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LETTER TO THE EDITOR

Anomalous behaviour of the in-plane electrical conductivity of the layered superconductor κ-(BEDT-TTF)$_2$Cu(NCS)$_2$

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Abstract
The apparent quasiparticle scattering rates in high-quality crystals of the quasi-two-dimensional superconductor κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ are studied using the Shubnikov–de Haas effect and megahertz penetration-depth experiments. The width of the superconducting transition observed in the megahertz experiments, taken in conjunction with the field dependence of the Shubnikov–de Haas oscillations, gives evidence that the broadening of the Landau levels is primarily caused by spatial inhomogeneities. This indicates a quasiparticle lifetime for the Landau states $\gg$3 ps. The megahertz data can also be used to derive an apparent scattering time (0.14–0.56 ps) from the skin depth. This is much *shorter* than the Landau-state lifetime, in strong contrast to the expectations of Landau Fermi-liquid theory. The simplest explanation for the data is that only a fraction of the crystal contributes to the metallic conductivity, an observation which may be related to the recently observed 'glassy' transition in κ-(BEDT-TTF)$_2$Cu(NCS)$_2$.

(Some figures in this article are in colour only in the electronic version)

In this letter, we study the apparent quasiparticle scattering rates in the high-quality single crystals of the very extensively studied organic metal κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ [1]. It is usually assumed that the low-temperature properties of this quasi-two-dimensional correlated-electron system conform to Landau Fermi-liquid theory (LFLT) [2, 3]; its magnetic quantum oscillations and magnetoresistance are apparently describable in terms of a Fermi surface with

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a simple and well-defined topology [4–6]; moreover, it is often regarded as a *model* layered metal [4] and superconductor (see [7, 8] and references therein). However, our data show that the *apparent* scattering rates differ greatly from the expectations of simple LFLT. The simplest explanation for these data is that only a small fraction of the crystal is metallic at low temperatures, suggesting that the assumption that κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ is a relatively homogeneous metal may be misplaced.

Our experiments measure the Shubnikov–de Haas (SdH) oscillations, the skin depth at megahertz frequencies and the width of the superconducting transition. They were carried out on single crystals of κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ ($\sim$0.5 $\times$ 0.25 $\times$ 0.1 mm^3; mosaic spread $\lesssim$0.1°), produced using electrocrystallization [9]. The resistance of the samples was measured using standard four-wire ac techniques (frequency 15–180 Hz, current 1–20 μA) [4]. Contacts were applied to the upper and lower faces of the crystals so that the current was directed and the voltage measured in the interlayer direction; this gives a resistance R_{zz} proportional to the interlayer resistivity, ρ_{zz} [10].

Penetration and skin depths were inferred by placing a sample in a small coil which is the inductive element of a resonant tank circuit [11]. The exclusion of flux from the sample, and hence the coil, decreases the inductance of the circuit; consequently the resonant frequency, $f = 1/(2\pi\sqrt{LC})$, increases. Through simple geometrical relations [11] it can be shown that

$$\Delta\lambda = -\frac{R^2}{r_{\mathrm{s}}}\frac{\Delta f}{f_0} \tag{1}$$

where $\Delta\lambda$ is the change in penetration/skin depth causing a change in frequency Δf, R is the effective radius of the coil, r_{s} is the effective sample radius and f_0 is the unperturbed resonant frequency. The samples were aligned with their highly conducting planes perpendicular to the axis of the coil to $\pm 2°$. Analysis of this geometry following [12] shows that the response is dominated by *in-plane* screening currents, even if the samples are slightly misaligned; this is a consequence of the large resistivity anisotropy in κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ [1, 10, 13]. It is for this reason that we have chosen this technique to measure the skin depth and hence the in-plane conductivity (cf [10]).

Results from all samples yielded similar behaviour of the apparent scattering rates; for ease of comparison, we discuss data from a single sample ($r_{\mathrm{s}} \approx 0.13$ mm).

Figure 1 shows R_{zz} ($\propto \rho_{zz}$) for a κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ sample; note that SdH oscillations due to the α pocket of the Fermi surface (see inset) are observed in R_{zz} at fields above the superconducting transition; the observation of oscillations at fields as low as 6 T (figure 1(b)) indicates the high quality of the samples. Scattering rates in metallic systems are often derived from the rate at which such quantum oscillations grow in amplitude with increasing field [14]. This is parametrized by the *Dingle temperature* T_{Ds}, which is traditionally related to the scattering rate τ_{D}^{-1} by [14]

$$T_{\mathrm{Ds}} = \frac{\hbar}{2\pi k_{\mathrm{B}}\tau_{\mathrm{D}}}. \tag{2}$$

Here, $\hbar\tau_{\mathrm{D}}^{-1}$ is usually taken to represent the energy width of the Landau levels associated with their finite lifetime due to scattering.

However, other processes can contribute towards the energy width of the Landau levels [13, 15, 16]. For example, screening is less effective in systems with low densities of quasiparticles (such as organic metals), compared with that in elemental metals [1]; hence variations in the potential experienced by the quasiparticles can lead to a spatial variation of the Landau-level energies. Even in the (hypothetical) complete absence of scattering [15], this spatial variation would give the Landau level a finite energy width and therefore lead to a

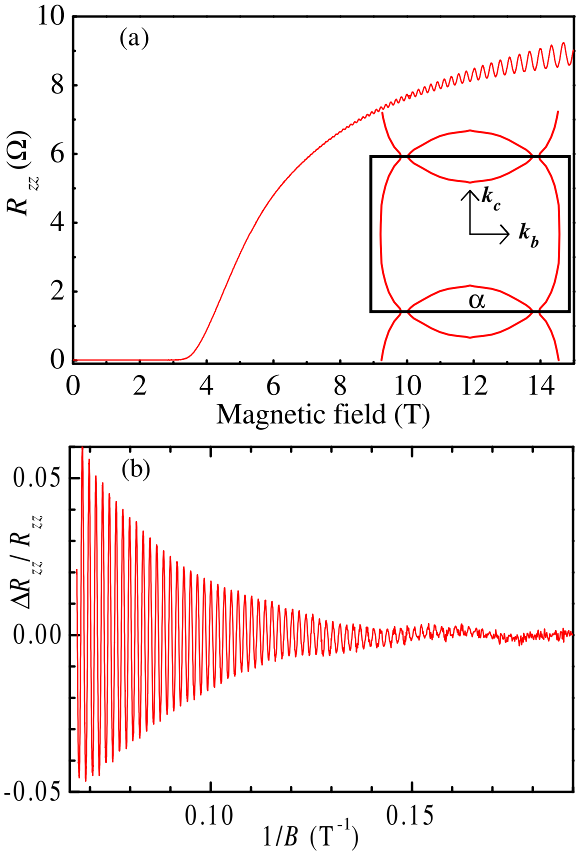

Figure 1. Resistance R_{zz} ($\propto \rho_{zz}$) of a κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ crystal with magnetic field perpendicular to the quasi-two-dimensional planes. (a) Raw data ($T = 0.8$ K); the superconducting to normal transition is visible, as are SdH oscillations due to the closed α pocket of the Fermi surface (see inset). (b) Normalized oscillatory component of magnetoresistance $\Delta R_{zz}/R_{zz}$ obtained by dividing R_{zz} by a function fitted to the non-oscillatory background and subtracting 1. Inset; Brillouin zone and Fermi surface cross-section showing the α pocket and the pair of Q1D sheets (based on [4]).

contribution T_{Di} to to the apparent Dingle temperature [16],

$$T_{\mathrm{Di}} = \frac{\bar{x}[1-\bar{x}]F'(\bar{x})^2 a}{\pi k_{\mathrm{B}} m^*}\sqrt{\frac{\hbar e^3}{2F}}. \tag{3}$$

Here F is the magnetic quantum oscillation frequency, and $F' = \mathrm{d}F/\mathrm{d}x$; x represents the (local) fractional variation of the quasiparticle density due to the potential fluctuations and $\bar{x}$ is its mean [16]. Hence, the Dingle temperature measured in experiments (T_{Dm}) will be a combination of T_{Ds} and T_{Di}; in the spirit of Matthiesson's rule [3], we propose[5] [15] $T_{\mathrm{Dm}} \approx T_{\mathrm{Di}} + T_{\mathrm{Ds}}$.

At the moderate fields employed in figure 1, the oscillatory magnetoresistance is much less than the non-oscillatory component (see figure 1(b)), and magnetic breakdown is a minor

[5] In some quasi-two-dimensional systems, it appears that the Landau-level broadening due to spatial inhomogeneities is decoupled from that due to scattering events; i.e. the inhomogeneities do not appear to induce a measurable increase in the scattering rate (see [15]).

consideration [5]. Hence, the Lifshitz–Kosevich formula may be used to extract the effective mass m_α^* and T_{Dm} for the α pocket, (the latter is denoted $T_{\mathrm{Dm}\alpha}$) from the temperature (T) and field dependences of the oscillation amplitude [14]. This yields $m_\alpha^* = 3.5 \pm 0.1 m_e$, in agreement with published values [1, 4], and $T_{\mathrm{Dm}\alpha} \approx 0.42$ K; the latter is equivalent to $T_{\mathrm{Dm}\beta} \approx 0.76$ K for the β breakdown orbit [5]. If this Landau-level broadening is interpreted as being solely due to the scattering rate (equation (2)), it yields $\tau_{\mathrm{D}} \approx 3$ ps. On the other hand, if interpreted as being due solely to inhomogeneities (equation (3)), it gives $\bar{x} \approx 0.0017$.

In order to separate the two contributions to T_{Dm}, it is necessary to have an independent gauge of $\bar{x}$. This can be estimated by modelling the width of the superconducting transition using a mean-field approach [17]. This relies solely on measured bandstructure parameters, and the experimental dependence of the superconducting critical temperature T_c on pressure; the only adjustable parameter is $\bar{x}$ [17]. Figure 2(a) shows data from the megahertz coil system; the background T dependence of the empty coil has been subtracted and equation (1) used to obtain $\Delta\lambda$. The superconducting to normal transition is observed as a change from skin-depth to penetration-depth limited coupling of the sample to the megahertz fields [8, 11, 17], which results in a shift in f. Because the mean-field approach predicts a Gaussian broadening of the transition, we fit the differential $\mathrm{d}\Delta\lambda/\mathrm{d}T$ of the data to a Gaussian centred on $T_c \approx 9.37$ K (figure 2(b)); a full width at half-maximum of $\Delta T_c \approx 0.71$ K is obtained. Substituting the value of $\bar{x} \approx 0.0017$ from the SdH data into the formulae of [17] gives $\Delta T_c \approx 0.6$ K, in good agreement with $\Delta T_c \approx 0.71$ K extracted from the megahertz experiments.

The implication of these data is that *virtually all* of the measured T_{Dm} may be accounted for by spatial inhomogeneities [15], and therefore that the true value of τ_{D} is probably much greater than 3 ps. The fact that spatial inhomogeneities are the dominant contribution to T_{Dm} in κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ is in qualitative agreement with cyclotron resonance (CR) experiments on other BEDT-TTF salts [13, 18][6]. CR involves an optically induced 'vertical' transition between Landau levels [13]; the energy of the transition depends on the separation of adjacent Landau levels at *a particular point in space* [14, 15]. Hence, the chief contribution to the CR linewidth is due to the broadening of the Landau levels caused by scattering; there is little or no contribution from spatial inhomogeneities [13–15].

In α-(BEDT-TTF)$_2$NH$_4$Hg(SCN)$_4$, the scattering rate from the CR linewidth was about 0.1 times that from T_{Dm} [13]. Similarly, CR data in β''-(BEDT-TTF)$_2$SF$_5$CH$_2$CF$_2$SO$_3$ [18] indicate a Landau-level broadening about three to four times smaller than that deduced from T_{Dm} [21]. Finally, the magnetoresistance of β-(BEDT-TTF)$_2$IBr$_2$ suggests that spatial inhomogeneities are the dominant contribution to T_{Dm} [22].

We now turn to the intralayer conductivity, $\sigma_\parallel$. Measurements of $\sigma_\parallel$ in layered systems yield another scattering time, τ_c [23]. However, it is necessary to add a note of caution about $\sigma_\parallel$ measurements in organic metals [1, 10]. Whilst interlayer resistivity (ρ_{zz}) measurements are simple [10], the large resistivity anisotropy $\sim 10^3$–10^5 [1] makes quantitative measurements of the intralayer resistivity $\rho_\parallel$ using wires and contacts very difficult [10]; there is often a substantial component of ρ_{zz} present, which will lead to an *underestimate* of $\sigma_\parallel$. Optical reflectivity potentially provides a more reliable method of obtaining $\sigma_\parallel$; however, there are severe problems associated with the use of Drude extrapolations of high-frequency conductivity to zero frequency [24]. We note also that the optical data are chiefly determined by the surface layers of the crystal, which are not perfect. Taking what are probably the best of the optical [24] and conventional [25] measurements yields values of $\sigma_\parallel$ in the range (1–4) $\times 10^5\ \Omega^{-1}\ \mathrm{m}^{-1}$ at $T \approx 12$ K.

[6] In spite of strenuous efforts, it has not been possible to observe CR in κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ [19, 20]. This may be associated with the very small interlayer corrugations of the Fermi surface [4]; the experimental configurations used for CR are sensitive to the interlayer corrugations of the Fermi surface [12, 13].

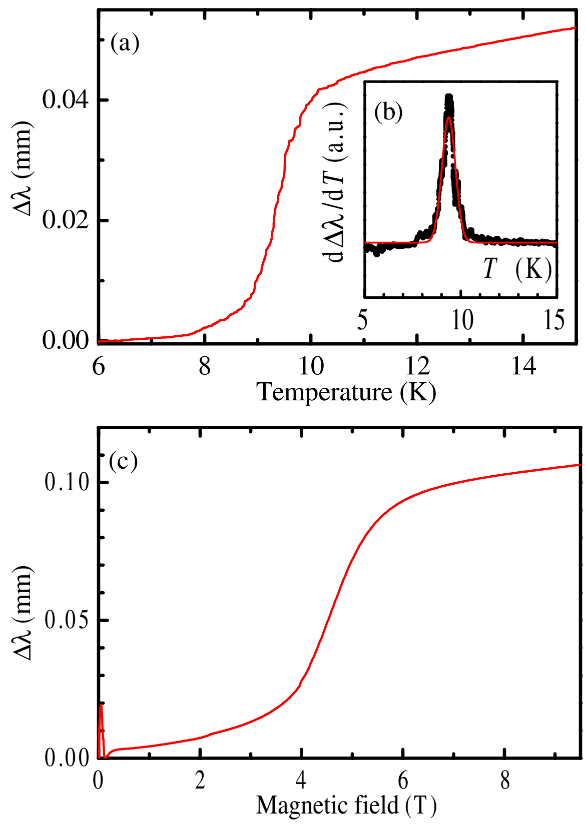

Figure 2. (a) Zero-field megahertz penetration data for a single crystal of κ-(BEDT-TTF)$_2$ Cu(NCS)$_2$, shown as change in penetration $\Delta\lambda$ versus T. (b) The differential $\mathrm{d}\Delta\lambda/\mathrm{d}T$ of the data (points) fitted to a Gaussian (curve) centred on $T_c = 9.37$ K, with a full width at half-maximum of $\Delta T_c \approx 0.71$ K. (c) Change in penetration/skin depth $\Delta\lambda$ versus magnetic field B ($T = 0.7$ K).

In view of the uncertainty surrounding these measurements, we have used the megahertz coil system to estimate $\sigma_\parallel$ of κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ from the behaviour of $\Delta\lambda$ (equation (1))[7] as a function of T and field B. In the superconducting state ($B = 0$, low T), $\Delta\lambda$ is limited by the small intralayer superconducting penetration depth [8] ($\lambda \approx 0.54\ \mu$m; see the summary of data in [1]) which means that very little of the sample is penetrated by the megahertz fields [8, 12]. On entering the normal state, the penetration of the megahertz fields is limited by the skin depth, $\delta_\parallel$ [11, 12].

Figure 2(a) shows $\Delta\lambda$, deduced using equation (1) with $f = 38$ MHz, $R^2 = 0.485$ mm^2 and $r_s = 0.13$ mm, as a function of T; on crossing from superconducting to normal, $\Delta\lambda$ increases to 0.05 mm at 12 K. This implies that the megahertz fields are penetrating a substantial fraction of the sample, so that the conductivity estimated will be characteristic of the *bulk*, rather than the surface, of the sample. Using[8] $\delta_\parallel \approx \Delta\lambda = 0.05$ mm, and the standard relationship $\delta_\parallel = (\pi\sigma_\parallel f\mu)^{-1/2}$ (e.g. [26]), we obtain $\sigma_\parallel = 2.7 \times 10^6\ \Omega^{-1}\ \mathrm{m}^{-1}$, rather higher than the values obtained by other methods [24, 25]. As mentioned earlier, the difference might be due to the contamination of the conventional $\rho_\parallel$ values [25] by ρ_{zz} [10], or, in the case

[7] As $f \approx 38$ MHz $\ll \tau_c^{-1}$, the megahertz skin-depth experiment is effectively a dc measurement of $\sigma_\parallel$ [24].
[8] This is in reasonable agreement with low-temperature values $\delta_\parallel \sim 1\ \mu$m obtained at a frequency $\approx$70 GHz [35]. Scaling by the square root of the frequency gives $\delta_\parallel \sim 0.05$ mm at 30 MHz.

of the optical data, the difficulties with the Drude extrapolation or problems with the sample surface [24].

The values of T_{Dm} quoted earlier were for $T \lesssim 1$ K and at finite B. We therefore obtain another estimate of $\sigma_{\|}$ by using B to drive the sample normal at $T = 0.7$ K. Figure 2(c) shows that $\Delta\lambda$ grows to 0.11 mm, indicating both that $\delta_{\|}$ is around this size, and that the sample ($r_{\mathrm{s}} \approx 0.13$ mm) is almost completely penetrated. Further support for this comes from the fact that no SdH oscillations are observed in $\Delta\lambda$ (compare figure 2(c) with 1), whereas their observation is possible with the megahertz coil system in larger samples [27] or by increasing f. Again using $\delta_{\|} = (\pi\sigma_{\|} f \mu)^{-1/2}$, we obtain $\sigma_{\|} = 6.8 \times 10^5\ \Omega^{-1}\ \mathrm{m}^{-1}$, somewhat smaller than the value deduced at $B = 0$ and 12 K, and similar to, but slightly higher than, the values obtained by other methods [24, 25]. This reduction in $\sigma_{\|}$ in higher fields may be associated with magnetoresistance.

We now use the observed value of $\sigma_{\|}$ to yield an *apparent* value of τ_{c}, *under the assumption that the whole of the sample contributes uniformly to the electrical conductivity*. As far as electrical conductivity is concerned, the Fermi surface of κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ is essentially a slightly distorted cylinder [4, 6] (inset, figure 1). We therefore use the Drude expression [3], $\sigma_{\|} = ne^2\tau_{\mathrm{c}}/m^*$ to infer τ_{c}; here, the quasiparticle density n is given by $n \approx 2eF_\beta/ha = 1.17 \times 10^{27}\ \mathrm{m}^{-3}$, where $F_\beta = 3920$ T is the magnetic quantum oscillation frequency of the β breakdown orbit, $m^* \approx 6.5m_{\mathrm{e}}$ is its effective mass [4, 5] and $a = 16.248$ Å is the interlayer lattice parameter [1]. This yields $\tau_{\mathrm{c}} \approx 0.14$ ps at $T = 0.7$ K and $B \approx 6$ T, and $\tau_{\mathrm{c}} \approx 0.56$ ps at $T = 12$ K and $B = 0$. (Note that the values for $\sigma_{\|}$ from [24, 25] give even smaller values of τ_{c}.) Qualitatively similar results were obtained on other samples of κ-(BEDT-TTF)$_2$Cu(NCS)$_2$.

Taken together, the SdH and megahertz experiments, carried out on the same, high-quality crystals of κ-(BEDT-TTF)$_2$Cu(NCS)$_2$, suggest that the dominant contribution to the Landau-level broadening is due to spatial inhomogeneities; this is similar to findings in three other BEDT-TTF salts [13, 18, 22] and in semiconductor heterostructures [15, 23]. The experiments therefore imply that the true $\tau_{\mathrm{D}} \gg 3$ ps. By contrast, the skin-depth experiments in κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ yield an apparent scattering time $\tau_{\mathrm{c}} \approx 0.14$ ps under the same conditions of field and temperature, if it is assumed that the whole crystal contributes to the metallic conductivity. At $B = 0$ and $T = 12$ K, similar assumptions yield $\tau_{\mathrm{c}} \approx 0.56$ ps. Even allowing for experimental uncertainties, there appears to be a large difference between τ_{D} and the apparent[9] τ_{c}.

The result $\tau_{\mathrm{D}} \gg \tau_{\mathrm{c}}$ contradicts the expectations of simple LFLT [2, 3, 23]. A quasiparticle will be removed from a Landau-level eigenstate by *any* scattering event, small- or large-angle [23]. Conversely, small-angle scattering events hardly affect the conductivity, because they cannot randomize a quasiparticle's excess forward momentum [3, 23]. In the most thoroughly studied systems of reduced dimensionality, semiconductor heterostructures and Si MOSFETs [15, 23], it is found that $\tau_{\mathrm{c}} \approx \tau_{\mathrm{D}}$ in systems in which large-angle scattering dominates. On the other hand, in systems in which small-angle scattering is pre-eminent, $\tau_{\mathrm{c}} \gg \tau_{\mathrm{D}}$ [23]. The case observed here, $\tau_{\mathrm{D}} \gg \tau_{\mathrm{c}}$, has never been observed in either system.

However, as has been remarked above, the other low-temperature properties of κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ appear to conform to LFLT [1]. Firstly, the SdH oscillations are apparently describable in terms of a Fermi surface with well-defined topology, populated by quasiparticles with a measurable effective mass [1]. Secondly, ρ_{zz} [1] and $\rho_{\|}$ [24] follow

[9] A similar observation about the discrepancy between the scattering times from T_{Dm} and $\sigma_{\|}$ was made for the organic superconductor β-(BEDT-TTF)$_2$IBr$_2$ [6, 28]. However, $\sigma_{\|}$ data in [28] should be treated with caution, as they were based on a contact geometry liable to yield an admixture of $\rho_{\|}$ and ρ_{zz} [10].

a T^2 dependence at $T \lesssim 40$ K, suggestive of a LFLT description[10]. Instead, it is more likely that the discrepancy between τ_c and τ_{Dm} is suggesting that κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ is *not* a homogeneous metal. Following this idea, we propose that the samples are made up of interconnected metallic domains which obey LFLT, separated by insulating or electrically inactive regions. In a conventional resistance measurement such as that shown in figure 1, these metallic regions will 'short out' the insulating parts of the sample, completely dominating the electrical transport. As long as the insulating regions remain electrically inactive at low temperatures (say $T \lesssim 40$ K), resistivity data will therefore be *qualitatively* in agreement with the expectations of LFLT; the measured resistivity will be proportional to T^2 and the SdH oscillations will obey the expected temperature dependence [1]. Similarly, quantities such as T_{Dm} and $\bar{x}$ deduced from such data will reflect *only the properties of these metallic domains*; the finite value of $\bar{x}$ obtained above shows that these regions are themselves somewhat inhomogeneous. By contrast, a *quantitative* measurement of the electrical conductivity will give a value that is smaller than if the whole sample were metallic. Consequently, the apparent scattering time τ_c deduced under the latter assumption will be too small, in agreement with the data reported in this letter.

A clue as to the origins of the coexisting metallic and insulating domains in κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ is given by the phase diagram for the κ-(BEDT-TTF)$_2$X series of organic metals in figure 3, compiled from recent data [29–32]. On cooling, these materials first undergo a 'glassy transition' at $T \approx 70$ K; this appears to be associated with freezing-in of disorder in the orientational degrees of freedom of the terminal ethylene groups of the BEDT-TTF molecules (see [30] and references therein). Subsequently, below $T = T^* \approx 40$ K, magnetization data indicate the coexistence of a density-wave-like phase and a paramagnetic metal [31]. Sasaki *et al* [31] suggest that this observation is due to the nesting of the quasi-one-dimensional (Q1D) sheets (figure 1) of the Fermi surface (giving the density-wave), whilst the α pocket remains ungapped (resulting in the metallic behaviour).

However, the interpretation of [31] *cannot* explain the smooth evolution with field of magnetic breakdown oscillations (due to tunnelling between the α pocket and Q1D sheets) that is observed in κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ [5]; such data may *only* be understood in terms of the unnested Fermi surface shown in figure 1 [5]. Instead, we suggest that the glassy transition represents κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ forming two distinct types of domain, each characterized by one of the two different possible configurations ('staggered' or 'eclipsed') of the terminal ethylene groups of the BEDT-TTF molecules [30]. The two possible arrangements will cause slight differences in Fermi-surface topology, leaving one type of domain with a Fermi surface prone to nest (probably at $T \approx T^*$), and the other with a Fermi surface that remains unnested down to low temperatures; models of κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ have shown that very small differences in Fermi-surface topology can affect the degree of nesting dramatically (see [33] and references therein).

The remaining question to be answered is the typical lengthscale of the distinct metallic and insulating regions. The data reported in this paper suggest that the true quasiparticle scattering time in the metallic regions is $\gg \tau_{Dm}$; using the typical Fermi velocity [4], this translates to a mean-free path of $\gg 0.2$ μm, representing a minimum size for the metallic domains. On a much shorter lengthscale within these metallic regions, a small fraction ($\bar{x} = 0.0017$) of the cation or anion molecules of κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ appears to be in some way defective [17], leading to the inhomogeneous broadening of the Landau levels which is the chief contribution to the measured value of T_{Dm}.

[10] The observation that *both* ρ_{zz} and $\rho_{\|}$ follow almost the same T dependence up to 40 K [24] is actually quite a surprise, given that the interlayer transfer integral is $\sim$0.5 K [4].

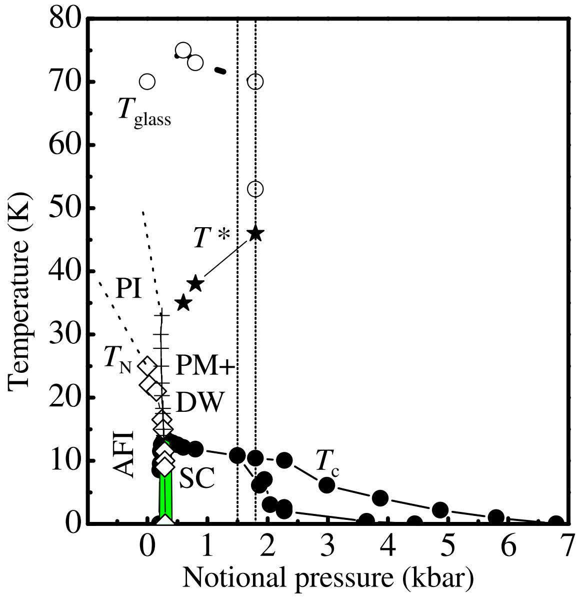

Figure 3. Phase diagram of κ-(BEDT-TTF)$_2$X, including boundaries suggested by recent data. AFI = antiferromagnetic insulator below T_N (◇) [29]; PI = paramagnetic insulator [29]; PM = paramagnetic metal; DW = proposed density wave below T^* (⋆) [31]; SC = superconductivity below T_c (●) [32]; T_{glass} (○) = proposed glassy structural transition [30]. 'Notional pressure' combines chemical pressure caused by changing anion X [30, 31] and applied hydrostatic pressure [32]; '0' is ambient pressure for κ-(BEDT-TTF)$_2$Cu[N(CN)$_2$]Cl; the vertical lines are the ambient pressure positions of deuterated (left) and undeuterated (right) κ-(BEDT-TTF)$_2$Cu(NCS)$_2$; note that T_c has a different pressure dependence for these two salts [32].

In summary, we have measured the apparent scattering times in κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ using a variety of techniques. A large discrepancy is encountered with the expectations of LFLT, suggesting that only a small volume fraction of (otherwise very high quality) κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ samples contribute to the observed in-plane conductivity. This might be an explanation for the great difficulty in observing de Haas–van Alphen oscillations in κ-(BEDT-TTF)$_2$Cu(NCS)$_2$; the signals seem much weaker than those from samples of other BEDT-TTF salts [34]. Our data are most easily interpreted in terms of a model in which the crystals comprise interconnected 'metallic' and 'insulating' domains, which originate at the recently observed glassy transition at around 70 K. In view of this conclusion, the assumption that κ-(BEDT-TTF)$_2$Cu(NCS)$_2$ is a homogeneous 'model' layered superconductor [7] should be questioned. Moreover, the observation of the glassy transition in other κ-(BEDT-TTF)$_2$X salts (figure 3) may indicate that they too suffer from similar inhomogeneities.

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Note that the electron density used in this paper is erroneous
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