Nodeless superconductivity in Ir$_{1-x}$Pt$_x$Te$_2$ with strong spin-orbital coupling

S. Y. Zhou$^1$, X. L. Li$^1$, B. Y. Pan$^1$, X. Qiu$^1$, J. Pan$^1$, X. C. Hong$^1$, Z. Zhang$^1$, A. F. Fang$^2$, N. L. Wang$^2$ and S. Y. Li$^1$(a)

1 State Key Laboratory of Surface Physics, Department of Physics, and Laboratory of Advanced Materials, Fudan University - Shanghai 200433, PRC
2 Beijing National Laboratory for Condensed Matter Physics, Institute of Physics, Chinese Academy of Sciences Beijing 100190, PRC

received 28 August 2013; accepted in final form 28 October 2013
published online 22 November 2013

PACS 74.25.fc - Electric and thermal conductivity
PACS 74.40.Kb - Quantum critical phenomena
PACS 74.25.Op - Mixed states, critical fields, and surface sheaths

Abstract – The thermal conductivity $\kappa$ of the superconductor Ir$_{1-x}$Pt$_x$Te$_2$ ($x = 0.05$) single crystal with strong spin-orbital coupling was measured down to 50 mK. The residual linear term $\kappa_0/T$ is negligible in zero magnetic field. In low magnetic field, $\kappa_0/T$ shows a slow field dependence. These results demonstrate that the superconducting gap of Ir$_{1-x}$Pt$_x$Te$_2$ is nodeless, and the pairing symmetry is likely a conventional $s$-wave, despite the existence of strong spin-orbital coupling and a quantum critical point.

Copyright © EPLA, 2013

Introduction. – The effect of strong spin-orbital coupling (SOC) on superconductivity has recently attracted much attention. One example is the topological superconductor, such as the candidate Cu$_x$Bi$_2$Se$_3$ in which Cu atoms are intercalated into the topological insulator Bi$_2$Se$_3$ with strong SOC [1]. A novel superconducting state was suggested in Cu$_x$Bi$_2$Se$_3$ by the point-contact spectra and superfluid density measurements [2,3]. Another example is the noncentrosymmetric superconductor, such as Li$_2$Pt$_3$B, in which the spatial inversion symmetry is broken [4]. The strong SOC in Li$_2$Pt$_3$B gives a large spin-triplet pairing component and produces line nodes in the superconducting gap [5,6].

More recently, superconductivity was discovered in the layered compound IrTe$_2$ by Pd intercalation (Pd$_x$IrTe$_2$) [7], Pd substitution (Ir$_{1-x}$Pd$_x$Te$_2$) [7], Pt substitution (Ir$_{1-x}$Pt$_x$Te$_2$) [8], or Cu intercalation (Cu$_x$IrTe$_2$) [9]. Since the SOC is proportional to $Z^4$, where $Z$ is the atomic number, the superconductivity in doped IrTe$_2$ must be associated with strong SOC due to the large $Z$.

Furthermore, the parent compound IrTe$_2$ exhibits an intriguing structural phase transition from a high-temperature trigonal to a low-temperature monoclinic phase near 270 K [7]. Initially it was related to a charge-density wave (CDW) induced by the Ir 5$d$ $t_{2g}$ orbitals [7,10]; however, later no CDW gap was detected from the optical spectroscopy [11] and angle-resolved photoemission spectroscopy (ARPES) [12] measurements. At this moment, the origin of the transition is still under hot debate, with proposals such as the crystal field effect [13], Pd substitution (Ir$_{1-x}$Pd$_x$Te$_2$) [14], or Cu intercalation (Cu$_x$IrTe$_2$) [15]. Such a phase diagram of doped IrTe$_2$ is reminiscent of high-$T_c$ cuprates and some heavy-fermion superconductors, in which superconductivity appears close to a magnetic quantum critical point (QCP). This means that there likely exists a QCP under the superconducting dome of doped IrTe$_2$, and the superconductivity may be unconventional [15]. Therefore, it is of great interest to investigate whether there is a novel superconducting state in doped IrTe$_2$.

The ultra-low–temperature thermal-conductivity measurement is a bulk tool to study the gap structure of superconductors [16]. The existence of a finite residual linear term $\kappa_0/T$ in zero field is usually considered as the
signature of a nodal superconducting gap. Further information of nodal gap, gap anisotropy, or multiple gaps may be obtained from the field dependence of $\kappa_0/T$ [16]. Previously, single-gap $s$-wave superconductivity near the QCP of CDW has been clearly shown in Cu$_2$TiSe$_2$ by thermal-conductivity measurements [17].

In this paper, we probe the superconducting gap structure of the Ir$_{1-x}$Pt$_x$Te$_2$ ($x = 0.05$) single crystal by measuring the thermal conductivity $\kappa$ down to 50 mK. The residual linear term $\kappa_0/T$ is negligible in zero magnetic field. The field dependence of $\kappa_0/T$ is slow at low field, unlike that of a nodal superconductor. Both results suggest nodeless superconducting gap in Ir$_{1-x}$Pt$_x$Te$_2$.

**Experimental details.** – Single crystals of Ir$_{1-x}$Pt$_x$Te$_2$ were grown via the self-flux method [11]. The dc magnetic susceptibility was measured by using a SQUID (MPMS, Quantum Design). The heat capacity was measured in a physical property measurement system (PPMS, Quantum Design) via the relaxation method. The Ir$_{0.95}$Pt$_{0.05}$Te$_2$ single crystal was cut to a rectangular shape of dimensions 2.0 $\times$ 0.55 mm$^2$ in the $ab$ plane and of 20 $\mu$m thickness along the $c$-axis. Four silver wires were attached to the sample surface with silver paint, which were used for both in-plane resistivity and thermal-conductivity measurements. The contacts are metallic with a typical resistance of 15 m$\Omega$ at 2 K.

In-plane thermal conductivity was measured in a dilution refrigerator, using a standard four-wire steady-state method with two RuO$_2$ chip thermometers, calibrated in situ against a reference RuO$_2$ thermometer. Magnetic fields were applied along the $c$-axis and perpendicular to the heat current. To ensure a homogeneous field distribution in the sample, all fields for resistivity and thermal-conductivity measurements were applied at a temperature above $T_c$.

**Results and discussion.** – Figure 1(a) presents the dc magnetic susceptibility of the Ir$_{0.95}$Pt$_{0.05}$Te$_2$ single crystal. It was measured in the magnetic field $H = 20$ Oe parallel to the $ab$ plane, with zero-field–cooled (ZFC) and field-cooled (FC) processes. A sharp superconducting transition with $T_c \approx 3.0$ K and an about 100% shielding volume fraction were observed for the ZFC process, suggesting the homogeneous bulk superconductivity in our sample.

The heat capacity was measured on three pieces of Ir$_{0.95}$Pt$_{0.05}$Te$_2$ single crystals, with a total mass of 8.8 mg. It is plotted in fig. 1(b), as $C/T$ vs. $T^2$. Above $T_c$, the data can be well fitted by $C/T = \gamma + \beta T^2$, giving the electronic specific-heat coefficient $\gamma = 9.20$ mJ mol$^{-1}$ K$^{-2}$. The significant jump was observed at $T_c \approx 3.0$ K, which also indicates the high quality of our single crystals.

Figure 2(a) shows the resistivity of the Ir$_{0.95}$Pt$_{0.05}$Te$_2$ single crystal in zero field. No resistivity anomaly is observed above $T_c$, suggesting that no structural transition occurs in this sample near optimal doping. The data between 3.5 and 31 K can be fitted to $\rho(T) = \rho_0 + AT^n$, with $\rho_0 = 4.34 \pm 0.002 \mu\Omega$cm and $n = 3.01 \pm 0.03$. Figure 2(b) plots $\rho$ vs. $T^{3.01}$ and the solid line represents the fitting curve data in fig. 2(a). Such a temperature dependence of $\rho(T) \sim T^n$ with $n \approx 2.8$ has been observed in the Ir$_{1-x}$Pt$_x$Te$_2$ polycrystal and attributed to phonon-assisted interband scattering [8], as in TiSe$_2$ [18].

Previously, the upper critical field $H_{c2}(0) \approx 0.17$ T has been determined for the Ir$_{0.96}$Pt$_{0.04}$Te$_2$ polycrystal by resistivity measurements [8]. In order to obtain the $H_{c2}(0)$ of our Ir$_{0.95}$Pt$_{0.05}$Te$_2$ single crystal, we also measure its resistivity with magnetic field parallel to the $c$-axis up to $H = 0.75$ T and down to 50 mK, shown in fig. 2(c). In zero field, the resistivity drops to zero at 2.94 K with a narrow transition width of 0.06 K. The superconducting transition is gradually suppressed in magnetic fields. In fig. 2(d), we plot the temperature dependence of $H_{c2}$, where $T_c$ is defined at the resistivity $\rho$ dropping to zero on the curves in fig. 2(c). The dashed line is a guide to the
Fig. 2: (Color online) (a) Temperature dependence of the resistivity $\rho$ for the Ir$_{0.95}$Pt$_{0.05}$Te$_2$ single crystal. The data between 3.5 and 31 K can be fitted to $\rho(T) = \rho_0 + aT^n$, as shown by the solid line, with $\rho_0 = 4.34 \pm 0.002 \mu\Omega\text{cm}$ and $n = 3.01 \pm 0.03$. (b) The resistivity $\rho$ as a function of $T^{0.1}$ below 31 K. The solid line is the fitting curve in (a). (c) Low-temperature resistivity of the Ir$_{0.95}$Pt$_{0.05}$Te$_2$ single crystal in magnetic fields up to 0.75 T. (d) Temperature dependence of the upper critical field $H_{c2}$, defined at the point $\rho$ dropping to zero on the curves in (c). The dashed line is a guide to the eye, which points to $H_{c2}(0) \approx 0.09$ T.

eye, which points to $H_{c2}(0) \approx 0.09$ T. This value is only about half of that in the Ir$_{0.96}$Pt$_{0.04}$Te$_2$ polycrystal [8]. It is not surprising since in the polycrystal, due to the random grain orientation, the $H_{c2}$ represents the maximum value for all field configurations, presumably $H \parallel ab$.

The thermal conductivities of the Ir$_{0.95}$Pt$_{0.05}$Te$_2$ single crystal in zero and magnetic fields up to $H = 0.08$ T are plotted in fig. 3, as $\kappa/T$ vs. $T$. We fit all the curves to $\kappa/T = a + bT^{n-1}$, in which the two terms $aT$ and $bT^{n-1}$ represent contributions from electrons and phonons, respectively [19,20]. The power $a$ of the second term contributed by phonons is typically between 2 and 3 for single crystals, due to the specular reflections of phonons at the boundary [19,20]. In zero field, the fitting gives a residual linear term $\kappa_0/T \equiv a = -7 \pm 18 \mu\text{W K}^{-2} \text{cm}^{-1}$, with $\alpha = 2.91 \pm 0.04$. Since our experimental error bar is $5 \mu\text{W K}^{-2} \text{cm}^{-1}$, the $\kappa_0/T$ in zero field is essentially zero, compared to the normal-state Wiedemann-Franz law expectation $L_0/\rho_0 = 5.65 \text{mW K}^{-2} \text{cm}^{-1}$, with the Lorenz number $L_0 = 2.45 \times 10^{-8} \text{W} \text{K}^{-2} \text{cm}^{-1}$ and $\rho_0 = 4.34 \mu\Omega\text{cm}$.

For s-wave nodeless superconductors, there are no fermionic quasiparticles to conduct heat when $T \to 0$ since all electrons become Cooper pairs. Therefore, there is no residual linear term of $\kappa_0/T$, as seen in Nb [21] and InBi [22]. However, a finite $\kappa_0/T$ in zero field is usually observed in a superconductor with nodal gap, coming from the nodal quasiparticles [16]. For example, $\kappa_0/T = 1.41 \text{mW K}^{-2} \text{cm}^{-1}$ for the overdoped cuprate Tl$_2$Ba$_2$CuO$_{6+\delta}$ (Tl-2201), a $d$-wave superconductor with $T_c = 15$ K [23], and $\kappa_0/T = 17 \text{mW K}^{-2} \text{cm}^{-1}$ for the ruthenate Sr$_2$RuO$_4$, a $p$-wave superconductor with $T_c = 1.5$ K [24]. The magnitude of $\kappa_0/T$ is determined by the ratio of quasiparticle velocities parallel ($v_F$) and perpendicular ($v_D$) to the Fermi surface near the nodes [25,26]. For a two-dimensional $d$-wave superconductor with a gap maximum $\Delta_0$ and a density of $n$ planes per unit cell of height $c$, one has the formula [26]

$$\kappa_0/T \approx \frac{k_B^2 n}{6} \frac{m}{c} k_F \frac{v_F}{\Delta_0},$$

assuming $v_F \gg v_D$, where $k_B$ is Boltzmann constant, $k_F$ is the Fermi wave vector. This formula works well for cuprate superconductors [27,28].

For Ir$_{0.95}$Pt$_{0.05}$Te$_2$, we estimate the values of $k_F$, $v_F$ and $\Delta_0$ by using $H_{c2}(0) = \frac{\lambda_0 c_4^2}{2\pi^2}, \Delta(0) = 2.14k_B T_c$ (weak-coupling approximation), $\xi_0 = \frac{\hbar v_F}{\pi \Delta(0)}$, $m = \hbar k_F/v_F$ with $T_c = 2.94$ K and $H_{c2}(0) \approx 0.09$ T, where $\xi_0$ is the coherence length and $m$ is the electronic mass. With the obtained $\xi_0 \approx 605 \AA$, $v_F = 1.57 \times 10^6 \text{m/s}$, $k_F = 0.135 \AA^{-1}$ and $\Delta(0) = 0.54 \text{meV}$, eq. (1) gives $\kappa_0/T = 1.8 \text{mW K}^{-2} \text{cm}^{-1}$. This estimation is several orders of magnitude larger than the measured value and the error bar of $\kappa_0/T$. Therefore, unconventional superconductivity with gap nodes in Ir$_{0.95}$Pt$_{0.05}$Te$_2$ can be ruled out and a nodeless superconducting gap is suggested.

The field dependence of $\kappa_0/T$ will give more information on the superconducting gap structure. For a clean type-II s-wave superconductor with isotropic gap, $\kappa_0/T$ should grow exponentially with the field (above $H_{c1}$), as in Nb [21]. In the case of a nodal superconductor, $\kappa_0/T$
increases rapidly ($\sim H^{1/2}$) in the low field due to the Volovik effect [29], as in TI-2201 [23].

In fig. 3, all the curves in the magnetic field are roughly linear. Therefore, we fit these curves to $\kappa / T = a + b T^x$ with $a$ fixed to 2. We note that the temperature dependence and magnitude of phonon conductivity $b T^x$ change in the magnetic field. In zero field, the paired electrons do not scatter phonons, while in the magnetic field, the normal electrons inside the vortex core will scatter the phonons. This may explain the change of phonon conductivity in the field. When the field is increased to $H = 0.068 \text{T}$ and $0.08 \text{T}$, $\kappa / T = 5.53 \pm 0.03$ and $5.49 \pm 0.05 \text{mW K}^{-2} \text{cm}^{-1}$ are obtained, respectively. Both values are close to the normal-state Wiedemann-Franz law expectation $L_0 / \rho_0 = 5.65 \text{mW K}^{-2} \text{cm}^{-1}$. We take $H = 0.068 \text{T}$ as its bulk $H_{c2}(0)$. A slightly different $H_{c2}(0)$ does not affect our discussion on the field dependence of $\kappa_0 / T$ below.

In fig. 4, the normalized $\kappa_0 / T$ of the Ir$_{0.95}$Pt$_{0.05}$Te$_2$ single crystal is plotted as a function of $H / H_{c2}$. For comparison, we also plot the data of the clean $s$-wave superconductor Nb [21], the dirty $s$-wave superconducting alloy InBi [22], the single-gap $s$-wave superconductor Cu$_{0.06}$TiSe$_2$ [17], the multi-band $s$-wave superconductor NbSe$_2$ [30], and an overdoped $d$-wave cuprate superconductor TI-2201 [23].

![Fig. 4: (Color online) Normalized residual linear term $\kappa_0 / T$ of Ir$_{0.95}$Pt$_{0.05}$Te$_2$ as a function of $H / H_{c2}$. For comparison, similar data are shown for the clean $s$-wave superconductor Nb [21], the dirty $s$-wave superconducting alloy InBi [22], the single-gap $s$-wave superconductor Cu$_{0.06}$TiSe$_2$ [17], the multi-band $s$-wave superconductor NbSe$_2$ [30], and an overdoped $d$-wave cuprate superconductor TI-2201 [23].](image)

Both InBi and Cu$_{0.06}$TiSe$_2$ are in the dirty limit $l < \xi_0$ [17,22]. To check whether Ir$_{0.95}$Pt$_{0.05}$Te$_2$ is also in the dirty limit, we need to estimate its electron mean free path $l$. Since the normal-state thermal conductivity $\kappa N / T = \frac{1}{2}C v_F l$, we have $l = 3(\kappa N / T)/(\gamma v_F)$ [17]. Here, $\kappa N / T$ is the Wiedemann-Franz law expectation $L_0 / \rho_0 = 5.65 \text{mW K}^{-2} \text{cm}^{-1}$, $\gamma = 9.20 \text{mJ mol}^{-1} \text{K}^{-2}$, and $v_F = 1.29 \times 10^5 \text{m/s}$ is calculated through $\rho_0 = \frac{\hbar v_F}{2 \Delta_{\text{min}}}$ and $\Delta(0) = 1.76 k_B T_c$ (in the BCS $s$-wave scenario). We obtained $l \approx 60 \text{Å}$, one order of magnitude smaller than $\xi_0$. This confirms that Ir$_{0.95}$Pt$_{0.05}$Te$_2$ is indeed in the dirty limit.

A nodeless gap does not essentially mean conventional superconductivity. For example, the nodeless gap observed in an optimally doped iron-based superconductor may be an unconventional $s_\pm$-wave resulting from antiferromagnetic spin fluctuations [31]. However, here in Ir$_{1-x}$Pt$_x$Te$_2$, the nodeless gap is unlikely a $s_\pm$-wave. One obvious reason is that Ir$_{1-x}$Pt$_x$Te$_2$ does not have that kind of multiple Fermi surfaces as in iron-based superconductors [7,31]. Therefore, the pairing symmetry in Ir$_{1-x}$Pt$_x$Te$_2$ is likely a conventional $s$-wave. In this sense, the appearance of superconductivity may have little relationship with quantum fluctuations near QCP. This is not surprising since the single-gap $s$-wave superconductivity in Cu$_x$TiSe$_2$ is also not related to quantum fluctuations of CDW.

From the aspect of strong SOC, an unconventional odd-parity pairing state was claimed for the topological superconductor candidate Cu$_x$Bi$_2$Se$_3$ [2,3]. Both nodeless gap or gap with point nodes are allowed for the odd-parity superconducting state [2]. In this context, more experiments such as point-contact spectra are needed to completely exclude novel superconductivity in Ir$_{1-x}$Pt$_x$Te$_2$.

**Conclusion.** – In summary, we investigated the superconducting gap structure of the Ir$_{1-x}$Pt$_x$Te$_2$ ($x = 0.05$) single crystal by thermal-conductivity measurements. The $\kappa_0 / T$ in zero field is negligible, and the field dependence of $\kappa_0 / T$ is slow in the low field. Both of them suggest nodeless superconductivity in Ir$_{1-x}$Pt$_x$Te$_2$. The pairing symmetry is likely a conventional $s$-wave, although the odd-parity superconducting state cannot be completely excluded from our measurements.

We thank J. J. YANG, Y. S. OH, and S.-W. CHEONG for providing Ir$_{0.96}$Pt$_{0.04}$Te$_2$ polycrystal and nominal Ir$_{0.8}$Pt$_{0.2}$Te$_2$ single crystal to initialize this study. This work is supported by the Natural Science Foundation of China, the Ministry of Science and Technology of China (National Basic Research Program Nos. 2009CB922903 and 2012CB821402), and the Program for Professor of
Special Appointment (Eastern Scholar) at Shanghai Institutions of Higher Learning.

REFERENCES

[1] Hor Y. S., Williams A. J., Checkelsky J. G., Roushan P., Seo J., Xu Q., Zandbergen H. W., Yazdani A., Ong N. P. and Cava R. J., Phys. Rev. Lett., 104 (2010) 057001.

[2] Sasaki S., Krienier M., Segawa K., Yada K., Tanaka Y., Sato M. and Ando Y., Phys. Rev. Lett., 107 (2011) 217001.

[3] Krienier M., Segawa K., Sasaki S. and Ando Y., Phys. Rev. B, 86 (2012) 180505(R).

[4] Badica P., Kondo T. and Togano K., J. Phys. Soc. Jpn., 74 (2005) 1014.

[5] Yuan H. Q., Agerber D. F., Hayashi N., Badica P., Vanderwelde D., Togano K., Sigrist M. and Salamon M. B., Phys. Rev. Lett., 97 (2006) 017006.

[6] Nishiyama M., Inada Y. and Zheng G.-Q., Phys. Rev. Lett., 98 (2007) 047002.

[7] Yang J. J., Choi Y. J., Oh Y. S., Hogan A., Horibe Y., Kim K., Min B. I. and Cheong S.-W., Phys. Rev. Lett., 108 (2012) 116402.

[8] Pyon S., Kudo K. and Nohara M., J. Phys. Soc. Jpn., 81 (2012) 053701.

[9] Kamitani M., Bhiramy M. S., Arita R., Seki S., Arima T., Tokura Y. and Ishiwata S., Phys. Rev. B, 87 (2013) 180501(R).

[10] Ootsuki D., Wakisaka Y., Pyon S., Kudo K., Nohara M., Arita M., Anzai H., Namatame H., Tanguichi M., Saini N. L. and Mizokawa T., Phys. Rev. B, 86 (2012) 014519.

[11] Fang A. F., Xu G., Dong T., Zheng P. and Wang N. L., Sci. Rep., 2 (2013) 1153.

[12] Ootsuki D., Pyon S., Kudo K., Nohara M., Horio M., Yoshida T., Fujimori A., Arita M., Anzai H., Namatame H., Tanguichi M., Saini N. L. and Mizokawa T., J. Phys. Soc. Jpn., 82 (2013) 093704.

[13] Oh Y. S., Yang J. J., Horibe Y. and Cheong S.-W., Phys. Rev. Lett., 110 (2013) 127209.

[14] Cao H. B., Chakoumakos B. C., Yan J.-Q., Zhou H. D., Custelcean R. and Mandrus D., arXiv:1302.5369.

[15] Monthoux P. and Lonzarich G. G., Phys. Rev. B, 69 (2004) 064517.

[16] Shakeripour H., Petrovic C. and Taillefer L., New J. Phys., 11 (2009) 055065.

[17] Li S. Y., Wu G., Chen X. H. and Taillefer L., Phys. Rev. Lett., 99 (2007) 107001.

[18] Kusmartseva A. F., Sipos B., Berger H., Forró L. and Tutiš E., Phys. Rev. Lett., 103 (2009) 236401.

[19] Sutherland M., Hawthorn D. G., Hill R. W., Ronning F., Wakimoto S., Zhang H., Proust C., Boaknin E., Lupien C., Taillefer L., Liang R. X., Bonn D. A., Hardy W. N., Gagnon R., Hussey N. E., Kimura T., Nohara M. and Takagi H., Phys. Rev. B, 67 (2003) 174520.

[20] Li S. Y., Bonnemaison J.-B., Payeur A., Fournier P., Wang C. H., Chen X. H. and Taillefer L., Phys. Rev. B, 77 (2008) 134501.

[21] Lowell J. and Sousa J., J. Low. Temp. Phys., 3 (1970) 65.

[22] Willis J. and Ginsberg D., Phys. Rev. B, 14 (1976) 1916.

[23] Proust C., Boaknin E., Hill R. W., Taillefer L. and Mackenzie A. P., Phys. Rev. Lett., 89 (2002) 147003.

[24] Suzuki M., Tanatar M. A., Kikugawa N., Mao Z. Q., Maeno Y. and Ishiguro T., Phys. Rev. Lett., 88 (2002) 227004.

[25] Graf M. J., Yip S.-K. and Sauls J. A., Phys. Rev. B, 53 (1996) 15147.

[26] Durst A. C. and Lee P. A., Phys. Rev. B, 62 (2000) 1270.

[27] Hawthorn D. G., Li S. Y., Sutherland M., Boaknin E., Hill R. W., Proust C., Ronning F., Tanatar M. A., Paglione J., Peets D., Liang R., Bonn D. A., Hardy W. N. and Kolesnikov N. N., Phys. Rev. B, 75 (2007) 104518.

[28] Chiao M., Hill R. W., Lupien C., Taillefer L., Lambert P., Gagnon R. and Fournier P., Phys. Rev. B, 62 (2000) 3554.

[29] Volovik G. E., JETP Lett., 58 (1993) 469.

[30] Boaknin E., Tanatar M. A., Paglione J., Hawthorn D., Ronning F., Hill R. W., Sutherland M., Taillefer L., Sonier J., Hayden S. M. and Brill J. W., Phys. Rev. Lett., 90 (2003) 117003.

[31] Hirschfeld P. J., Korshunov M. M. and Mazin I. I., Rep. Prog. Phys., 74 (2011) 124508.