Optical properties of the iron-chalcogenide superconductor FeTe_{0.55}Se_{0.45}

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Abstract

The complex optical properties of the iron-chalcogenide superconductor FeTe_{0.55}Se_{0.45} with $T_c = 14$ K have been examined over a wide frequency range for light polarized in the Fe-Te(Se) planes above and below $T_c$. At room temperature the optical response may be described by a weakly-interacting Fermi liquid; however, just above $T_c$ this picture breaks down and the scattering rate takes on a linear frequency dependence. Below $T_c$ there is evidence for two gap features in the optical conductivity at $\Delta_1 \approx 2.5$ meV and $\Delta_2 \approx 5.1$ meV. Less than 20% of the free carriers collapse into the condensate for $T \ll T_c$, and this material is observed to fall on the universal scaling line for a BCS dirty-limit superconductor in the weak-coupling limit.

Keywords: Superconductivity, Infrared spectroscopy, Optical properties, Iron chalcogenides, Order parameter

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1. Introduction

There have been many surprises in the field of superconductivity in the last 25 years. First and foremost was the discovery of superconductivity at elevated temperatures in the copper-oxide materials \cite{1}. For conventional metals and alloys, in the model developed by Bardeen, Cooper and Schrieffer (BCS) superconductivity is mediated through lattice vibrations where electrons form bound pairs \cite{2}; the condensation below the critical temperature ($T_c$) is also accompanied by the formation of an isotropic $s$-wave energy gap at the Fermi surface. Within this framework, it was thought that the critical temperature could not exceed approximately 30 K \cite{3}. With $T_c$’s in excess of 130 K at ambient pressure and an unusual $d$-wave energy gap with nodes at the Fermi surface, the pairing mechanism in the cuprates remains unresolved. The discovery of superconductivity in MgB$_2$ with the surprisingly high transition temperature of $T_c = 39$ K \cite{4} initially suggested an unusual pairing mechanism; however, the isotope effect established that the superconductivity in this material is likely phonon mediated \cite{5}. In this particular case, the high phonon frequencies in MgB$_2$ are likely responsible for the unusually large value for $T_c$ \cite{4}.

The discovery of superconductivity in the iron-arsenic LaFeAsO$_{1-x}$F$_x$ (1111) pnictide compound \cite{7} was surprising because iron had long been considered detrimental to superconductivity. Different rare earth substitutions in this material quickly raised $T_c \gtrsim 50$ K \cite{8,9,10}. While such high values for $T_c$ do not definitively rule out a phonon-mediated pairing mechanism, the presence of magnetic order close to the superconductivity in these compounds \cite{11} has led to the suggestion that the pairing in this class of materials may have another origin \cite{12,13}. The highest $T_c$’s have been observed in the 1111-family of materials; however, large single crystals have only recently been obtained \cite{14}; the extended unit cell of non-superconducting LaFeAsO is shown in Fig. 1(a). As a consequence, much of the focus has shifted to the structurally-simpler AFe$_2$As$_2$ (122) iron-pnictides, shown in Fig. 1(b), and the FeTe(Se) (11) iron-chalcogenide materials, shown in Fig. 1(c), where large single crystals are available. The metallic AFe$_2$As$_2$ materials (where $A = Ca$, Ba or Sr) have been extensively studied; in BaFe$_2$As$_2$ the application of pressure results in $T_c \approx 29$ K, while Co- and Ni-doping yields $T_c \approx 23$ K at ambient pressure \cite{15,16,17}. Superconductivity has been observed in the arsenic-free iron-chalcogenide FeSe compound with $T_c = 8$ K, which increases to $T_c \approx 37$ K under pressure \cite{18,19,20,21}. Through the substitution of Se for Te the critical temperature at ambient pressure reaches a maximum $T_c = 14$ K in FeTe$_{0.55}$Se$_{0.45}$. Despite the structural differences of the iron-pnictides and the iron-chalcogenides illustrated in Fig 1 the band structure of these materials is remarkably similar, with a minimal description consisting of an electron band ($\beta$) at the $M$ point and a hole band ($\alpha$) centered at the $\Gamma$ point of the Brillouin zone \cite{22}.

There have been a number of studies of the Fe$_{1-x}$Te and FeTe$_{1-x}$Se$_x$ materials, including transport \cite{23,24,25,26}, tunneling \cite{27}, and angle-resolved photoemission (ARPES) \cite{28,29}, with particular attention on the magnetic properties \cite{29,30,31,32,33,34,35,36}. While the optical properties of the superconducting iron-pnictides have been investigated in considerable detail \cite{37,38,39,40,41,42,43,44,45}, the iron-chalcogenide materials are relatively unexplored \cite{20}. In this work we examine the in-plane complex optical properties of a single crystal of superconducting FeTe$_{0.55}$Se$_{0.45}$ above and below $T_c$. At room temperature this material may be described as a weakly-interacting Fermi liquid and the transport is Drude-like. However, just above $T_c$ this interpretation breaks...
down and the scattering rate adopts a linear frequency dependence. Below \( T_c \) there are clear signatures of the superconductivity in the reflectance and the optical conductivity. Less than 20% of the free carriers collapse into the superconducting condensate, suggesting that this material is in the dirty limit, and this material is observed to fall on the scaling line predicted for a BCS dirty-limit superconductor in the weak-coupling limit. In addition, there is evidence for two gap features in at \( \Delta_1 \approx 2.5 \text{ meV} \) and \( \Delta_2 \approx 5.1 \text{ meV} \). Some of these results have been discussed in a previous work [46].

2. Results and Discussion

Single crystals with good cleavage planes (001) were grown by a unidirectional solidification method with a nominal composition of Fe\( \text{Te}_{0.55}\text{Se}_{0.45} \). The normal-state resistivity is in good agreement with literature values [24]. The critical temperature determined from magnetic susceptibility is \( T_c = 14 \text{ K} \) with a transition width of \( \lesssim 1 \text{ K} \). The temperature dependence of the reflectance has been measured at a near-normal angle of incidence on a freshly-cleaved surface above and below \( T_c \) over a wide frequency range (\( \approx 2 \text{ meV} \) to 3.5 eV) for light polarized in the \( a-b \) planes using an \textit{in situ} overcoating technique [47]. The reflectance in the terahertz and far-infrared region is shown in Fig. [2] the reflectance at 295 K is shown over a wider region in the inset. At room temperature, the reflectance displays the \( R \propto 1 - \sqrt{\omega} \) response characteristic of a metal in the Hagen-Rubens regime; however, at low temperature just above above \( T_c \) (18 K), \( R \propto 1 - \omega \), which is reminiscent of the reflectance in a marginal Fermi liquid [48, 49]. The development of the superconducting state has a clear signature in the reflectance. However, the reflectance is a complex quantity consisting of an amplitude and a phase, \( \mathbf{R} = \mathbf{R}e^{i\theta} \). Normally, only the amplitude \( R = |\mathbf{R}| \) is measured so it is not intuitively obvious what changes in the reflectance imply. For this reason, the complex optical properties have been calculated from a Kramers-Kronig analysis of the reflectance [50].

The temperature dependence of the real part of the infrared optical conductivity is shown in Fig. [3]. At room temperature the conductivity is flat and relatively featureless except for an infrared-active \( E_g \) mode at 204 cm\(^{-1}\) which involves the in-plane displacements of the Fe-Te(Se) atoms [51]. As the temperature is lowered there is a shift in the spectral weight from high to low frequency, where the spectral weight is defined here as the area under the conductivity curve over a given interval

\[
N(\omega, T) = \int_{0}^{\omega} \sigma_{1}(\omega', T) d\omega'.
\]

This is the expected response for a metallic system where the scattering rate decreases with temperature. The optical conductivity over most of the temperature range is described quite well by a simple Drude-Lorentz model for the dielectric function \( \tilde{\epsilon} = \epsilon_1 + i\epsilon_2 \),

\[
\tilde{\epsilon}(\omega) = \epsilon_\infty - \frac{\omega_{p,D}^2}{\omega^2 + i\omega/\tau_D} + \sum_j \frac{\Omega_j^2}{\omega_j^2 - \omega^2 - i\omega \gamma_j},
\]

where \( \epsilon_\infty \) is the real part of the dielectric function at high frequency, \( \omega_{p,D}^2 = 4\pi e^2 n/m^* \) and \( 1/\tau_D \) are the square of the plasma frequency and scattering rate for the delocalized (Drude) carriers, respectively; \( \omega_j \), \( \gamma_j \) and \( \Omega_j \) are the position, width, and strength of the \( j \)th vibration or excitation. The complex conductivity is \( \sigma(\omega) = \sigma_1 + i\sigma_2 = -i\omega|\tilde{\epsilon}(\omega) - \epsilon_\infty|/4\pi \). The optical conductivity may be reproduced using this model at 295, 200 and 100 K, with fitted values of \( \omega_{p,D} \) of 7200 cm\(^{-1}\) and \( 1/\tau_D = 414, 363 \) and 317 cm\(^{-1}\), respectively (\( \pm 5\% \)). To fit the midinfrared component, Lorentzian oscillators at the some-
what arbitrary positions of 650 and 3200 cm\(^{-1}\) have been introduced; the results of the fit to the data at 100 K are shown in the inset in Fig. 3. An alternative method that has been used in some of the pnictide materials considers two Drude components \[44, 52\]. If we apply this approach to the data at 100 K, we note that the “narrow” Drude component is similar to that obtained from the Drude-Lorentz fits, while the “broad” Drude component has a width and strength similar to that of the oscillator \(\omega_0\) at \(\approx 650\) cm\(^{-1}\). While the scattering rates are expected to change with temperature, the plasma frequencies should remain relatively constant. However, just above \(T_c\) at 18 K the plasma frequency of the narrow Drude component has unexpectedly decreased by more than a factor of two. In addition, at 18 K neither the Drude-Lorentz or the two-Drude model accurately the shape of the low-frequency conductivity. To address this problem, we consider the extended-Drude model in which both the shape of the low-frequency conductivity. To address this problem, we consider the extended-Drude model in which both

\begin{equation}
\frac{1}{\tau(\omega)} = \frac{\omega_p^2}{4\pi} \frac{1}{\sigma(\omega)} \frac{1}{\tau(\omega)} ,
\end{equation}

and

\begin{equation}
m^*(\omega) \equiv \frac{\omega^2}{4\pi m} \frac{1}{\sigma(\omega)} \frac{1}{\tau(\omega)} .
\end{equation}

In this instance we set \(\omega_p \equiv \omega_{p,D}\) and \(\epsilon_\infty = 4\) (although the choice of \(\epsilon_\infty\) has little effect on the scattering rate or the effective mass in the far-infrared region). The temperature dependence of \(1/\tau(\omega)\) is shown in Fig. 4 and the inset shows the temperature dependence of \(m^*(\omega)/m_b\). At 295, 200 and 100 K the scattering rate displays little frequency dependence, and moreover \(1/\tau(\omega) \to 0\) \(\approx 1/T_D\). This self-consistent behavior indicates that within this temperature range, the transport may be described as a weakly-interacting Fermi liquid (Drude model). However, just above \(T_c\) at 18 K the scattering rate develops a linear frequency dependence \(\approx 200\) cm\(^{-1}\), suggesting the presence of electronic correlations. This may be due in part to magnetic correlations \[31\] that arise from the suppression of the magnetic transition in Fe\(_{1-x}\)Te at \(T_N \approx 70\) K in response to Se substitution \[24\]. We note that similar behavior of the scattering rate is observed in many optimally-doped cuprate superconductors where the electronic correlations may have a similar origin \[55\]. Dramatic changes are also observed in 1/\(\tau(\omega)\) below \(T_c\) where the scattering rate is suppressed at low frequencies, but increases rapidly and overshoots the normal-state (18 K) value at about 60 cm\(^{-1}\), finally merging with the normal-state curve at about 200 cm\(^{-1}\); this behavior is in rough agreement with a recently proposed differential sum rule for the scattering-rate \[56, 57, 58\]. We note that the overshoot in 1/\(\tau(\omega)\) below \(T_c\) is in general more characteristic of a material with an s-wave gap rather than a higher-order d-wave gap \[59\]. In the normal state just above \(T_c\), the low-frequency limit for the effective mass is \(m^*(\omega \to 0)/m_b \approx 7 - 8\), which is comparable with ARPES estimates of \(m^*/m_b \approx 6 - 20\) \[60\]. This mass enhancement is too large to be caused by electron-phonon interactions or coupling to spin-fluctuations alone \[61\], suggesting that electronic correlations play a dominant role in the low-energy excitations in this material \[62, 63\]. This result is quite different than the behav-

Figure 3: The real part of the optical conductivity of FeTe\(_{0.55}\)Se\(_{0.45}\) in the infrared region for light polarized in the a-b planes at several temperatures above and below \(T_c\). In the normal state there is a gradual shift in spectral weight from high to low frequency. Below \(T_c\) there is a dramatic loss of the low-frequency scattering rate and \(\epsilon_\infty\). Inset: The Drude-Lorentz fit to the data at 100 K. The Drude component (dashed line) and Lorentzian oscillators (long-dashed lines) combine (dotted line) to reproduce the data (solid line) quite well.

Figure 4: The in-plane frequency-dependent scattering rate of FeTe\(_{0.55}\)Se\(_{0.45}\) for several temperatures above and below \(T_c\) in the far-infrared region. The values for \(1/\tau(\omega)\) are shown at 295 (\(\Delta\)), 200 (\(\bullet\)) and 100 K (\(\large\triangle\)), respectively, where the scattering rate displays little temperature dependence. For \(T \gtrsim T_c\) (18 K) at low frequency \(1/\tau(\omega) \approx \omega_\infty\) while for \(T < T_c\) large changes in the scattering rate are observed in response to the formation of superconducting gap(s). Inset: The frequency dependence of the effective mass.
The low-frequency conductivity is shown in more detail in Fig. 5. For $T < T_c$ there is a dramatic suppression of the low-frequency conductivity with a commensurate loss of spectral weight. This “missing area” is associated with the formation of a superconducting condensate, whose spectral weight $N_c$ may be calculated from the Ferrell-Glover-Tinkham sum rule \[^{[66, 67]}\] for an isotropic superconducting energy gap $\Delta_0$.

$$N_c \equiv N(\omega_c, T \approx T_c) - N(\omega_c, T < T_c) = \omega_{p,S}^2 / 8.$$  

Here $\omega_{p,S}^2 = 4\pi n_F e^2 / m^*$ is the square of the superconducting plasma frequency and superfluid density is $\rho_0 \equiv \omega_{p,S}^2$; the cutoff frequency $\omega_c \approx 150 \text{ cm}^{-1}$ is chosen so that the integral converges smoothly. The superconducting plasma frequency has also been determined from $\epsilon_1(\omega)$ in the low frequency limit where $\epsilon_1(\omega) = \epsilon_{\infty} - \omega_{p,S}^2 / \omega^2$. Yet another method of extracting $\omega_{p,S}$ from $\epsilon_1(\omega)$ is to determine $[-\omega^2 \epsilon_1(\omega)]^{1/2}$ in the $\omega \rightarrow 0$ limit \[^{[68]}\]. All three techniques yield $\omega_{p,S} = 3000 \pm 200 \text{ cm}^{-1}$, indicating that less than one-fifth of the free-carriers in the normal state have condensed ($\omega_{p,S}^2 / \omega_{p,D}^2 \lesssim 0.18$), implying that this material is not in the clean limit. The superfluid density can also be expressed as an effective penetration depth $\lambda_0 = 5300 \pm 300 \text{ Å}$, which is in good agreement with recent tunnel-diode \[^{[69]}\] and muon-spin spectroscopy \[^{[70]}\] measurements on FeTe$_{0.65}$Se$_{0.35}$ and FeTe$_{0.5}$Se$_{1.5}$, respectively.

The strong suppression of the conductivity for $T < T_c$ below $\approx 120 \text{ cm}^{-1}$ is characteristic of the opening of a superconducting energy gap in the density of states at the Fermi surface; in addition, there also appears to be a shoulder at $\approx 60 \text{ cm}^{-1}$. As previously noted, the formation of a gap leads to a transfer of spectral weight into the condensate. Below $T_c$, the optical conductivity has been calculated using a Mattis-Bardeen approach \[^{[71]}\] for the contribution from the gapped excitations \[^{[72]}\]. This method assumes that $l \lesssim \xi_0$, where the mean-free path $l = v_F \tau$ ($v_F$ is the Fermi velocity), and the coherence length $\xi_0 = h v_F / n \Delta_0$ for an isotropic superconducting energy gap $\Delta_0$; this may also be expressed as $1/\tau \gtrsim 2\Delta_0$. The dirty-limit approach is consistent with the observation that less than 20% of the free carriers collapse into the condensate. Initially, only a single isotropic gap $\Delta_0 \approx 4.5 \text{ meV}$ for $T < T_c$ was considered; however, even with a moderate amount disorder scattering ($1/\tau = 4\Delta_0$) and the addition of the low-frequency tail from the bound Lorentzian oscillators this fails to accurately reproduce the low-frequency conductivity, as Fig. 5(a) demonstrates. To properly model the conductivity two gap features have been considered, $\Delta_1 \approx 2.5 \text{ meV}$ and $\Delta_2 \approx 5.1 \text{ meV}$ with $1/\tau_1 = 4\Delta_j$ for $T < T_c$. The observation of two gap features is consistent with recent optical \[^{[40, 41]}\] and microwave \[^{[42]}\] measurements on the pnictide compounds, as well as some theoretical works that propose that s-wave gaps form on the hole ($\alpha$) and electron ($\beta$) pockets, possibly with a sign change between them \[^{[43, 44]}\], the so-called $s^\pm$ symmetry state. In the $s^\pm$ model, the gap on the electron pocket ($\beta$) may be an extended s-wave and have nodes on its Fermi surface \[^{[73]}\]. While there is some uncertainty associated with the low-frequency conductivity in this work, the apparent lack of residual conductivity in the terahertz region for $T < T_c$ suggests the absence of nodes. It is possible that disorder may lift the nodes, resulting in a nodeless extended s-wave gap \[^{[80, 81]}\]. Without the optical gaps provide estimates of the gap amplitudes, they do not distinguish between $s^\pm$ and extended s-wave. The optical gaps at $2\Delta_j \approx 40$ and $82 \text{ cm}^{-1}$ are either similar to or slightly larger than the low-frequency scattering rate observed at 18 K, $1/\tau(\omega \rightarrow 0) \approx 40 \text{ cm}^{-1}$. This might seem to suggest that the Mattis-Bardeen approach should not be used; however, the linear frequency dependence of the scattering rate complicates matters. If we consider the value of the scattering rate in the region of the optical gaps where the scattering should be important, then from Fig. 4 we have

$$\frac{1/\tau_j(2\Delta_j)}{2\Delta_j} \approx 3$$

which is actually larger than the ratio of 2 that was assumed in the calculation, indicating that the Mattis-Bardeen approach is valid. While the value of $2\Delta_1/k_B T_c \approx 4$ is close the value of 3.5 expected for a BCS superconductor in the weak-coupling limit, $2\Delta_2/k_B T_c \approx 8.4$ is significantly larger.
It was recently noted that in a number of the pnictide materials [82] the superfluid density $\rho_{s0}$ falls on a recently proposed empirical scaling relation for the cuprate superconductors shown by the dashed line in Fig. 6 [83, 84].

$$\rho_{s0}/8 \approx 4.4 \sigma_{dc} T_c \tag{8}$$

From the estimate of $\sigma_{dc} \equiv \sigma(\omega \to 0) = 3500 \pm 400 \mathrm{cm}^{-1}$ for $T \gtrsim T_c$ (determined from Fig. 3 as well as Drude-Lorentz fits), and the previously determined value of $\rho_{s0} = 9 \pm 1 \times 10^6 \mathrm{cm}^{-2}$, we can see that FeTe$_{0.55}$Se$_{0.45}$ also falls close to this scaling line. In fact, in a BCS dirty-limit superconductor in the weak-coupling limit, the numerical constant in the scaling relation is calculated to be slightly larger [83, 84]:

$$\rho_{s0}/8 \approx 8.1 \sigma_{dc} T_c \tag{9}$$

(dotted line in Fig. 6); the result for this material is actually closer to the BCS dirty-limit line than the one established for the cuprate superconductors.

3. Conclusion

To summarize, the optical properties of FeTe$_{0.55}$Se$_{0.45}$ ($T_c = 14$ K) have been examined for light polarized in the Fe-Te(Se) planes above and below $T_c$. Well above $T_c$, the transport may be described by a weakly-interacting Fermi liquid (Drude model); however, this picture breaks down close to $T_c$ when the scattering rate takes on a linear frequency dependence, similar to what is observed in the cuprate superconductors. Below $T_c$, less than one-fifth of the free carriers collapse into the condensate ($\delta_0 = 5300 \AA$), indicating that this material is in the dirty limit, and indeed this material falls on the general scaling line predicted for a BCS dirty-limit superconductor in the weak coupling limit. To successfully model the optical conductivity in the superconducting state, two gaps of $\Delta_1 \approx 2.5$ meV and $\Delta_2 \approx 5.1$ meV are considered using a Mattis-Bardeen formalism (with moderate disorder scattering), suggesting either an s- or a nodeless extended s-wave gap.

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References

[1] J. G. Bednorz, K. A. Muller, Z. Phys. B 64 (1986) 189.
[2] J. Bardeen, L. N. Cooper, J. R. Schrieffer, Phys. Rev. 108 (1957) 1175.
[3] W. L. McMillan, Phys. Rev. 167 (1968) 331.
[4] J. Nagamatsu, N. Nakagawa, T. Muranaka, Y. Zenitani, J. Akimitsu, Nature (London) 410 (2001) 63.
[5] S. L. Bud’ko, G. Lapertot, C. Petrovic, C. E. Cunningham, N. Anderson, P. C. Canfield, Phys. Rev. Lett. 86 (2001) 1877.
[6] J. Kontus, I. I. Mazin, K. D. Belashchenko, V. P. Antropov, L. L. Boyer, Phys. Rev. Lett. 86 (2001) 4656.
[7] Y. Kamihara, T. Watanabe, M. Hirano, H. Hosono, J. Am. Chem. Soc. 130 (2008) 3296.
[8] Z.-A. Ren, J. Yang, W. Lu, W. Yi, X.-L. Shen, Z.-C. Li, G.-C. Che, X.-L. Dong, L.-L. Sun, F. Zhou, Z.-X. Zhou, EPL 82 (2008) 57002.
[9] Z.-A. Ren, G.-C. Che, X.-L. Dong, J. Yang, W. Lu, W. Yi, X.-L. Shen, Z.-C. Li, L.-L. Sun, F. Zhou, Z.-X. Zhou, EPL 83 (2008) 17002.
[10] K. Ishida, Y. Nakai, H. Hosono, J. Phys. Soc. Japan 78 (2009) 062001.
[11] C. de la Cruz, Q. Huang, J. W. Lynn, J. Li, W. Ratcliff II, J. L. Zarestky, H. A. Mook, G. F. Chen, J. L. Luo, N. L. Wang, P. Dai, Nature (London) 453 (2008) 899.
[12] L. Boeri, O. V. Dolgov, A. A. Golubov, Phys. Rev. Lett. 101 (2008) 026403.
[13] D. C. Johnston (2010). arXiv:1005.4392 [unpublished]
[14] J.-Q. Yan, S. Nandi, J. L. Zarestky, W. Tian, A. Kreyssig, B. Jensen, A. Kracher, K. W. Dennis, R. J. McKeeveney, A. I. Goldman, R. W. McCartney, T. A. Lograsso, Appl. Phys. Lett. 95 (2009) 222504.
[15] W. L. Alireza, Y. T. C. Ko, J. K. Budai, C. M. Petrone, J. M. Cole, G. G. Lonzarich, S. E. Sebastian, J. Phys.: Condens. Matter 21 (2009) 012208.
[16] A. S. Sefat, R. Jin, M. A. McGuire, B. C. Sales, D. J. Singh, D. Mandrus, Phys. Rev. Lett. 101 (2008) 117004.
[17] L. J. Li, Q. B. Wang, Y. K. Luo, H. Chen, Q. Tao, Y. K. Li, M. He, Z. W. Zhu, G. H. Cao, Z. A. Xu, New J. Phys. 11 (2009) 025008.
[18] F.-C. Hsu, J.-Y. Luo, K.-W. Yeh, T.-K. Chen, T.-W. Huang, P. M. Wu, Y.-C. Lee, Y.-L. Huang, Y.-Y. Chu, D.-C. Yan, M.-K. Wu, Proc. Natl. Acad. Sci. U.S.A. 105 (2008) 14262.
[19] M. H. Fang, H. M. Pham, B. Qian, T. J. Liu, E. K. Vehstedt, Y. Liu, Z. W. Zhu, G. H. Cao, Z. A. Xu, New J. Phys. 11 (2009) 025008.
[20] S. Raghu, X.-L. Qi, C.-X. Liu, D. J. Scalapino, S.-C. Zhang, Phys. Rev. B 77 (2008) 220503(R).
[21] M. H. Fang, H. M. Pham, B. Qian, T. J. Liu, E. K. Vehstedt, Y. Liu, L. Spinu, Z. Q. Mao, Phys. Rev. B 78 (2008) 224503.
[22] B. C. Sales, A. S. Sefat, M. A. McGuire, R. Y. Jin, D. Mandrus, Y. Mozharivskyj, Phys. Rev. B 79 (2009) 094521.
[23] T. Taen, Y. Tsuchiya, Y. Nakajima, T. Tamegai, Phys. Rev. B 80 (2009) 092502.
[24] G. F. Chen, Z. G. Chen, J. Dong, W. Z. Hu, G. Li, X. D. Zhang, P. Zheng, J. L. Luo, N. L. Wang, Phys. Rev. B 79 (2009) 140509(R).
