Photoemission spectroscopy and the unusually robust one-dimensional physics of lithium purple bronze

L Dudy$^{1,2}$, J D Denlinger$^3$, J W Allen$^1$, F Wang$^{1,4}$, J He$^5$, D Hitchcock$^5$, A Sekiyama$^6$ and S Suga$^6$

$^1$Randall Laboratory, University of Michigan, Ann Arbor, MI 48109, USA
$^2$Physikalisches Institut, Universität Würzburg, D-97074 Würzburg, Germany
$^3$Advanced Light Source, Lawrence Berkeley National Laboratory, Berkeley, CA 94270, USA
$^4$BTG Pactual, 601 Lexington Avenue—57th floor, New York, NY 10022, USA
$^5$Department of Physics and Astronomy, Clemson University, Clemson, SC 29534, USA
$^6$Department of Material Physics, Graduate School of Engineering Science, Osaka University, 1–3 Machikaneyama, Toyonaka, Osaka 560-8531, Japan

E-mail: lenart.dudy@physik.uni-wuerzburg.de

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Abstract

Temperature-dependent photoemission spectroscopy in Li$_{0.9}$Mo$_6$O$_{17}$ contributes to evidence for one-dimensional (1D) physics that is unusually robust. Three generic characteristics of the Luttinger liquid are observed: power law behavior of the $k$-integrated spectral function down to temperatures just above the superconducting transition, $k$-resolved lineshapes that show holon and spinon features, and quantum critical (QC) scaling in the lineshapes. Departures of the lineshapes and the scaling from expectations in the Tomonaga–Luttinger model can be partially described by a phenomenological momentum broadening that is presented and discussed. The possibility that some form of 1D physics obtains even down to the superconducting transition temperature is assessed.

(Some figures may appear in colour only in the online journal)

1. Introduction

Li$_{0.9}$Mo$_6$O$_{17}$, the so-called lithium purple bronze (hereafter called LiPB), has emerged as a paradigm quasi-one-dimensional (quasi-1D) material for studying the stability of Luttinger liquid chains against single-particle electron hopping $t_\perp$ between the chains [1]. The usual theoretical understanding [2, 3] and empirical experience is that quasi-1D materials can display such 1D behavior at sufficiently elevated temperatures but that, with decreasing temperature ($T$), $t_\perp$ leads to a crossover to some kind of 3D behavior, e.g. a Fermi liquid, or an ordered state such as a charge density wave (CDW) or spin density wave (SDW). Because of the excellent Fermi surface (FS) nesting that is typical of quasi-1D materials, density waves are especially likely. Focusing on photoemission spectroscopy but considering other data as well, this paper explores the possibility that LiPB, in spite of excellent FS nesting, does not develop a density wave and crosses over to 3D behavior only when it becomes a superconductor (SC). If such is true then the SC could well be unconventional, as suggested in a recent paper [4] reporting that the superconducting critical field is far above the Pauli limit.

LiPB is a material with overall 3D bonding but whose valence band electronic structure is derived from two parallel zig-zag Mo–O chains per unit cell. The chains lie in well separated planes and receive electrons donated by Li ions located out of the planes. Early tight binding [5] and more recent LDA [6] band calculations agree that the Mo–O orbitals of the chains give rise to four bands, two of which (A, B) are always below the Fermi energy ($E_F$) and two of which (C, D) cross $E_F$. The resulting FS is quasi-1D but the LDA calculations also show some splitting and warping due to single-particle hopping $t_\perp$ between chains both within a unit cell and in neighboring unit cells. Angle resolved photoemission spectroscopy (ARPES) observes all
four bands with the general behavior that is predicted in LDA calculations [7–9], including the quasi-1D FS [7]. Figures presented in later discussions of ARPES data show the four bands (figure 2(a) in section 3), and the FS and Brillouin zone (figure 4(c) in section 4). The magnitude of $t_\perp$ is discussed in section 6 of the paper.

The physical properties of LiPB were first investigated in the 1980s [10–20]. Electrical transport [10] was found to be highly anisotropic and superconductivity (SC) [10, 11, 15, 16, 18] was observed below $T_{SC} = 1.9$ K, albeit not in all samples [19] and with a variation of $T_{SC}$ [10, 11, 15, 16, 19] that might involve the Li stoichiometry [21] or disorder [15]. More recent studies [22–27] confirm the anisotropic transport but with differences in the absolute values of the resistivities along various axes, and also in the magnitudes of the anisotropy ratios. These differences may reflect variation in samples and in the method [23] of the measurement.

2. Resistivity upturn and evidence against a density wave

The early experimental work found that the resistivity is metallic with decreasing $T$ only down to a minimum at $T_{min} \approx 26$ K, below which there is an abrupt upturn. This finding has been robust over time. Figure 1(a) shows early data [10] and more recent data [25, 26] for the upturn as measured along the 1D direction. An upturn at the same temperature is also seen for the two directions perpendicular to the chains. There is yet no consensus as to the origin of this upturn because each proposed explanation requires that one or another piece of data be ignored, be given less weight in the argument, or be left unexplained. Early discussions ascribed the upturn to Anderson localization due to disorder [12–15] or to a single-particle gap, most likely due to CDW formation [10, 11, 18–20]. Both direct and nuanced lines of argument [10–12, 14, 13, 18–20, 22–27] for and against these two hypotheses have been presented many times in the transport literature. Reference [22] is the most comprehensive recent overview of the issue.

The disorder hypothesis was proposed [12, 14] because a number of important expected signatures of a density wave are missing in LiPB. X-ray diffraction studies that have revealed CDW ordering in related materials [28] have not observed a CDW in LiPB in spite of repeated attempts over time [29]. Low-$T$ optical spectroscopy that has observed single-particle gaps in various CDW and SDW materials [17, 30, 31] finds [17, 22] in LiPB no gap even down to 6 K and 1 meV. This finding would seem to be conclusive for the magnitude of gap ($\approx 8$ meV) expected in a mean field description of a CDW transition at $T_{min}$, as seen from the BCS curve of figure 1(b). Comparison to mean field is done here for lack of a better alternative. On the one hand, the actual transition temperature for a quasi-1D system is likely controlled by a small $t_\perp$, and so can be much less than the mean field temperature for a single chain [32]. In this case the actual gap could be much larger than that obtained in this estimate. On the other hand the single-chain gap for $T = 0$ can be much reduced from its mean field value by LL fluctuations on the chain, in which case the gap could be smaller than in this estimate. In fact figure 1(b) shows that the $T$-dependent gap that is deduced [26] by describing the upturn as an Arrhenius law is slightly smaller than 1 meV at low $T$ and so could have gone undetected in the optical work. Such a small gap was invoked as an explanation for the recent and very exciting finding [26] that a sufficiently large magnetic field applied along the 1D direction can entirely suppress the upturn. The explanation put forth is that the magnetic field quenches the CDW and hence the gap...
Figure 2. ARPES spectra of Li$_{10.9}$Mo$_6$O$_{17}$. (a) Intensity map of 300 K spectra of a $T$-dependent data-set analyzed quantitatively in [42]. D is shown as a dashed line because it is strong [7–9] only for different experimental geometries and $k$-paths. (b) Overplot of spectra in the range of box in (a). Different colors represent different $k$ values with increments 3.6 % of $\Gamma$–$Y$. (c) Spectra for $k = k_F$ normalized as described in [42] and in the text. Reproduced with permission from [42]. Copyright 2009 American Physical Society.

because Zeeman splitting of the bands disrupts FS nesting. The discussion in [26] also took up a previous proposal [24] for a purely electronic CDW, driven entirely by Coulomb interactions in a strong coupling limit. This proposal was made to overcome indications from optical spectroscopy [22] and thermal expansion [24] measurements that the lattice has no strong involvement with the upturn, which implies that a conventional CDW involving a periodic lattice distortion does not occur. While such a purely electronic CDW would be more difficult to detect in the conventional x-ray diffraction performed for LiPB to date [29], it is certainly not impossible as shown by such an observation for a Bechgaard salt [33].

The great beauty of the experimental magnetoresistance result [26] and the neatness of the explanation notwithstanding, there is no CDW theory that would explain such an unusual gap function and one must also ignore another important missing signature of a single-particle gap [10, 12, 22], that the dc magnetic susceptibility measured by three groups over time is entirely unaffected by the resistivity upturn at $T_{\text{min}}$. Reference [22] points out that because the magnitude of the measured susceptibility involves a near cancelation of several large negative and positive contributions, the measurement is particularly sensitive to a change in any one of them, e.g. that of the conduction electrons. This is perhaps the strongest single piece of all the many pieces of evidence that collectively weigh heavily against a CDW. The lack of CDW signatures in LiPB is all the more striking because the closely related compounds KMo$_6$O$_{17}$ and NaMo$_6$O$_{17}$ do exhibit CDW behavior clearly seen with x-ray diffraction [28] and optics [17]. These materials are also planar but have higher symmetry such that the planes contain three Mo–O chains oriented at 120° to one another. The CDW $Q$-vector is such as to gap out the quasi-1D bands associated with two of the three chains. Finally the possibility of an SDW in LiPB has been excluded by muon spin relaxation measurements [34]. The last section of the paper returns to the issue of the upturn and to the other parts of figure 1.

3. Signatures of Luttinger liquid behavior

Early workers did not include the ideas of the Luttinger liquid (LL) in their thinking. LL behavior in LiPB was first demonstrated by photoemission experiments. Indeed, these experiments were undertaken specifically because of the missing signatures for CDW formation. Pioneering studies [35] had observed in angle integrated photoemission of certain quasi-1D materials one of the characteristic features expected of LL behavior, that the angle integrated single-particle spectrum approaches $E_F$ as a power law. These studies suffered from some ambiguity because the materials studied all manifested low $T$ CDW formation and CDW fluctuations at higher $T$ can in principle cause pseudogap behavior [32, 36] that might mimic the expected LL behavior.
Thus LiPB was conceived [7] as a non-CDW quasi-1D paradigm for studying LL behavior. It was also envisioned that ARPES experiments could distinguish features of the single-particle spectral function that are unique to LL behavior. A series of ARPES studies dating from 1999 [7–9, 37–42] has born out this motivation. This section gives a guide to the key results of these various past ARPES studies and points out two other measurements using different techniques that have added very importantly to the case for LL behavior. Sections 4 and 5 present new photoemission data and a new phenomenological analysis of ARPES data, respectively.

ARPES measures the single-particle spectral function. For the paradigm one band Tomonaga–Luttinger (TL) model [43] the spectral function [44] shows three characteristic signatures of LL behavior. The ARPES measurements described below have demonstrated all three of these signatures. First, as already mentioned, the momentum \((k)\)-summed spectrum approaches \(E_F\) as a power law characterized by an anomalous exponent \(\alpha\). Second, the \(k\)-resolved single-particle lineshape is composed of two features, a holon peak with a leading spinon edge, the two dispersing with different velocities, \(v_c, v_s\) respectively. This lineshape signifies the absence of quasi-particles and the fractionalization of the electron into density fluctuation modes of charge (holon) and spin (spinon). Third, the TL-model is quantum critical (QC). QC systems display scale invariance, i.e. at \(T = 0\) K and for large distances and long times, the correlation functions lack a characteristic scale and take the simplest possible scale-free, functional form, a power law. Departures from \(T = 0\) K satisfy simple scaling laws [45, 46] with the only scale. For the case of a spin rotationally invariant interaction the TL-model spectral function [44] explicitly obeys the ideal scaling form \(A(k, \omega, T) = T^n A(\nu k / T, \omega / T)\) where \(A\) is a universal scaling function, \(k\) is measured from the Fermi momentum \(k_F\), \(\omega\) is measured from the Fermi energy \(E_F\), and \(\nu\) is a constant with units of velocity. Within the one band TL-model the \(T\)-scaling exponent \(\eta\) is then also determined by a scaling relation to be \(\eta = (\alpha - 1)\).

The initial ARPES studies [7] were performed at relatively high temperatures of 200–300 K to avoid any influence of the putative CDW transition at \(T_{\text{min}}\) but the TL-model theoretical lineshapes then available [47] for comparison [7, 37] to the ARPES data were for \(T = 0\) K. Over time the quality of the spectra improved and nonzero-\(T\) theory lineshapes [44] became available. The latter were particularly important because including the effect of temperature changed the parameters deduced from the \(T = 0\) K comparison such as to improve the internal consistency of the description. It was found thereby [38, 39] that the high \(T\) ARPES spectra could be generally well described by the theory lineshapes for the measurement \(T\) and for an \(\alpha\) value the same as determined directly from the angle integrated spectrum. These results can be seen in figures 7 and 8 of [38] and figures 4, 5 and 6 of [39]. Figure 7 of [38] sketches the bands and shows ARPES data for \(T = 250\) K. Figure 8 of [38] compares the data to theoretical lineshapes calculated for \(T = 250\) K, for \(\alpha = 0.9\) and for a range of values of \(\nu / v_c\). The lineshapes include the experimental broadening in \(k\) and \(\omega\). The choice of \(\nu / v_c = 2\) gives the best agreement with the data. Figure 6 of [39] shows a similar comparison of the data to theory for a range of values of \(\alpha\) and one can see that \(\alpha = 0.9\) gives the best agreement for the rate of falloff of the holon peak intensity as \(k\) approaches \(k_F\). Figure 7 of [38] is interesting for showing explicitly that \(k\)-integration of ARPES data for a Fermi liquid material yields a Fermi edge whereas one obtains a power law at \(E_F\) for LiPB.

Figure 8 of [38] (the same as figure 5 of [39]) makes an important point concerning the ARPES lineshapes. For \(\alpha > 0.5\), as is the case for LiPB, the spinon feature of the lineshape is an edge singularity rather than a peak singularity. Including also the broadening due to temperature and experimental resolutions yields a lineshape which at first glance does not have two distinct features and so could arouse skepticism as to the claim of observing spin–charge separation. Nonetheless the various panels of the figure make it clear that the two features are indeed present and are still very visible because the spinon edge disperses at a rate that differs from that of the holon peak, controlled by varying \(\nu / v_c\). Thus these lineshapes are quite unique, are characteristic of LL behavior, and are much different from the usual dispersing peaks seen in ARPES.

Two other important findings for the ARPES lineshapes of LiPB are that the same spectra are obtained [40] for samples made using each of the two main crystal growth techniques and that, within experimental resolutions, the same spectra are obtained when measured [41] at a photon energy \(h\nu = 500\) eV as when measured with the lower photon energies \(h\nu\) between 20 and 30 eV that were used for all the other ARPES summarized here. The significance of the former is that it dispels the possibility of the sample growth method being the origin of an early ARPES report [48] of Fermi liquid lineshapes for \(T\) above \(T_{\text{min}}\) and of an 80 meV gap for \(T\) below \(T_{\text{min}}\). This anomalous finding was disputed [49, 50] and since then has never been replicated. The significance of the latter is that the higher photon energy spectra are more bulk sensitive and the agreement of the spectra is consistent with the likelihood based on the crystal structure that the quasi-1D chains lie two layers below the cleavage plane and are hence well protected from surface effects, the same situation that has made ARPES relevant to the bulk properties of many superconducting cuprate materials.

Scanning tunneling spectroscopy (STS) [51] has made an important contribution to the case for LL behavior. This work is notable for its \(T\)-range from 55 K down to 5 K, only slightly above \(T_{\text{SC}}\). The experimental resolution deduced from the STS spectra is 9 meV, essentially the same as the gap value implied by a mean field transition at \(T_{\text{min}}\). STS also observes a power law density of states at \(E_F\), although the values of \(\alpha \approx 0.6\) are smaller than those measured in the high temperature ARPES. No change was detected in the \(T\)-range of the resistivity upturn, consistent with the lack of a gap found in optical spectroscopy. As with ARPES it is likely that the quasi-1D chains being probed in this surface sensitive spectroscopy are protected from surface effects by lying well below the cleavage plane.

Photoemission measurements extended to lower temperatures have continued to support LL physics, but have
also revealed interesting disagreements with TL-model predictions. Lineshapes from angle integrated ARPES data from 300 K down to 30 K are very well fitted by TL lineshapes but only if \( T \) is \( T \)-dependent [9]. The \( T \)-dependence of \( \alpha \) nicely connects the high \( T \) value of \( \approx 0.9 \) to the low \( T \)-value of \( \approx 0.6 \) found in STS. However, the fact that \( \alpha \) is \( T \)-dependent is clearly outside the one band TL-model. A microscopic theory offered in [9] is that Coulomb interactions involving the two bands of the two chains per unit cell (implying four collective modes, two holons, two spinons [3, 52]) give rise to a \( T \)-dependent renormalization of \( \alpha \).

\( T \)-dependent ARPES lineshapes measured over the \( T \) range 300–30 K showed QC scaling [42]. However the \( T \)-dependence observed was not in agreement with expectations from the TL-model. As shown in figures 1(b)–(d) of [42], although the lineshapes sharpened considerably with decreasing \( T \), the sharpening was not as great as predicted, even when taking account of the measured \( T \)-dependence of \( \alpha \). An alternative microscopic theory [53] is that the low \( T \) value can have a purely electronic origin involving the coupled chains, but that the increase from 0.6 to 0.9 arises from the influence [54] of phonons on the underlying bands.

We now summarize aspects of the \( T \)-dependent ARPES study that are important for the later discussion in section 5. Figure 2(a) shows the 300 K spectra for wavevector \( k \) varying along the \( \Gamma \)–\( Y \) direction of the Brillouin zone. Bands A and B approach \( E_F \) no closer than 0.12 eV. Bands C and D merge and disperse to cross \( E_F \) together. For the particular \( k \)-path shown, the B band is too weak to observe and so its dispersion is sketched as a dashed line based on data [7–9] from other \( k \)-paths. Figure 2(b) overplots the spectra in the range of the box of figure 2(a) to show the dispersing holon peak and spinon edge approaching \( E_F \). The need to measure for multiple \( T \) values within the lifetime of the sample necessitated noticeably poorer statistics for these spectra than for the spectra (cited above [9, 38, 39]) that were compared in detail to the TL theory lineshapes, but apart from the poorer statistics the general features of the spectra are the same.

The \( T \)-dependent spectra were tested for QC behavior as follows. If the general scaling form of the spectral function holds, then \( T^{-\delta}\alpha(k, \omega, T) \) is independent of \( T \) if \( k \) is chosen for each \( T \) so that \( k/T \) does not change, i.e. \( k = 0 \) (the \( k_F \) spectra) or \( k = c T \) where \( c \) is a constant. As described in [42] the \( k_F \) spectra were normalized to one another by matching the leading edges (0.25–0.4 eV) of the B band peak, which lies relatively far from \( E_F \) and has no apparent \( T \) dependence. Figure 2(c) shows the normalized spectra vertically offset for clarity. The spectral intensities at \( E_F \) were then matched with a multiplicative \( T \)-dependent factor whose inverse has a power law dependence from which \( \eta \) can be deduced, as shown in figure 3(a). The deviations between the data and the power law fit have no systematic pattern and are consistent with the effects of the noise in the spectra (see figure 2(b)) giving uncertainty to the normalization process leading to figure 2(c). When the \( k_F \) \( T \)-scaled spectra are plotted versus \( (E - E_F)/k_B T \) the spinon edges align. The scaling behavior of the spectra can be visualized from figure 3(b) which shows for \( k = 0 \) the result of a phenomenology that is presented in section 5 and that gives a good description of the data of [42]. We see that the spinon edges scale, but not the holon peaks. Unscaled \( k \)-dependent spectra were normalized and \( T \)-scaled by exactly the same factors as already found for the \( k_F \) spectra and also showed the same general scaling result. In addition to the lack of scaling of the holon peaks, one sees from figure 3(a) that \( \eta \) has a value that is essentially \( \alpha \) rather than the value \( (\alpha - 1) \) that would be expected from the TL-model [44]. The difference between the fitted power law 0.56 and the measured value of \( \alpha \) is not deemed significant in view of the uncertainty in the normalization process described above. On the one hand the observed scaling behavior confirms a generic property of LL physics, but on the other hand the deviations from expectations in the TL-model raise difficult questions for 1D theory [1].

This section closes by calling attention to other very important evidence of LL behavior that has recently been reported [55]. A pioneering measurement of the ratio of the thermal and electrical Hall conductivities in the \( T \)
Figure 4. (a) Angle integrated photoemission spectra of lithium purple bronze for $T = 4, 6$ and $30$ K. The spectra at $T = 4$ and $30$ K were taken with a resolution of $5$ meV with $h\nu = 8.4$ eV. For reference a gold spectrum with the same settings is also shown. The angle integrated spectra for $T = 6$ K stems from the same data-set as the Fermi surface depicted in (c). This data-set was taken with energy resolution of $12$ meV and $h\nu = 30$ eV. All the spectra are generally well fitted by the TL-model lineshape. The gold spectrum fits well with $\alpha$ essentially zero, corresponding mathematically to a Fermi edge. Panel (b) shows the normalized residuals for each fit. Panel (c) shows the Fermi surface produced by integrating ARPES spectra $\pm 6$ meV around the Fermi energy. The green box in the upper left represents the FWHM of the resolution function in $k$-space. The black color means higher intensity. One sees clearly a very straight Fermi surface represented by the two vertical black lines. The upper BZ shows the Fermi surface as calculated by Popović [6]. Compared to the theoretical FS, there is no sign of a splitting or warping of the experimental Fermi surface.

range 300–25 K observed a spectacular violation of the Wiedemann–Franz law. As explained in [55] the spin–charge separation of LL physics offers a very natural explanation for this failing because heat transfer proceeds by both spin and charge degrees of freedom whereas electric current involves charge alone. This work is notable for being clearly bulk sensitive and for demonstrating a general transport technique capable of distinguishing between Fermi liquid and LL behaviors.

4. New photoemission data just above $T_{SC}$

Up to now the only single-particle spectroscopy data for temperatures not far above $T_{SC} = 1.9$ K has come from the STS measurement cited above [51] for which the lowest $T$ was $5$ K. Our previous $T$-dependent photoemission studies, angle integrated spectra shown in figure 2 of [9], and ARPES spectra shown in figure 1 of [42], extended from $300$ K down to only $30$ K. In this section we present and discuss two new sets of photoemission data taken at temperatures comparable to those of the STS data. These new data provide a much higher resolution view of the FS than we have previously published in figure 1 of [7] and also allow a more direct comparison to the low $T$ STS results. As discussed below, the new data also augment our previous bulk-sensitive study in [41]. We note that it is very challenging to achieve very low temperatures with ARPES because of the complexities of sample holders with multiple mechanical degrees of freedom, and the difficulties of providing thermal shielding of the sample while also allowing for photon excitation and electron collection.

The new data were taken with two different ARPES setups where the lowest temperatures were $T = 4$ K and $5$ K respectively. One setup is situated in the laboratory of the Department of Material Physics, Graduate School of Engineering Science of the University of Osaka. Here we obtained angle integrated data with a photon energy of $8.4$ eV and a resolution of $5$ meV. The other setup is the MERLIN Beamline 4.0.3 at the Advanced Light Source (ALS) synchrotron. Here we measured with $h\nu = 30$ eV and a resolution of $12$ meV. For both experiments, single-crystal Li$_{0.9}$Mo$_6$O$_{17}$ samples were grown using the temperature gradient flux method [10]. The instrumentation in the Osaka laboratory consists of an MB Scientific MBS $T = 1$ microwave excited rare gas lamp monochromatized by use of ionic crystals (CaF$_2$ for Kr and sapphire for Xe) [56]. Here Xe was used for the lamp which gives a photon energy of $h\nu = 8.4$ eV. Other instrumentation consists a low temperature closed-cycle He cryostat manipulator and a Scienta SES2002 electron energy analyzer driven by an MBS A-1 power supply. The beamline at the ALS utilizes a low temperature six-axis sample manipulator cooled with an open-cycle He
flow cryostat and a Scienta R8000 kinetic energy analyzer. The MERLIN beamline has an elliptically polarized undulator which allows one to choose arbitrary polarizations of the incident light. Due to the longer elastic escape depth of lower kinetic energy electrons, the bulk sensitivity of data from the Osaka measurements with \( hν = 8.4 \text{ eV} \) is expected to be higher than that for the synchrotron measurements with \( hν = 30 \text{ eV} \) Thus this study complements our previous bulk-sensitivity study [41] but the resolution is much higher than could be achieved using the very high photon energy of [41]. The drawback of that enhanced bulk sensitivity is a lower cross-section of the photoemission process, which requires longer acquisition times compared to the synchrotron in order to achieve sufficiently good statistics. In the present case this difficulty was exacerbated by the desire for high resolution and the small size of the sample relative to the photon spot size of the laboratory setup, so that the acquisition time difference was almost ten times longer than that of the synchrotron. On the other hand, the temperature for the laboratory measurement is essentially the same as the lowest \( T \) of the STS measurement while its resolution is even higher and would be capable of detecting the \( \approx 8 \text{ meV} \) mean field gap of a transition at \( T_{\text{min}} \).

We first present the new ARPES results obtained at the ALS. The sample was orientated such that the 1D chains (\( b \) direction) were along the angular axis of the detector, which is along the \( \Gamma–Y \) direction in figure 4(c) and was cleaved in situ on the cold cryostat in a vacuum better than \( 8 \times 10^{-11} \text{ Torr} \). The photon polarization was chosen with the electric field vector perpendicular to the \( b \) direction. The temperature was maintained at \( T = 6 \text{ K} \). A 3D spectroscopic data-set \( I(k_{f}, k_{\perp}, E) \) was obtained by rotating the sample around the polar direction. The FS map shown in figure 4(c) was produced by integrating the photoemission signal \( \pm 6 \text{ meV} \) around \( E_{F} \), essentially the width of the energy resolution. The \( k \)-resolution is \( 0.01 \text{ Å}^{-1} \) along the analyzer slit (\( \Gamma–Y \)) and \( 0.03 \text{ Å}^{-1} \) perpendicular to the slit (along \( \Gamma–X \)), much better than for our previously published FS map [7]. Black means high intensity in the map and we see two vertical black lines which represent the FS. In contrast to the theoretical calculations by Popović et al [6], we see at this resolution no splitting of the FS. Also it appears that the FS is essentially straight at this resolution. A Lorentzian peak-fit estimates the value of \( k_{\perp} \) along the FS to vary by not more than \( 0.01 \text{ Å}^{-1} \), which is less than the \( k \)-space resolution and less than the variation of the LDA calculation.

In order to compare with the STS data and for comparison with the Osaka data presented below, we angle integrated the spectrum is indicated by the solid line through the data points. This fit was made using the \( k \)-integrated spectral weight of a spin-rotational invariant TL-model [44] with \( \nu = 2 \), consistent with previous results described above. The theoretical spectrum was broadened by the experimental energy resolution of \( 12 \text{ meV} \). In panel (b) of figure 4 one sees the normalized residual of the fit, i.e. the difference between the fit and the actual spectrum divided by the actual spectrum. The residual is essentially smooth. The \( \alpha \) value of \( \approx 0.65 \) is comparable to that obtained at 30 K in our previous photoemission experiments [9, 42]. As found with STS [51], there is no significant difference for \( T \) above and below \( T_{\text{min}} \) at the resolution of this measurement.

We discuss the Osaka results next. The samples were oriented by x-ray diffraction with the chains (\( b \) direction) along the angular axis of the analyzer. They were cleaved in situ on the cold cryostat manipulator in a vacuum better than \( 2 \times 10^{-10} \text{ Torr} \). As the acceptance angle of the spectrometer is \( \pm 7° \) corresponding to \( \pm 0.11 \text{ Å}^{-1} \) at \( hν = 10 \text{ eV} \), and the Brillouin zone extends along \( \Gamma–Y \) to \( 0.57 \text{ Å}^{-1} \) in the chain direction, corresponding to \( \approx 37° \), we measured at angles of \( 0° \), \( 14° \), and \( 28° \) with the same conditions of time, photon flux, and temperature, and afterwards added the spectra to obtain angle integration. After finishing the measurements, the Fermi energy was determined by evaporating gold onto the sample. The spectra of LiPB for \( T = 4 \text{ and } 30 \text{ K} \) are shown in figure 4(a) together with the corresponding gold spectrum at \( T = 4 \text{ K} \). One sees the considerably poorer statistics of the LiPB Osaka data relative to that of the ALS data or the gold spectrum. Nonetheless the sharp contrast between the reduced weight near \( E_{F} \) in the LiPB spectra and the Fermi edge of the gold spectrum can clearly be seen. Solid lines through the data indicate line-fits of the spectra using the same TL-model as for the ALS \( k \)-integrated spectrum. The theoretical spectrum was broadened by the experimental energy resolution of \( 5 \text{ meV} \) which was determined from the gold spectrum. The \( \alpha \) value of the fit is written beside each spectrum together with the standard deviation resulting from the \( \chi^{2} \)-fit. For the gold spectrum, we expected a Fermi edge and therefore \( \alpha = 0 \). This small deviation from zero and also the peak above \( E_{F} \) in the residual shown in figure 4(b) might be seen as indicating some systematic error.

In panel (b) of figure 4 one sees the normalized residuals of the fits. The numbers next to the curves indicate how much the curves were shifted in the vertical direction for improved clarity of presentation. Comparison of the residuals again shows clearly the poorer statistics of the Osaka LiPB data. Both the \( T = 4 \text{ and } 30 \text{ K} \) spectra show a deviation from the TL lineshape at around \( 5 \text{ meV} \) binding energy. This deviation is larger for \( T = 4 \text{ K} \) than for \( T = 30 \text{ K} \). Since the resolution of \( 5 \text{ meV} \) is better here than in the STS experiment, it is logically possible that such a feature could have escaped detection in the STS spectrum. But a skeptic might point out that because it appears also in the 30 K spectrum it does not correlate directly with the \( T \)-dependent transport properties. A skeptic might also claim a hint of this feature even in the gold spectrum residual. Thus the question of whether this deviation is real or is due to a small systematic error, i.e. not perfectly linear behavior of the detector or a contribution from the sample holder because the photon spot is large, or is simply the result of insufficient acquisition time to obtain better statistics, cannot be absolutely answered. We present the spectra as a further confirmation of the overall large energy scale power law behavior and regard them as ambiguous on

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8 The kinetic energy dependence of the elastic escape depth has a so-called U-shape, being large for very low and very high energies.
the possibility of structure at the 5 meV energy scale of the resolution. We take it as significant that all the new LiPB data-sets presented here give essentially the same $\alpha$ values $\approx 0.65–0.7$. This is very much consistent with the range of values found in STS and perhaps consistent with a slight low $T$ upturn of $\alpha$ found previously, as shown in figure 3(a) of [9].

### 5. Phenomenological description of $T$-dependent ARPES data

The findings of [42] show that Li$_{10.9}$Mo$_6$O$_{17}$ is a QC system, but with important differences relative to expectations from the one band TL-model, specifically in the exponent of the temperature prefactor and in the lack of the full sharpening predicted for decreasing $T$, which is part of the non-scaling of the holon peaks.

These differences can be partially described by a phenomenological momentum broadening of the TL spectral function, but the required broadening greatly exceeds the experimental momentum resolution, as discussed below. The fact that the experimental scaling prefactor is $T^\alpha$ rather than $T^{(\alpha−1)}$ is a basic motivation to try an integration of the theoretical TL spectral function because that will draw a factor of $1/T$ outside the integral. Convolving the theoretical spectral function (for which $\eta = (\alpha − 1)$) with a momentum window function $R(p/p_0)$, where $p_0$ is a $T$-independent width, and using the change of variables $\tilde{p} = \frac{(\eta p)}{T}$ (recall that $v$ is a constant with units of velocity), gives

$$A_{\text{test}}(k, \omega, T) = T^\alpha \int_{-\infty}^{\infty} \tilde{R} \left( \frac{\tilde{p}}{p_0} \right) \tilde{A} \left( \tilde{k} - \tilde{p}, \frac{\omega}{T} \right) d\tilde{p}. \tag{1}$$

A constant factor of $1/v$ has been drawn into a redefined $\tilde{A}$. $A_{\text{test}}(k, \omega, T)$ has the $T^\alpha$ prefactor observed experimentally but the rest of the expression is no longer a universal function of $\omega$ because $T$ also enters in the form $p_0 = \frac{v}{p_0 T}$ and $A_{\text{test}} = T^\delta f \left( \frac{v p_0}{T}, \frac{\omega}{T} \right)$. It can be noted in passing that this exercise also shows the way in which a fixed experimental resolution has an increasingly large effect on a QC spectrum as temperature decreases.

Choosing $R$ to be a simple normalized Gaussian with $p_0 = 0.065$ Å$^{-1}$ and choosing the low temperature value of $\alpha = 0.6$, consistent with the experimental prefactor, yields a reasonable description of the experimental data below the binding energy for which the intensity from band B becomes important. Figure 3(b) shows the general scaling behavior of $A_{\text{test}}$, including the experimental energy resolution. After the intensity is multiplied by $1/T^\alpha$ and the energy axis is scaled by $k_B T$, the theoretical curves at $k_F$ have their edges falling on each other while the peak parts do not, exactly what is observed in experiment. The slight deviation from the scaled edge at low temperature above $E_F$ is seen in the data of [42], and as shown there, is due to the experimental resolutions.

Figure 5(a) compares the $k_F$ data of figure 2(c) for ten values of $T$ with the phenomenological $A_{\text{test}}$ and figures 5(b) and 5(c) for $T = 300$ K and 100 K respectively. Parameters are: $v_c/v_s = 2$ and $\alpha = 0.6$, $v_s = 1.9$ eV Å and $p_0 = 0.065$ Å$^{-1}$, the same for all panels. The unit of $k$ is Å$^{-1}$.

![Figure 5](image-url)

**Figure 5.** Fitting the ARPES data of [42] to $A_{\text{test}}$. (a) Fitting for $k = k_F$ at various temperatures. (b), (c) Fitting to both $k = k_F$ and $k$ away from $k_F$ at 300 K and 100 K respectively. Parameters are: $v_c/v_s = 2$ and $\alpha = 0.6$, $v_s = 1.9$ eV Å and $p_0 = 0.065$ Å$^{-1}$, the same for all panels.
and (c) make the comparison for four values of \( k \), \( T \)-scaled away from \( k_F \) at two temperatures. In these figures, the amplitudes of the experimental curves have been adjusted to remove the slight (maximum 15%) prefactor differences from \( T^\alpha \) seen in figure 3(a) and ascribed in the section 3 discussion of that figure to uncertainty in the normalization process. By eye there is generally good agreement in these comparisons until reaching the binding energy at which the higher lying bands begin to contribute intensity in the experimental spectra. From the spectra of figure 2(a) and taking account of the width of band B, this energy is 0.1–0.12 eV. For \( k \)-values away from \( k_F \) the binding energy of band B increases and so the deviation also moves to higher binding energies, nearly 0.2 eV. In spite of this agreement, it is important to point out that this phenomenology does not produce the \( T \)-dependence of \( \alpha \) that is measured from \( T \)-dependent \( k \)-integrated spectra. In fact it works only for \( k \)-values rather near to \( k_F \) whereas a much larger range of \( k \) contributes to the \( k \)-integrated spectrum that determines \( \alpha \). Thus the phenomenology is not a substitute for a satisfactory microscopic model.

What is the meaning of this phenomenology? The simplest possibility is of course that the momentum is indeed broadened. One could argue that the QC scaling predicted in the TL-model is being observed except for being modified by the experimental resolution. However, the value of \( p_0 \) is five times larger than the instrumental \( k \) resolution, computed from the instrumental angle resolution as 0.013 Å\(^{-1}\). Such a large extra broadening would have to be ascribed to sample surface quality. For example, our relatively large spot size might be extra broadening would have to be ascribed to sample surface quality. For example, our relatively large spot size might be illuminating regions with slightly different orientations on the sample surface. The size of the discrepancy makes this interpretation more difficult to defend, but nonetheless the phenomenology shows that the issue of finite momentum resolution is very important. ARPES scaling data might evolve with progress in sample growth and surface preparation techniques.

Another possibility is that the phenomenology tries to catch some of the intrinsic physics. For example it is clear from the \( T \)-dependence of \( \alpha \) that the one band TL-model is missing interactions of importance and we know from band theory and the ARPES data that there are two bands crossing \( E_F \), associated with the two chains per unit cell. Perhaps the phenomenology can be a guide for future theory. It would be very desirable to test the data against the spectral function for a two-band TL-model if it were available.

6. Robust 1D physics

The combination of data from ARPES and STS along with the Wiedemann–Franz law violation cited above firmly establish generic signatures of LL physics in LiPB for \( T \) at least down to \( T_{\text{min}} \). How does the system achieve the 3D character that must obtain in the SC state? At this point it is not possible to make a unified interpretation of all the relevant data. One can only speculate in a general way, guided by current 1D theory. Recent discussions [24–27, 57] focus on the idea that the crossover to 3D occurs before \( T_{\text{SC}} \) is reached; either the resistivity rise is due to CDW formation [24, 26], which is certainly of a 3D nature, or a CDW does not occur [25, 27, 57] and that crossover is more subtle in the resistivity data [27] and in the mechanism whereby it occurs [25, 57]. Here we speculate that 1D physics may be more robust, with crossover to 3D occurring only with the transition to SC, i.e. that above \( T_{\text{SC}} \) there is non-Fermi liquid behavior rooted somehow in 1D with no single-particle gap or density wave. In this connection one can take note of recent arguments [4] for unconventional SC, possibly spin-triplet, based on the finding that the critical field is much larger than the Pauli limit.

Because the 1D physics of LiPB appears to lie outside standard 1D theory [1] we can only try to identify generic possibilities that may be robust for this line of thinking. Looking first to single-particle spectroscopy, the power laws observed for temperatures only slightly above \( T_{\text{SC}} \) in the STS and ARPES data with 12 meV resolution are an initial temptation to consider this hypothesis. As discussed in section 4, the 4 K angle integrated spectra are ambiguous on the possibility of spectral structure on a lower energy scale of the 5 meV resolution, but we note that even this energy scale corresponds to a temperature of 58 K, much higher than the \( T \) of the measurement or the \( T_{\text{min}} \) of the resistivity upturn. The energy scale can perhaps be pushed lower by the 6 K optical spectroscopy [17] performed down to 1 meV, corresponding to 11.6 K. This work found no single-particle gap and no change associated with the resistivity upturn but did employ a Drude description of the infrared data. A more recent study [22], albeit a measurement down to \( T = 10 \) K and with a lower energy limit of 6 meV, also found no gap but in addition showed that the infrared behavior is actually non-Drude. Taken together, these two studies suggest no single-particle gap and non-Fermi liquid behavior down to an energy scale of 1 meV, but definitive low temperature spectroscopic measurements probing to energy scales well below 1 meV are clearly needed.

The resistivity upturn, with the possibility of a gap less than 1 meV, as obtained in [26] and shown in figure 1(a), is the primary transport evidence suggesting crossover due to a small single-particle gap. However, the highly novel gap function and the \( T \)-independent constant dc magnetic susceptibility measured down to 2 K give us cause to consider an alternative description of the upturn.

Power laws are characteristic features of the QC nature of LL physics. Indeed a power law is expected for the resistivity of an LL system. Figure 1(a) shows that the data of [26] can be described by two power laws, one with a positive exponent for the metallic part\(^9\) and one with a negative exponent for the upturn. Figure 1(b) shows the \( T \)-dependent gap that results from recasting this fit as an Arrhenius law. The gap so obtained is very similar to that obtained in [26] and the difference is very likely due to the difficulty of accurately digitizing the upturn in the resistivity data. In this view the unusual gap function is the result of forcing an Arrhenius description onto a power law. In particular one notes that at

\(^9\) In contrast to the data of figure 1(a), there are instances [16, 55] where the resistivity above \( T_{\text{min}} \) was found to be strictly linear with \( T \).
low $T$ the gap must roll over and then actually decrease in order to map the very fast rising exponential onto the slower rising power law.

Such a two-power-law description of LiPB resistivity was first put forth in Ref. [25], where the two power laws were ascribed to two independent LLs, one for each of the two bands crossing $E_F$ with each having different values of $\alpha$ and quite different roles in the crossover physics. This interpretation neglects the Coulomb interaction that couples the chains to give symmetric and antisymmetric holon and spinon modes, as discussed in Refs. [3, 52] and applied to LiPB in Ref. [9, 53]. An alternate interpretation for the power law upturn is the proposal of Anderson localization due to disorder. In the context of LL physics, the crossover is then a change of the sign of the resistivity power law exponent at low $T$ due to disorder [58]. A great concern of early discussions [12–15] was that disorder and localization could be incompatible with SC. But, as suggested in Ref. [15], the increased sensitivity to disorder in 1D could make it possible for the energy scale of the disorder to be much less than $T_{SC}$.

As seems true of all hypotheses concerning the resistivity upturn, this one also leaves some aspects of the data unexplained, why the same $T_{\text{min}}$ would be found for the resistivity along all three axes, why hydrostatic pressure suppresses the upturn (and also enhances the SC) [18], and why the upturn can be suppressed [26] by a magnetic field. For these aspects of the data one sees the attraction of postulating a gap that can be affected by pressure or a magnetic field. Whatever is the proper understanding of the upturn, a case can be made that by virtue of a power law dependence on $T$ it signifies a continuation of 1D behavior rather than a loss. The highly directional nature of the magnetic field suppression of the upturn [26] emphasizes the importance of the 1D character in this temperature range and in this connection one notes also the remarkable recent observation [26] that for a non-superconducting crystal a sufficiently large magnetic field applied specifically along the 1D axis appeared to restore the SC with a transition temperature considerably higher than 1.9 K.

To return to the theme of the opening paragraph of the paper, the greatest issue to be confronted in any line of thinking that emphasizes 1D physics down to $T_{SC}$ for LiPB is the role and the energy scale of $t_{\perp}$. The data of figure 4(c) show that the FS is unsplit and straight to a greater extent than predicted in current LDA calculations. If one accepts the current lack of evidence for a CDW, and that is the stance of this paper, then there is a need to understand why the good nesting of the FS does not result in a CDW having all the standard properties that are readily detected for CDWs in related materials. Whatever the mechanism is it would apply down to $T_{SC}$. Within present theory the situation is very difficult but perhaps not quite impossible. We now discuss three possibilities.

In a one band LL model it is well known that the easy route to avoiding crossover due to $t_{\perp}$ is a gap in the spinon mode. The gap energy must be overcome for single-particle hopping between the chains, which renders $t_{\perp}$ to be irrelevant in the renormalization group sense [3]. But pair hopping that is second order in $t_{\perp}$ remains relevant and can lead to SC [3]. If the energy scale of the spin gap is greater than $T_{SC}$ there can be a transition directly from the spin-gapped LL (a Luther–Emery liquid) to a SC [59]. Such a gap in LiPB would seem to be precluded by the $T$-independent dc magnetic susceptibility which has been measured down to 2 K [22]. However, the two-chain model with four modes, as mentioned in the preceding paragraph, has both a symmetric and an antisymmetric spin mode. The symmetric mode that is probed in the dc susceptibility can be ungapped for some parameter ranges of this model [52]. One could think of the possibility of a gap in the antisymmetric spin mode that has not been probed in any transport measurement to date and is small enough to have escaped spectroscopic detection but is nonetheless larger than $T_{SC}$. Such a gap would also be a barrier against single-particle hopping between the chains and so might render $t_{\perp}$ to be irrelevant.

A second possibility is the so-called sliding Luttinger liquid (SLL) [60–62]. This is a lattice model for an array of coupled chains and specifically includes a mechanism for suppressing a purely electronic CDW along the chains. The physical idea [61, 62] is that Coulomb interactions between the chains lead to transverse incommensurate CDW fluctuations that modulate the charge density on the chains, and therefore the Fermi wavevector and therefore the longitudinal CDW $q$-vector, which frustrates locking of the longitudinal CDW. In this special regime the $T = 0$ ground state retains the LL properties of a single chain and is stable against $t_{\perp}$, against CDW fluctuations and SC fluctuations. A spin gap is not required but of course aids the SLL stability. But in any case the parameters of the model must be tuned to bring the system to be near such a transverse CDW instability, raising doubt as to whether the model could apply to any real material. For other parameter values the stable ground state is a CDW, a novel SC, or a Fermi liquid. In addition there are difficulties specific to its application to LiPB. The model does not address the possibility of a thermal phase transition from the SLL state to a SC. Also, as $T$ goes to zero for the SLL state, the resistivity is predicted to become infinite transverse to the chains and to become zero parallel to the chains. For the transverse direction the resistivity upturn is then consistent with this prediction but for the parallel direction one must invoke some other mechanism such as disorder. But one must then explain why $T_{\text{min}}$ is the same for all directions and one must take account of how the transverse interactions may change the effect of disorder [62]. Nonetheless the concept that the chain CDW can be suppressed by frustration arising from competition between parallel and transverse CDW fluctuations may be more robust than the details of the model calculations that have been done and so the SLL remains interesting for LiPB.

In the absence of a spin gap or a scenario like the SLL, the only possibility for the line of thinking in this section would seem to be that the energy scale of $t_{\perp}$ is less than $T_{SC}$. In that case there is presently no conclusive theory for or against direct crossover of LL behavior to SC. The question entails the interplay of the pair tunneling terms needed for SC and the extent to which Coulomb interactions might suppress
or enhance the pair tunneling. Could the energy scale of $t_\perp$ be this small in LiPB? Current LDA calculations [6, 63] suggest values of $t_\perp \approx 30$ meV that would preclude this possibility. However, 1D fluctuations on the chains are known to produce a suppression of $t_\perp$ [2, 3, 64] to an effective value $t_\perp(t_\perp/t_\parallel)^{\alpha/(1-\alpha)}$, where $t$ is the intra-chain hopping. For $t \approx 800$ meV suggested by LDA hands for LiPB and the measured low $T$ value of $\alpha \approx 0.6$, one obtains the very small effective hopping value of 0.22 meV, essentially the same as $T_{SC}$. Although this estimate is extremely sensitive to the value of $\alpha$, which is not known with such great precision, nonetheless it shows that the effective hopping could be very greatly suppressed, consistent with the FS data in figure 4(c).

To pursue this idea, our current ARPES research aims at quantifying the single-particle interchain hopping and setting a more precise bound on the extent of warping and splitting of the FS through yet higher resolution measurements just above $T_{SC}$. New LDA calculations using the Nth-order muffin-tin orbitals (NMTO) method [65] and downfolding can be used to characterize the $t_\perp$ hopping in greater detail. It is anticipated that the theoretical $t_\perp$S can then be adjusted in a realistic way to describe the ARPES data and that the adjustments required can be compared to expectations from 1D theory. Thereby it is hoped to obtain a more precise picture of the electronic structure that is giving rise to the remarkably robust 1D properties and perhaps unconventional SC of LiPB.

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References

[1] Giamarchi T 2009 Physics 2 78
[2] Boies D, Bourbonnais C and Tremblay A-M S 1995 Phys. Rev. Lett. 74 968
[3] Giamarchi T 2004 Quantum Physics in One Dimension (Oxford: Oxford University Press) chapter 8
[4] Mercure J F et al 2012 arXiv:1203.6672
[5] Whangbo M-H and Canadell E 1988 J. Am. Chem. Soc. 110 358
[6] Popović Z S and Satpathy S 2006 Phys. Rev. B 74 045117
[7] Denlinger J D et al 1999 Phys. Rev. Lett. 82 2540

[8] Gweon G-H et al 2001 J. Electron Spectrosc. Relat. Phenom. 117/118 481
[9] Wang F et al 2006 Phys. Rev. Lett. 96 196403
[10] Greenblatt M et al 1984 Solid State Commun. 51 671
[11] Schlenker C et al 1989 Phys. Rev. B 39 105120
[12] Matsuda Y et al 1986 J. Phys. C: Solid State Phys. 19 6039
[13] Matsuda Y et al 1986 Physica B+C 143 243
[14] Sato M, Matsuda Y and Fukuyama H 1987 J. Phys. C: Solid State Phys. 20 L137
[15] Ekino T et al 1987 Solid State Commun. 63 41
[16] Boujida M et al 1988 Physica C 153–155 465
[17] Degiorgi L et al 1988 Phys. Rev. B 38 5821
[18] Filippini C E et al 1989 Physica C 162–164 427
[19] Schlenker C, Dumas J, Eschrig-Filippini C and Guyot H 1989 Low Dimensional Electronic Properties of Molybdenum Bronzes and Oxides (Physics and Chemistry of Materials with Low-Dimensional Structures vol 11)
[20] Duns J and Schlenker C 1993 Int. J. Mod. Phys. B 7 4045
[21] He J, unpublished
[22] Choi J et al 2004 Phys. Rev. B 69 085120
[23] da Luz M S et al 2007 Phys. Rev. B 76 233105
[24] dos Santos C A M et al 2007 Phys. Rev. Lett. 98 266405
[25] dos Santos C A M et al 2008 Phys. Rev. B 77 193106
[26] Xu X et al 2009 Phys. Rev. Lett. 102 206602
[27] Chen H et al 2010 Europhys. Lett. 90 67010
[28] Fourny P and Pouget J P 1993 Int. J. Mod. Phys. B 7 3973
[29] Pouget J P, private communication
[30] Travaglini G and Wachter P 1984 Phys. Rev. B 30 1971
[31] Degiorgi L et al 1996 Phys. Rev. B 54 3838
[32] Lee P A et al 1973 Phys. Rev. Lett. 31 462
[33] Pouget J P and Ravy S 1997 Synth. Met. B 85 1523
[34] Chakravarty J et al 2005 Physica B 359–361 1333
[35] Darbell B et al 1991 Phys. Rev. Lett. 67 3144
[36] McKenzie R H and Scarratt D 1996 Phys. Rev. B 54 R12709
[37] Gweon G-H et al 2002 Physica B 312/313 584
[38] Allen J W 2002 Solid State Commun. 123 469
[39] Gweon G-H, Allen J W and Denlinger J D 2003 Phys. Rev. B 68 195117
[40] Gweon G-H et al 2004 Phys. Rev. B 70 153103
[41] Wang F et al 2006 Phys. Rev. B 74 113107
[42] Wang F et al 2009 Phys. Rev. Lett. 103 136401
[43] Tomonaga S 1950 Prog. Theor. Phys. 5 544
[44] Luttinger J M 1963 J. Math. Phys. 4 1154
[45] Orgad D 2001 Phil. Mag. B 81 377
[46] Hertz J A 1976 Phys. Rev. B 14 1165
[47] Sachdev S 1999 Quantum Phase Transitions (Cambridge: Cambridge University Press)
[48] Meden V and Schönhammer K 1992 Phys. Rev. B 46 15753
[49] Xue J et al 1999 Phys. Rev. Lett. 83 1325
[50] Gweon G-H et al 2000 Phys. Rev. Lett. 85 3985
[51] Smith K E et al 2000 Phys. Rev. Lett. 85 3986
[52] Hager J et al 2005 Phys. Rev. Lett. 83 186402
[53] Wu C et al 2003 Phys. Rev. B 68 115104
[54] Chudzinski P, Jarlborg T and Giamarchi T 2012 Phys. Rev. B 86 075147
[55] Jarlborg T, Chudzinski P and Giamarchi T 2012 Phys. Rev. B 85 235108
[56] Wakeham N et al 2011 Nature Commun. 2 396
[57] Suga S et al 2010 Rev. Sci. Instrum. 81 105111
[58] da Luz M S et al 2011 Phys. Rev. B 84 041407
[59] Giamarchi T and Schulz H J 1988 Phys. Rev. B 37 325
[60] An example of such a transition is given by, Carlson E W et al 2000 Phys. Rev. B 62 3422
[61] Emery V J et al 2000 Phys. Rev. Lett. 85 2160
[62] Vishwanath A and Carpenter D 2001 Phys. Rev. Lett. 86 676
[63] Mukhopadhyay R et al 2001 Phys. Rev. B 64 045120
[64] Satpathy S, private communication
[65] Bourbonnais C et al 1984 J. Physique Lett. 45 L755
[66] Anderson O K and Saha-Dasgupta T 2000 Phys. Rev. B 62 R16219

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