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Phys. Rev. B **93**, 125135 — Published 21 March 2016

DOI: [10.1103/PhysRevB.93.125135](https://doi.org/10.1103/PhysRevB.93.125135)

# Phonon anomalies in some iron-telluride materials

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The detailed temperature dependence of the infrared-active mode in  $\text{Fe}_{1.03}\text{Te}$  ( $T_N \simeq 68$  K) and  $\text{Fe}_{1.13}\text{Te}$  ( $T_N \simeq 56$  K) has been examined, and the position, width, strength, and asymmetry parameter determined using an asymmetric Fano profile superimposed on an electronic background. In both materials the frequency of the mode increases as the temperature is reduced; however, there is also a slight asymmetry in the line shape, indicating that the mode is coupled to either spin or charge excitations. Below  $T_N$  there is an anomalous decrease in frequency and the mode shows little temperature dependence, at the same time becoming more symmetric, suggesting a reduction in spin- or electron-phonon coupling. The frequency of the infrared-active mode and the magnitude of the shift below  $T_N$  are predicted reasonably well by first-principles calculations; however, the predicted splitting of the mode is not observed. In superconducting  $\text{FeTe}_{0.55}\text{Se}_{0.45}$  ( $T_c \simeq 14$  K) the infrared-active  $E_u$  mode displays asymmetric line shape at all temperatures, which is most pronounced between 100–200 K, indicating the presence of either spin- or electron-phonon coupling, which may be a necessary prerequisite for superconductivity in this class of materials.

PACS numbers: 63.20.-e, 71.27.+a, 78.30.-j

## I. INTRODUCTION

The discovery of superconductivity in the iron-based materials has prompted a thorough investigation of this class of materials in an effort to understand the mechanism responsible for the superconductivity, as well as the normal state from which it emerges<sup>1</sup>. While much of this work has focussed on the electronic and magnetic structure, lattice vibrations are also useful for probing electron-phonon coupling and the effects of chemical substitution<sup>2</sup>. Among the iron-based materials, iron telluride is of particular interest because it is thought to be among the most strongly correlated of the iron-chalcogenide materials<sup>3</sup>. At room temperature, the nearly stoichiometric  $\text{Fe}_{1+\delta}\text{Te}$  is a tetragonal, paramagnetic metal that undergoes a first-order structural and magnetic transition to a monoclinic, antiferromagnetic (AFM) metal at  $T_N \simeq 68$  K<sup>4–9</sup>. The introduction of excess iron leads to an increase in the resistivity and the suppression of  $T_N$ <sup>10,11</sup>. The substitution of Se for Te in  $\text{FeTe}_{1-x}\text{Se}_x$  suppresses the structural and magnetic transition and results in a superconducting phase for a broad range of compositions<sup>12,13</sup>, with a maximum critical temperature of  $T_c \simeq 14$  K<sup>14–17</sup>. The optical and transport properties of  $\text{Fe}_{1+\delta}\text{Te}$  and  $\text{FeTe}_{1-x}\text{Se}_x$  have been investigated<sup>17–23</sup> and in addition there have been a number of reports on the vibrational properties of these materials; however, only the Raman-active modes have been studied<sup>24–28</sup>, leaving the infrared-active modes largely unexplored.

In this work we examine the optical properties of  $\text{Fe}_{1.03}\text{Te}$  ( $T_N \simeq 68$  K) and  $\text{Fe}_{1.13}\text{Te}$  ( $T_N \simeq 56$  K) above and below  $T_N$ , as well as  $\text{FeTe}_{0.55}\text{Se}_{0.45}$  ( $T_c \simeq 14$  K), and determine the detailed temperature dependence of the position, width, strength and asymmetry parameter of the infrared-active mode using an asymmetric (Fano) line shape; in several cases the electronic properties of these

materials have been previously discussed by us<sup>19–21</sup>. In both  $\text{Fe}_{1+\delta}\text{Te}$  materials the infrared-active mode displays a slight asymmetry, suggesting either spin- or electron-phonon coupling, and increases in frequency with decreasing temperature, as expected for an anharmonic decay process; below  $T_N$  the mode undergoes an anomalous decrease in frequency at  $T_N$  and the asymmetry parameter decreases, indicating reduced coupling. The vibrational frequencies and atomic intensities have been calculated at the center of the Brillouin zone from first-principles methods for both the high-temperature tetragonal phase, and the low-temperature monoclinic phase. In the latter case, spin ordering is shown in some cases to alter the character of the mode resulting in a large predicted frequency shift below  $T_N$ . Interestingly, the predicted splitting of the infrared-active  $E_u$  mode below  $T_N$  is not observed. In  $\text{FeTe}_{0.55}\text{Se}_{0.45}$  the  $E_u$  mode displays an asymmetry at all temperatures, which is most pronounced between 100–200 K. The asymmetric profile is a signature of either spin- or electron-phonon coupling, which may be necessary condition for superconductivity in this class of materials.

## II. EXPERIMENT

Single crystals of  $\text{Fe}_{1.03}\text{Te}$  and  $\text{Fe}_{1.13}\text{Te}$  with good cleavage planes (001) have been grown by a unidirectional solidification process; these crystals undergo structural and magnetic transitions<sup>4–9</sup> at  $T_N \simeq 68$  K and  $\simeq 56$  K, respectively. Single crystals with a nominal composition of  $\text{FeTe}_{0.55}\text{Se}_{0.45}$  have also been grown using this similar process with  $T_c \simeq 14$  K and a transition width of  $\simeq 1$  K. The reflectance of these materials with the light polarized in the Fe–As ( $a$ - $b$ ) planes has been measured over a wide frequency range ( $\simeq 2$  meV to over 3 eV) using an *in situ* evaporation technique<sup>29</sup> for a wide range

of temperatures. The complex optical properties have been calculated from a Kramers-Kronig analysis of the reflectance<sup>30</sup>. The large-scale structure and temperature dependence of the optical conductivity of some of these materials have been previously discussed<sup>19–21</sup>.

### III. RESULTS AND DISCUSSION

Lattice vibrations in solids are often described using a Lorentzian oscillator superimposed on a linear (or polynomial) background. The complex dielectric function  $\tilde{\epsilon} = \epsilon_1 + i\epsilon_2$  for the Lorentz oscillator is

$$\tilde{\epsilon}(\omega) = \frac{\Omega_0^2}{\omega_0^2 - \omega^2 - i\gamma_0\omega}, \quad (1)$$

where  $\omega_0$ ,  $\gamma_0$  and  $\Omega_0$  are the position, width, and strength of the vibration, respectively. The complex conductivity is  $\tilde{\sigma}(\omega) = \sigma_1 + i\sigma_2 = -2\pi i\omega[\tilde{\epsilon}(\omega) - \epsilon_\infty]/Z_0$ , where  $\epsilon_\infty$  is a high-frequency contribution to the real part, and  $Z_0 \simeq 377 \Omega$  is the impedance of free space, yielding units for the conductivity of  $\Omega^{-1}\text{cm}^{-1}$ . The real part of the optical conductivity for the oscillator may then be written as

$$\sigma_1(\omega) = \frac{2\pi}{Z_0} \left[ \frac{\gamma_0\omega^2\Omega_0^2}{(\omega_0^2 - \omega^2)^2 + \gamma_0^2\omega^2} \right]. \quad (2)$$

While this approach accurately describes a symmetric line shape, the coupling of lattice vibrations to either the spin or electronic background may lead to an asymmetric line shape, often referred to as the Fano profile<sup>31</sup>. The Fano line shape is written as  $\sigma_1(\omega) = A[(x+q)^2/(1+x^2)]$ , where  $A$  is a constant,  $x = (\omega - \omega_0)/\gamma_0$ , and the asymmetry is described by the dimensionless parameter  $1/q^2$ . A more useful expression,

$$\sigma_1(\omega) = \frac{2\pi}{Z_0} \frac{\Omega_0^2}{\gamma_0} \frac{(q^2 + 4qx - 1)}{q^2(1 + 4x^2)}, \quad (3)$$

incorporates the oscillator strength; however, this from contains no information about  $\sigma_2(\omega)$  and does not satisfy the Kramers-Kronig relation. To resolve this issue a phenomenological dielectric function for a Fano-shaped Lorentz oscillator is employed in which is Kramers-Kronig consistent<sup>32,33</sup>,

$$\tilde{\epsilon}(\omega) = \frac{\Omega_0^2}{\omega_0^2 - \omega^2 - i\gamma_0\omega} \left( 1 + i\frac{\omega_q}{\omega} \right)^2 + \left( \frac{\Omega_0\omega_q}{\omega_0\omega} \right)^2 \quad (4)$$

where  $1/q = \omega_q/\omega_0$ . Unlike the simple Fano profile, a complex conductivity may be calculated which satisfies  $\tilde{\sigma}^*(\omega) = \tilde{\sigma}(-\omega)$ . The real part of the optical conductivity is then

$$\sigma_1(\omega) = \frac{2\pi}{Z_0} \frac{\Omega_0^2}{\gamma_0} \frac{[\gamma_0\omega^2 - 2(\omega^2\omega_0 - \omega_0^3)/q - \gamma_0\omega_0^2/q^2]}{(\omega^2 - \omega_0^2)^2 + \gamma_0^2\omega^2}. \quad (5)$$

In the  $1/q^2 \rightarrow 0$  limit a symmetric Lorentzian is recovered [Eq. (2)]; however, as  $1/q^2$  increases the line shape becomes increasingly asymmetric. We adopt this latter form for asymmetric line shapes to fit the infrared-active vibrations in  $\text{Fe}_{1+\delta}\text{Te}$  and  $\text{FeTe}_{0.55}\text{Se}_{0.45}$ .

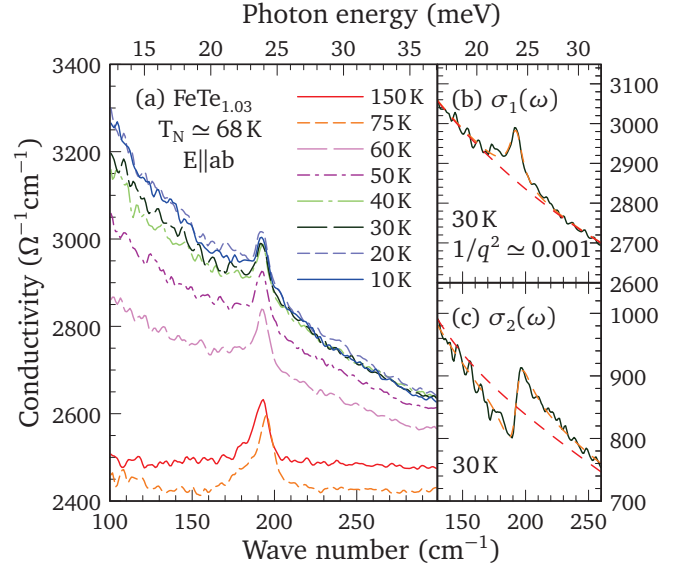


Figure 1. (Color online) (a) The real part of the optical conductivity of  $\text{Fe}_{1.03}\text{Te}$  above and below  $T_N \simeq 68 \text{ K}$  in the region of the infrared-active mode at  $\simeq 196 \text{ cm}^{-1}$ . There is no evidence that this mode splits below  $T_N$ . (b) Fano oscillator fit (dashed line) to  $\sigma_1(\omega)$  at 30 K with Drude-Lorentz background (long-dashed line), and (c) the simultaneous fit to  $\sigma_2(\omega)$ ; note the almost perfectly symmetric profile.

#### A. $\text{Fe}_{1+\delta}\text{Te}$

##### 1. Vibrational properties

The real part of the optical conductivity of  $\text{Fe}_{1.03}\text{Te}$  in the region of the infrared-active phonon is shown in Fig. 1(a). The rapid increase in the low-frequency conductivity below  $T_N$  has been attributed to the closing of the pseudogap below  $T_N$  leading to an increase in the free-carrier concentration<sup>19</sup>. In the room-temperature tetragonal phase, the  $P4/nmm$  space group with two formula units per unit cell yields the irreducible vibrational representation<sup>24</sup>

$$\Gamma_{\text{HT}} = A_{1g} + B_{1g} + 2E_g + A_{2u} + E_u.$$

Two infrared-active modes are expected, a doubly-degenerate  $E_u$  mode active in the planes, and an  $A_{2u}$  mode active along the  $c$  axis. In the low-temperature monoclinic phase, the  $P2_1/m$  space group with two formula units per unit cell yields the irreducible vibrational representation<sup>26</sup>

$$\Gamma_{\text{LT}} = 4A_g + 2B_g + A_u + 2B_u,$$

where the  $A_u$  mode is infrared active along the  $c$  axis and the  $E_u$  modes splits into two  $B_u$  modes that are active in the  $a$ - $b$  plane.

The infrared-active mode is fit simultaneously to the real and imaginary parts of the complex conductivity using the asymmetric line shape described by the dielectric function in Eq. (4) superimposed on a free-carrier

(Drude) response in combination with several symmetric Lorentz oscillators that describe the mid-infrared response. Because the iron-telluride compounds are multi-band materials<sup>34</sup> the minimal description of the dielectric function usually consists of at least two Drude components<sup>35</sup>,

$$\tilde{\epsilon}(\omega) = \epsilon_{\infty} - \sum_{j=1}^2 \frac{\omega_{p,D;j}^2}{\omega^2 + i\omega/\tau_{D,j}} \quad (6)$$

where  $\epsilon_{\infty}$  has been previously defined,  $\omega_{p,D;j}^2 = 4\pi n_j e^2 / m_j^*$  and  $1/\tau_{D,j}$  are the square of the plasma frequency and scattering rate for the delocalized (Drude) carriers, respectively, in the  $j$ th band, and  $n_j$  and  $m_j^*$  are the carrier concentration and effective mass. The values for  $\omega_{p,D;j}$  and  $1/\tau_{D,j}$  and the mid-infrared oscillators are initially fit to the complex conductivity up to at least 1 eV using a non-linear least-squares technique. Fits to the infrared-active mode encompass only a narrow region around the mode (typically  $\pm 50 \text{ cm}^{-1}$ ), and as a result in addition to the vibrational parameters, only the Drude components and the width and strength of the lowest mid-infrared oscillator are allowed to vary, while the other parameters remain fixed. The temperature dependence of  $\omega_{p,D;1}$ ,  $1/\tau_{D,j}$  and the mid-infrared oscillators in these materials have been previously discussed<sup>19,21</sup>. As Figs. 1(b) and 1(c) illustrate, both the background and the line shape are reproduced quite well using this approach.

The temperature dependence of  $\omega_0$ ,  $\gamma_0$ ,  $\Omega_0$ , and the asymmetry parameter  $1/q^2$  of the infrared-active mode in  $\text{Fe}_{1.03}\text{Te}$  are shown in Fig. 2. The frequency of the mode increases (hardens) in an almost linear fashion with

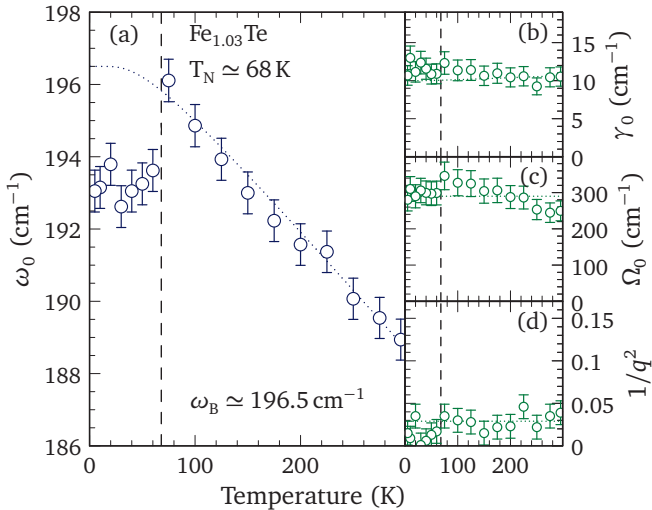


Figure 2. (Color online) The temperature dependence of the (a) infrared-active mode compared with the anharmonic decay model (dotted line), (b) linewidth and predicted dependence (dotted line), (c) oscillator strength, and (d) asymmetry parameter in  $\text{Fe}_{1.03}\text{Te}$  ( $T_N \simeq 68 \text{ K}$ ); the dotted lines in the last two panels are drawn as a guide to the eye.

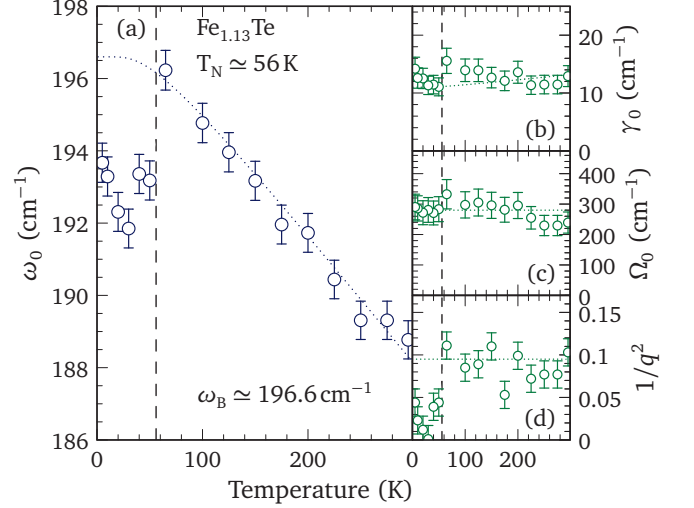


Figure 3. (Color online) The temperature dependence of the (a) infrared-active mode compared with the anharmonic decay model (dotted line), (b) linewidth and predicted dependence (dotted line), (c) oscillator strength, and (d) asymmetry parameter in  $\text{Fe}_{1.13}\text{Te}$  ( $T_N \simeq 56 \text{ K}$ ); the dotted lines in the last two panels are drawn as a guide to the eye.

decreasing temperature. At  $T_N$  there is an anomalous decrease (softening) in the frequency of the mode followed by a very weak temperature dependence; there is no evidence for a new infrared active  $B_u$  mode appearing below  $T_N$ . This is similar to the temperature dependence of the frequency of the  $B_{1g}$  mode in these materials<sup>26–28</sup>. The width of this mode shows very little temperature dependence either above or below  $T_N$ . While strength of this mode may increase slightly as the temperature is lowered to  $T_N$ , it is also possible that, within error, it is constant; below  $T_N$  there is very little change. Above  $T_N$ , the asymmetry parameter  $1/q^2 \simeq 0.03$ , suggesting weak spin- or electron-phonon coupling<sup>2</sup>. Below  $T_N$  the asymmetry parameter  $1/q^2$  decreases and the line shape becomes almost perfectly symmetric; this inferred reduction in the electron-phonon coupling may be related to the fact that in the AFM state the spin fluctuations are observed to decrease<sup>36</sup> and most signs of correlations disappear<sup>37</sup>.

The temperature dependence of the infrared-active mode in  $\text{Fe}_{1.13}\text{Te}$  has been fit using procedure described above, and the results are shown in Fig. 3 (the temperature dependence of the optical conductivity and the fit at 30 K is shown in Fig. S1 in the Supplemental Material). The behavior of  $\omega_0$ ,  $\gamma_0$ , and  $\Omega_0$  are quite similar to those of  $\text{Fe}_{1.03}\text{Te}$ , although there does tend to be a bit more scatter in the data points. However, above  $T_N$  the asymmetry parameter  $1/q^2 \simeq 0.1$ , suggesting that the spin- or electron-phonon coupling is stronger in this material; as in  $\text{Fe}_{1.03}\text{Te}$ , below  $T_N$  the asymmetry parameter  $1/q^2$  decreases, suggesting a decrease in correlations and a commensurate reduction in the coupling of the phonons to either the spin or charge.

Above  $T_N$  the behavior of  $\omega_0$  can be described quite

Table I. Calculated frequencies and Fe atomic intensities in the high- and low-temperature phases of FeTe for the zone center modes. For each phase, the results based on the relaxed and experimental (shown in parenthesis) unit cell parameters are considered; for the low-temperature phase, non-spin polarized (NSP) and spin polarized (SP) calculations are performed. All units are in  $\text{cm}^{-1}$  unless otherwise indicated.

$P4/nmm$ (HT)		Atomic intensity	$P2_1/m$ (LT)		Atomic intensity	Atomic intensity		Expt. <sup>a</sup>
Symmetry	Theory	Fe	Symmetry	Theory-NSP	Fe	Theory-SP	Fe	
$E_g$	98 (77)	0.11 (0.12)	$A_g$	98 (57)	0.10 (0.17)	84 (81)	0.13 (0.11)	160
$A_{1g}$	170 (141)	0.00 (0.00)	$B_g$	99 (72)	0.10 (0.13)	85 (79)	0.13 (0.08)	
$E_u$	249 (203)	0.70 (0.70)	$A_g$	173 (137)	0.00 (0.00)	145 (139)	0.15 (0.00)	
$B_{1g}$	272 (221)	1.00 (1.00)	$B_u$	245 (203)	0.70 (0.70)	241 (195)	0.70 (0.70)	196
$E_g$	292 (223)	0.89 (0.88)	$B_u$	300 (254)	0.70 (0.70)	306 (267)	0.70 (0.70)	
$A_{2u}$	304 (257)	0.70 (0.69)	$A_g$	270 (210)	1.00 (0.89)	167 (—)	0.85 (—)	203
			$A_g$	289 (226)	0.90 (0.92)	287 (243)	0.87 (0.89)	
			$B_g$	292 (215)	0.90 (0.87)	288 (245)	0.87 (0.92)	
			$A_u$	248 (196)	0.70 (0.70)	244 (189)	0.70 (0.70)	

<sup>a</sup> The Raman results are taken from Ref. 27; infrared is from this work.

well assuming a symmetric anharmonic decay into two acoustic modes with identical frequencies and opposite momenta<sup>38,39</sup>. A slightly modified functional form is employed here,

$$\omega_0(T) = \omega_B \left[ 1 - \frac{2C}{e^x - 1} \right], \quad (7)$$

$$\gamma_0(T) = \Gamma_0 \left[ 1 + \frac{2\Gamma}{e^x - 1} \right], \quad (8)$$

where  $\omega_B$  is the bare phonon frequency,  $\Gamma_0$  is a residual line width,  $C$  and  $\Gamma$  are constants, and  $x = \hbar\omega_B/(2k_B T)$ ; an advantage of this form is that the bare phonon frequency (residual line width) is recovered in the  $T \rightarrow 0$  limit. The model fits to the  $T > T_N$  data are indicated by the dotted lines in Figs. 2 and 3, yielding nearly identical values of  $\omega_B \simeq 196.5$  and  $196.6 \text{ cm}^{-1}$  for  $\text{Fe}_{1.03}\text{Te}$  and  $\text{Fe}_{1.13}\text{Te}$ , respectively.

The behavior of the infrared-active mode at  $T_N$  is also reminiscent of the  $E_u$  mode in  $\text{BaFe}_2\text{As}_2$  which also experiences an anomalous softening below the magnetic and structural transition at  $T_N \simeq 138 \text{ K}$ <sup>40</sup>; however, in that material the oscillator strength nearly doubles below  $T_N$ , whereas in  $\text{Fe}_{1+\delta}\text{Te}$  it remains more or less unchanged.

## 2. Lattice dynamics

The frequencies and atomic intensities (square of the vibrational amplitude of each atom that sum to unity) for the zone-center phonons have been calculated from first principles using the so-called frozen-phonon method, the details of which are described in the Appendix. The results for the high-temperature tetragonal phase using the relaxed (experimental) unit cell parameters (Table II in the Appendix) are shown in Table I and compared with experimental results; these results are in good agreement with other calculations<sup>24,25</sup>. It has been remarked that in

this material the frequencies calculated using the experimental unit cell parameters are in better agreement with the experimentally-observed frequencies than the values determined from a relaxed unit cell<sup>24</sup>, and that is indeed the case here. In the relaxed unit cell the Fe–Te bond is shorter, which has the effect of scaling up all of the frequencies by 20 – 30%; however, we note that the character of the atomic intensities is essentially unaffected. The  $A_{1g}$  and  $B_{1g}$  Raman modes are calculated using the relaxed (experimental) unit cell parameters at 170 (141) and 272 (221)  $\text{cm}^{-1}$ , respectively, are pure Te and Fe vibrations, whereas the  $E_u$  mode is a mixture of the two. Using the experimental unit cell, the calculated frequency for the  $E_u$  mode of 203  $\text{cm}^{-1}$  is quite close to the measured value of  $\simeq 196 \text{ cm}^{-1}$ , as are the values for the  $A_{1g}$  and  $B_{1g}$  modes<sup>27</sup>.

The frequencies and atomic intensities for the low-temperature monoclinic phase have also been calculated using the relaxed (experimental) unit cell parameters (Table III in the Appendix), and the results listed in Table I. As previously noted, the lower symmetry results in the splitting of the doubly-degenerate  $E_g \rightarrow A_g + B_g$  and  $E_u \rightarrow 2B_u$  modes. In the non-magnetic calculation (non-spin polarized), the calculated frequencies for the relaxed and experimental unit cells match quite well with their counterparts in the tetragonal phase, and the character of the modes does not change substantially. In the relaxed unit cell the splitting of the upper and lower  $E_g$  modes is negligible ( $\simeq 2 \text{ cm}^{-1}$ ); however in the experimental unit cell this splitting is much larger,  $\simeq 20 \text{ cm}^{-1}$ . In both the experimental and relaxed geometries, the  $E_u$  mode is predicted to split by  $\simeq 50 \text{ cm}^{-1}$ . A simple empirical force-constant model<sup>41</sup> reproduces the size of the splitting below  $T_N$  and indicates that the two  $B_u$  modes should have roughly the same strength. Using a splitting of  $\simeq 50 \text{ cm}^{-1}$ , a new mode might be expected at  $\simeq 240 \text{ cm}^{-1}$ ; however, there does not appear to be any experimental evidence for a new mode below  $T_N$  [Fig. 1(a) and Fig. S1(a) in the Supplementary Material].

The non-spin polarized calculation does not take into



consideration the antiferromagnetic ground state of this material<sup>4</sup>. Accordingly, the spins on the two Fe sites have been placed in an antiferromagnetic configuration while the Te atoms remain non-magnetic; the frequencies and atomic characters are then calculated for a spin-polarized case using both the relaxed and experimentally-determined unit cell parameters, and the results listed in Table I. The introduction of magnetic order has a significant effect on some modes changing both the frequency and the character of the vibration; however, we note that the  $E_u$ , upper  $E_g$  and  $A_{2u}$  modes experience only a small effect. The small predicted downward shift in the lower  $B_u$  mode of  $\simeq 4 - 7 \text{ cm}^{-1}$  for the spin-polarized case is in good agreement with the experimentally-observed softening of this mode of between  $2 - 4 \text{ cm}^{-1}$  in  $\text{Fe}_{1+\delta}\text{Te}$  below  $T_N$ .

### B. $\text{FeTe}_{0.55}\text{Se}_{0.45}$

The real part of the optical conductivity of superconducting  $\text{FeTe}_{0.55}\text{Se}_{0.45}$  in the region of the infrared-active phonon at  $\simeq 213 \text{ cm}^{-1}$  is shown in Fig. 4(a)<sup>21</sup>. The substitution of Se for Te in this material has the effect of shifting the frequency of the  $E_u$  mode up about 10%. The temperature dependence of the  $E_u$  mode has been fit to the complex conductivity using the asymmetric Fano line shape in Eq. (4) superimposed on a Drude-Lorentz background; as Figs. 4(b) and 4(c) illustrate, both the background and the line shape are reproduced quite well. In several cases, it appears as though there is a weak

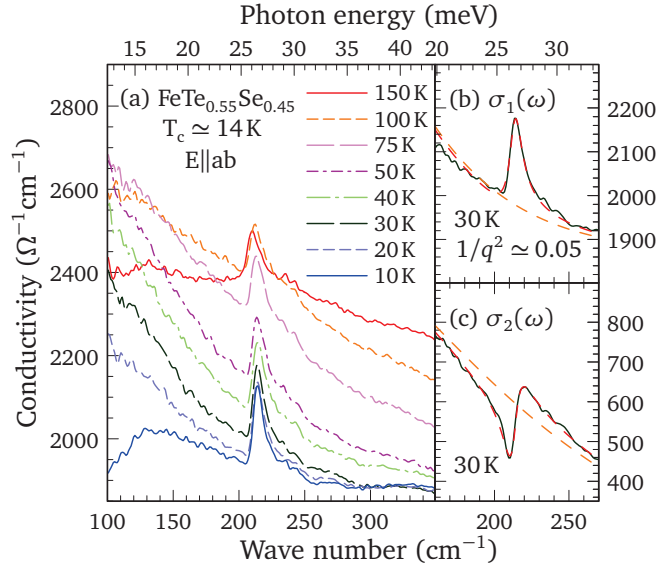


Figure 4. (Color online) (a) The real part of the optical conductivity of  $\text{FeTe}_{0.55}\text{Se}_{0.45}$  above and below  $T_c \simeq 14 \text{ K}$  in the region of the infrared-active  $E_u$  mode at  $\simeq 213 \text{ cm}^{-1}$ . (b) Fano oscillator fit (dashed line) to  $\sigma_1(\omega)$  at 30 K with Drude-Lorentz background (long-dashed line), and (c) the simultaneous fit to  $\sigma_2(\omega)$ .

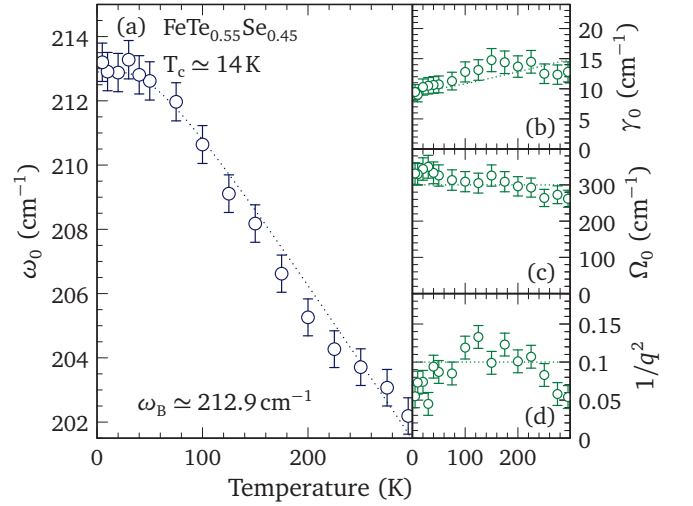


Figure 5. (Color online) The temperature dependence of the (a) infrared-active  $E_u$  mode compared with the anharmonic decay model (dotted line), (b) line width and predicted dependence (dotted line), (c) oscillator strength, and (d) asymmetry parameter in  $\text{FeTe}_{0.55}\text{Se}_{0.45}$  ( $T_c = 14 \text{ K}$ ); the dotted lines in the last two panels are drawn as a guide to the eye.

sideband at  $\simeq 240 \text{ cm}^{-1}$ ; however, no new mode is predicted in this material, suggesting that such features are the result of an asymmetric line shape.

The temperature dependence of the frequency and line width of the mode, shown in Figs. 5(a) and 5(b), respectively, both agree quite well with the anharmonic decay model. The strength of the mode shown in Figs. 5(c) may be increasing somewhat with decreasing temperature, but within the accuracy of the fits it is also possible that it is temperature independent. The asymmetry parameter is small at room temperature,  $1/q^2 \simeq 0.05$ , but increases with temperature, reaching a maximum of  $1/q^2 \simeq 0.13$  at  $\simeq 125 \text{ K}$  before decreasing at low temperature, as shown in Fig. 5(d). This would tend to indicate that the spin- or electron-phonon coupling of the  $E_u$  mode is becoming significant, and indeed it is considerably larger than what is observed in related iron-pnictide materials<sup>2</sup>. The decrease in  $1/q^2$  at low temperature is correlated with the rapid decrease in the Hall coefficient below  $100 \text{ K}$ <sup>42</sup>, suggesting that the electron-phonon coupling is related to a specific carrier pocket in this material. While electron-phonon coupling is present in this material at low temperatures, it is thought that the electron-phonon coupling constant  $\lambda$  in this material is quite small ( $\lambda \simeq 0.01$ )<sup>43</sup>. The low values for  $\lambda$  do not support the high  $T_c$ 's observed in iron-based superconductors, implying that these materials are not phonon-mediated superconductors<sup>44</sup>. However, a detailed analysis of the five-orbital Hubbard-Holstein model for iron pnictides has demonstrated that a relatively small electron-phonon interaction arising from the Fe-ion oscillation can induce substantial  $d$ -orbital fluctuations<sup>45,46</sup>, as well as the anion electronic polarization<sup>47</sup>. These orbital fluctuations

give rise to a strong pairing interaction and mediate an  $s_{++}$  superconducting state. In addition, a large Fe isotope effect has been reported in  $\text{SmFeAsO}_{1-x}\text{Fe}_x$  and  $\text{Ba}_{1-x}\text{K}_x\text{Fe}_2\text{As}_2$ <sup>48</sup>, indicating that electron-phonon coupling likely plays some role in the pairing mechanism. Our experimental results in the  $\text{Fe}_{1+y}\text{Te}_{1-x}\text{Se}_x$  system, in combination with the above theoretical and experimental studies, suggests that electron-phonon coupling may be a necessary prerequisite for superconductivity in this class of materials.

#### IV. CONCLUSIONS

The detailed temperature dependence of the infrared-active infrared-active mode has been fit in  $\text{Fe}_{1.03}\text{Te}$  and  $\text{Fe}_{1.13}\text{Te}$  using an asymmetric Fano line shape superimposed on a background of Drude free carriers and symmetric Lorentz oscillators. In both materials, above  $T_N$  the frequency increases as the temperature is lowered; this behavior is described quite well by an anharmonic decay process. Below  $T_N$ , there is an abrupt softening of this mode in both materials, below which it displays little temperature dependence. While frequency of the mode and the size of the shift below  $T_N$  are predicted reasonably well from first-principles calculations (with some caveats), the predicted splitting of the infrared-active mode is not observed. In both materials, the phonon has a slightly asymmetric line shape above  $T_N$ , suggesting that there is coupling of the phonons to either spin or charge; however,  $1/q^2$  decreases and the profile becomes more symmetric below  $T_N$ , suggesting this coupling is reduced. The  $E_u$  mode in superconducting  $\text{FeTe}_{0.55}\text{Se}_{0.45}$  has been fit in a similar fashion and displays an asymmetry at all temperatures, which is most pronounced between 100 – 200 K. The temperature dependence of this mode is once again described quite well by an anharmonic decay process over the entire temperature range. The asymmetric nature of the  $E_u$  mode is an indication of spin- or electron-phonon coupling to the electronic background, which may be a necessary prerequisite for superconductivity in the iron-based materials.

#### ACKNOWLEDGMENTS

We would like to acknowledge helpful discussions with T. Birol. We are indebted to S. V. Dordevic for referring us to a phenomenological Fano line shape, developed by A. Damascelli and A. Kuzmenko. This work is supported by the Office of Science, U.S. Department of Energy under Contract No. DE-SC0012704. J.S. and R.D.Z. are supported by the Center for Emergent Superconductivity, an Energy Frontier Research Center funded by the U.S. Department of Energy, Office of Science.

Table II. The experimental and theoretical lattice constants and atomic fractional coordinates for the relaxed structure of FeTe for the high-temperature tetragonal  $P4/nmm$  space group. The fractional coordinate for the Fe atom is  $(\frac{3}{4}\frac{1}{4}0)$ . The Fe-Te bond lengths are also shown, while the Fe-Fe and Te-Te bond lengths, 2.6960 Å and 2.8126 Å, respectively, are essentially identical for the two structures.

	Experiment <sup>a</sup>	GGA
$a$ (Å)	3.8127	3.8126
$c$ (Å)	6.2450	6.2448
Te( $\frac{1}{4}\frac{1}{4}z$ )	0.2813	0.2569
Fe-Te (Å)	2.5923	2.4915

<sup>a</sup> Ref 4 [Supplemental Material. Table I, Section (a)].

#### Appendix: Lattice dynamics

##### 1. $P4/nmm$ (HT)

The total energy of FeTe was calculated with the generalized gradient approximation (GGA) using the full-potential linearized augmented plane-wave (FP-LAPW) method<sup>49</sup> with local-orbital extensions<sup>50</sup> in the WIEN2k implementation<sup>51</sup>. An examination of different Monkhorst-Pack  $k$ -point meshes indicated that 175  $k$  points ( $6 \times 6 \times 4$  mesh) and  $R_{mt}k_{max} = 7.5$  was sufficient for good energy convergence. The geometry of the unit cell was refined through an iterative process whereby the volume was optimized with respect to the total energy while the  $c/a$  ratio remained fixed. The atomic fractional coordinates were then relaxed with respect to the total force, typically resulting in residual forces of less than 1 mRy/a.u. This procedure was repeated until no further improvement was obtained. A comparison of the experimental and calculated (relaxed) unit cell parameters are shown in Table. II. Note that while the unit cell parameters are nearly identical to the experimentally-determined values, the position of the Te atom has shifted somewhat.

The phonons have been determined using the direct method. To determine the phonons at the zone center, a  $1 \times 1 \times 1$  supercell is sufficient. To obtain a complete set of Hellmann-Feynman forces, a total of 4 independent displacements are required; because there are always some residual forces at the atomic sites we have considered symmetric displacements, which doubles this number, resulting in a total of 8 atomic displacements. In this case, displacement amplitudes of 0.03 Å were used. The calculations have converged when the successive changes for the forces on each atom are less than 0.1 mRy/au. The residual forces are collected for each set of symmetric displacements and a list of the Hellmann-Feynman forces are generated. Using the program PHONON<sup>52</sup> the cumulative force constants deconvoluted from the Hellmann-Feynman forces are introduced into the dynamical matrix, which is then diagonalized in order to obtain the phonon frequencies; the atomic intensities are further cal-

culated to describe the character of the vibration. The results are shown in Table I.

## 2. $P2_1/m$ (LT)

The procedure that was used to calculate the geometry of the room-temperature tetragonal phase has been repeated for the low-temperature monoclinic phase; the same  $k$  point mesh and  $R_{mt}k_{max}$  were used. The optimized unit cell parameters are compared with the experimental values and shown in Table III. Once again, the calculated unit cell parameters are nearly identical to the experimental values; however the Fe and Te atoms have undergone small shifts resulting in a decrease of the Fe–Te bond lengths. Interestingly, the Fe–Fe bond modulation observed experimentally is not captured in the non-spin polarized GGA result in which the Fe–Fe bonds are almost identical. To calculate the zone-center phonons, a  $1 \times 1 \times 1$  supercell is employed. In the monoclinic phase a total of 6 independent displacements are required; symmetric displacements double this number to 12. Displacement amplitudes of 0.03 Å were used.

Two different cases have been considered for this symmetry. In the first case the system is assumed to be non-magnetic and the spins on the Fe and Te sites are ignored and a non-spin polarized (NSP) calculation performed. The calculated phonon frequencies and atomic intensities which are listed in Table I. In the second case the spins on the Fe sites are specified to be in an antiferromagnetic configuration and the Te atoms are assumed to be non-magnetic and a spin-polarized calculation is performed. If the fractional coordinates are relaxed in this case, position of the Te atom is now closer to the experimental value, and a Fe–Fe bond modulation is now observed (Table III). However, given that this structure appears to be intermediate between the non-spin polarized and the experimental unit cell parameters, we have not performed a phonon calculation. Using the previously established criteria for convergence, we find the magnetic moments are about  $1.7 \mu_B/\text{Fe}$  for the (non-spin polarized) relaxed unit cell, which is very close to the expected value of  $2 \mu_B/\text{Fe}$  observed in other work<sup>53</sup>. When the experimental unit cell is used, the magnetic moments increase to  $2.2 \mu_B/\text{Fe}$ ; however, this results in a calculated frequency of less than zero for one of the  $A_g$  modes, this unphysical value indicates that within the scope of this calculation the structure is unstable.



Table III. The experimental and theoretical lattice constants and atomic fractional coordinates for the relaxed structure of FeTe for the low-temperature monoclinic  $P2_1/m$  space group. The Fe–Fe and Fe–Te bond lengths are also shown. Note that the GGA calculation does not reproduce the experimentally-observed Fe–Fe bond modulation.

	Experiment <sup>a</sup>		GGA (NSP)		GGA (SP)	
$a$ (Å)	3.8346		3.8391		3.8391	
$b$ (Å)	3.7836		3.7897		3.7897	
$c$ (Å)	6.2567		6.2649		6.2649	
$\beta$ (°)	89.21		89.24		89.24	
Fe( $x \frac{1}{4} z$ )	0.7612	0.0033	0.7501	0.0013	0.7507	0.0011
Te( $x \frac{1}{4} z$ )	0.2586	0.2822	0.2448	0.2559	0.2454	0.2717
Fe–Fe (Å)	2.633/2.756		2.697/2.698		2.693/2.701	
Fe–Te (Å)	2.582/2.603/2.604		2.487/2.495/2.496		2.552/2.559/2.562	

<sup>a</sup> Ref 8.

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