ -wave gap	9.7(6)/1.5(4)	2.3(2)/0.3(1)	49.8(8)	405(5)	0.81(3)	2.40
$\lambda_{ab,c}^{-2}$	s -wave gap	10(2)	2.3(5)	52(5)	516(7)	—	19.1
	d -wave gap	20(3)	4.9(6)	48(1)	497(5)	—	7.35
	$s + s$ -wave gap	17(2)/2.0(3)	4.2(5)/0.49(8)	47.0(6)	488(4)	0.74(2)	2.56

Here $\lambda^{-2}(0)$ is the magnetic penetration depth at zero temperature, and $f = [1 + \exp(-E/k_B T)]^{-1}$ is the Fermi-Dirac function. The superconducting gap function can be written as $\Delta(T, \phi) = \Delta(T)g(\phi)$, where the temperature dependent part is given by $\Delta(T) = \Delta(0) \tanh\{1.82[1.018(T_c/T - 1)]^{0.51}\}$, where $\Delta(0)$ is the zero temperature gap. The angular dependence $g(\phi)$ (ϕ is the azimuthal angle) is $g(\phi) = 1$ for an s -wave gap (Δ_s) and $g(\phi) = \cos(2\phi)$ for a d -wave gap (Δ_d).

Figure 3 displays the results of fitting for the superconducting gap structure with different models using Eq. (4). Both the nodeless single gap s -wave model and the d -wave model with line nodes do not adequately describe the data, where the former exhibits a large deviation over a wide temperature range, while the latter cannot account for the inflection point nor the low temperature saturation. To account for these features, $\lambda_{ab}^{-2}(T)$ was analyzed with a phenomenological two-gap $s + s$ wave model:

$$\tilde{n}(T) = x\tilde{n}_1(T, \Delta_{s,1}) + (1-x)\tilde{n}_2(T, \Delta_{s,2}), \quad (5)$$

where $\tilde{n}_i(T, \Delta_i)$ is the superfluid density component corresponding to the gap $\Delta_{s,i}$ calculated using Eq. (4), with a weight x ($0 \leq x \leq 1$) for $i = 1$ and $1 - x$ for $i = 2$. As shown in Fig 3, the two-gap $s + s$ model can well fit $\lambda_{ab}^{-2}(T)$, with fitted parameters $\lambda_{ab}(0) = 405(5)$ nm, $\Delta_{s,1}(0) = 9.7(6)$ meV, $\Delta_{s,2}(0) = 1.5(4)$ meV, $x = 0.81(3)$, and $T_c = 49.8(8)$ K. The fitted parameters for other models are given in Table I.

C. Transverse-field μSR with $H \perp c$

Figures 4(a) and 4(b) display the μSR spectra with a TF field of 10 mT perpendicular to the c axis at 57.4 and 1.6 K, respectively. The precession at a single frequency above T_c at 57.4 K and the depolarization at 1.6 K correspond to type-II superconductivity. The red solid lines show the fitting using Eq. (2), and σ_n is estimated to be $0.195(5) \mu\text{s}^{-1}$ from the normal state analysis. The temperature evolution of B_{int} and B_{bg} shown in Fig. 4(c) are also consistent with type-II superconductivity and the system being in the mixed state.

The temperature dependence of $\lambda_{ab,c}^{-2}(T)$ is displayed in Fig. 5. It shows a similar behavior to $\lambda_{ab}^{-2}(T)$, including a saturation behavior below about 6.4 K and a plateau with an inflection point at intermediate temperatures, which also indicate a nodeless multigap structure. Correspondingly, the single gap s -wave and d -wave models also poorly describe the data, as shown in Fig. 5. $\lambda_{ab,c}^{-2}(T)$ were fitted using a two-gap $s + s$ model [Eq. (5)], which well fits the data with $\lambda_{ab,c}(0) = 488(4)$ nm, $\Delta_{s,1}(0) = 17(2)$ meV, $\Delta_{s,2}(0) = 2.0(3)$ meV,

$x = 0.74(2)$, and $T_c = 47.0(6)$ K. The fitted parameters for other models are also summarized in Table I.

IV. DISCUSSION AND CONCLUSIONS

Our TF- μSR measurements for fields applied both parallel and perpendicular to the c axis evidence a nodeless two-gap order parameter in $(\text{TBA})_{0.3}\text{FeSe}$. The magnitudes of the larger of the two gaps obtained for the two field directions are $4.2(5)k_B T_c$ [$17(2)$ meV] and $2.3(2)k_B T_c$ ($9.7(6)$ meV) for

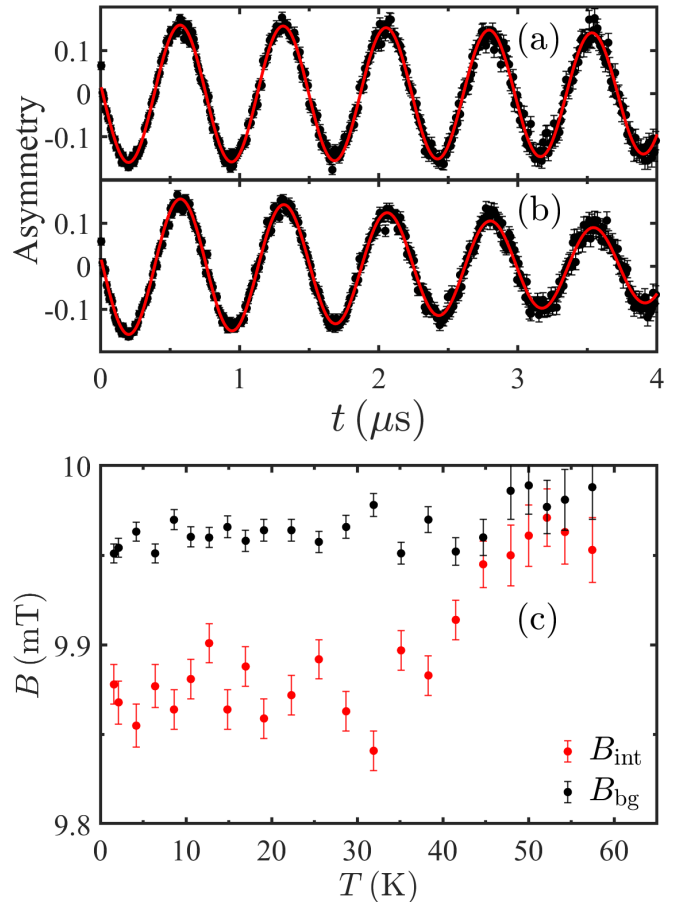


FIG. 4. Transverse-field μSR spectra in a field of 10 mT applied parallel to the ab plane at (a) 57.4 and (b) 1.6 K. The solid lines are the fitted curves using Eq. (2). (c) Temperature dependence of the internal magnetic field B_{int} and the magnetic field in the silver holder (background) B_{bg} obtained from the fitting.

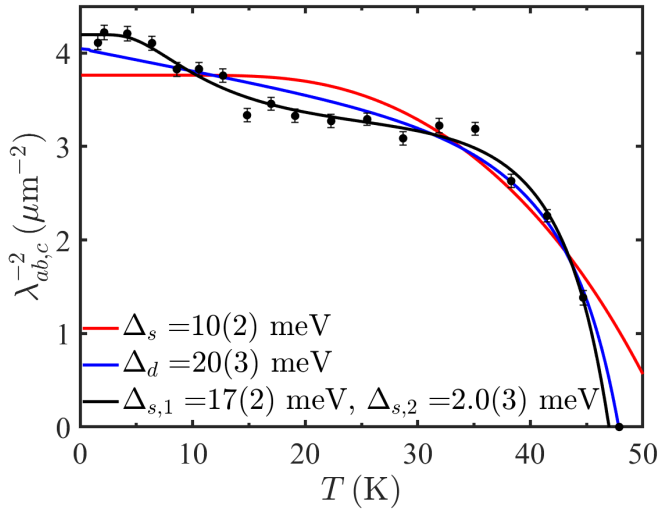


FIG. 5. Temperature dependence of $\lambda_{ab,c}^{-2}$ of $(\text{TBA})_{0.3}\text{FeSe}$ with a transverse field of 10 mT parallel to the ab plane. The solid lines are the fitting curves of a single s -wave model (red line), a single d -wave model (blue line) and a two-gap $s+s$ wave model (black line), respectively.

$\lambda_{ab,c}^{-2}$ and λ_{ab}^{-2} , respectively, where the former is very close to the gap magnitude of 16 meV from scanning tunneling spectroscopy (STS) [41]. On the other hand, the smaller gaps are $0.49(8)k_B T_c$ [2.0(3) meV] and $0.3(1)k_B T_c$ [1.5(4) meV] for the respective directions, which are significantly smaller than the large gaps. Note that although the χ^2 of the d -wave and $s+s$ model fitting are similar for $H \parallel c$ (Table I), it can be seen that the d -wave model systematically gives a poorer fit at low temperatures, failing to capture the inflection and low temperature flattening of the data. Moreover, for $H \perp c$ the χ^2 of the $s+s$ -wave fit is considerably lower.

The occurrence of two-gap nodeless superconductivity is in contrast to a TF- μ SR study of FeSe, where there is a lack of saturation and a linear temperature dependence of $\lambda_{ab}^{-2}(T)$, which was best described by a model with one nodal and one nodeless gap [23], although the presence of deep gap minima has also been suggested [25–27]. The nodal order parameter was proposed to correspond to an $s+d$ wave pairing state with nodes on the electron pocket [23]. The presence of nodes has also been inferred from STS measurements of FeSe thin-films [45], while such measurements on monolayer FeSe/SrTiO₃ consistently show U-shaped spectra [6,46] indicating nodeless pairing upon reducing the dimensionality.

Two-gap nodeless superconductivity has also been reported from penetration depth studies of the intercalated FeSe based superconductor $(\text{Li}_{0.84}\text{Fe}_{0.16})\text{OHFe}_{0.98}\text{Se}$ [36,37], which also has an enhanced transition temperature of $T_c = 42$ K. Similarly, nodeless superconductivity was inferred from angle-resolved photoemission spectroscopy measurements of the intercalated systems $(\text{Tl}, \text{K})\text{Fe}_{1.78}\text{Se}_2$ [47] and $\text{A}_x\text{Fe}_2\text{Se}_2$ ($\text{A} = \text{K}, \text{Cs}$) [48]. For $(\text{Li}_{0.84}\text{Fe}_{0.16})\text{OHFe}_{0.98}\text{Se}$, the two gap magnitudes deduced from the analysis of λ_{ab}^{-2} are more similar with $\Delta_2/\Delta_1 \approx 0.5 - 0.6$ [36,37], as opposed to a correspond-

ing ratio of 0.15 from our study of $(\text{TBA})_{0.3}\text{FeSe}$. On the other hand, a small gap of 1.05 meV is deduced from the analysis of the out-of-plane penetration depth of $(\text{Li}_{0.84}\text{Fe}_{0.16})\text{OHFe}_{0.98}\text{Se}$ [36], as compared to a larger gap of 9.7 meV. Interestingly these gap parameters are close to those obtained here from λ_{ab}^{-2} of $(\text{TBA})_{0.3}\text{FeSe}$. These similarities between the studies of $(\text{Li}_{0.84}\text{Fe}_{0.16})\text{OHFe}_{0.98}\text{Se}$ and $(\text{TBA})_{0.3}\text{FeSe}$ suggest that such multigap nodeless structures are a common feature of quasi-two-dimensional FeSe-based superconductors, seemingly different from bulk FeSe.

In $(\text{Li}_{0.84}\text{Fe}_{0.16})\text{OHFe}_{0.98}\text{Se}$, the opening of a small gap only for the out-of-plane direction is ascribed to superconductivity in the spacer layers being induced by proximity to the FeSe layers [36]. A two-dimensional nature of the superconductivity in $(\text{TBA})_{0.3}\text{FeSe}$ is inferred from the observation of a Berezinskii-Kosterlitz-Thouless transition, as well as a pseudogap phase above the superconducting T_c , where strong phase fluctuations prevent zero-resistance [41], and the resistivity ratio ρ_c/ρ_{ab} of around 10^5 exceeds that of ≈ 2500 for $(\text{Li}_{0.84}\text{Fe}_{0.16})\text{OHFe}_{0.98}\text{Se}$ [49]. However, in $(\text{TBA})_{0.3}\text{FeSe}$, this two-dimensionality is not as readily inferred from the anisotropy of the superfluid density. Another notable feature of $(\text{Li}_{0.84}\text{Fe}_{0.16})\text{OHFe}_{0.98}\text{Se}$ is that clear evidence for magnetic ordering is found from the observation of coherent oscillations in ZF- μ SR [36]. No such oscillations are found in our ZF- μ SR measurements of $(\text{TBA})_{0.3}\text{FeSe}$, and there is only a weak temperature dependence of $\Lambda(T)$ [(Fig. 1(b))]. Note that muons stopping in the large organic spacer layers may not be sensitive to the magnetic fields generated in the FeSe layers, and therefore the ZF- μ SR results may not reflect the intrinsic spin dynamics of the iron ions.

In summary, we performed μ SR measurements to probe the superconducting order parameter of $(\text{TBA})_{0.3}\text{FeSe}$. No evidence is found for time-reversal symmetry breaking below T_c from ZF- μ SR, while the analysis of the temperature dependence of the superfluid density deduced from TF- μ SR with magnetic fields $H \parallel c$ and $H \perp c$ shows that there are multiple nodeless superconducting gaps, which is in line with the findings for the intercalated FeSe based superconductor $(\text{Li}_{0.84}\text{Fe}_{0.16})\text{OHFe}_{0.98}\text{Se}$. Given that there are conflicting reports as to whether $(\text{Li}_{0.84}\text{Fe}_{0.16})\text{OHFe}_{0.98}\text{Se}$ exhibits a sign-changing [38,50,51], or sign preserving [35] s -wave pairing state, it is of particular interest to examine this aspect of the order parameter of $(\text{TBA})_{0.3}\text{FeSe}$, such as by utilizing quasiparticle interference techniques or by looking for a spin-resonance mode.

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