

Influence of carrier lifetime on quantum criticality and superconducting T_c of $(\text{TMTSF})_2\text{ClO}_4$ Abdelouahab Sedeki,¹ Pascale Auban-Senzier,² Shingo Yonezawa,³ Claude Bourbonnais,⁴ and Denis Jerome²¹*Université Dr Tahar Moulay Saïda, B.P. 138 cité ENNASR 20000, Saida, Algeria*²*Laboratoire de Physique des Solides (UMR 8502), Université Paris-Sud, 91405 Orsay, France*³*Department of Physics, Graduate School of Science, Kyoto University, Kyoto 606-8502, Japan*⁴*Regroupement Québécois sur les Matériaux de Pointe and Institut Quantique, Département de physique, Université de Sherbrooke, Sherbrooke, Québec, Canada J1K-2R1*

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This work presents and analyzes electrical resistivity data on the organic superconductor $(\text{TMTSF})_2\text{ClO}_4$ and their anion-substituted alloys $(\text{TMTSF})_2(\text{ClO}_4)_{1-x}(\text{ReO}_4)_x$ along the least conducting c^* axis. Nonmagnetic disorder introduced by finite size domains of anion ordering on non-Fermi-liquid character of resistivity is investigated near the conditions of quantum criticality taking place at $x = 0$. The evolution of the T -linear resistivity term with anion disorder shows a limited decrease in contrast with the complete suppression of the critical temperature T_c as expected for unconventional superconductivity beyond a threshold value of x . The resulting breakdown of scaling between both quantities is compared to the theoretical predictions of a linearized Boltzmann equation combined to the scaling theory of umklapp scattering in the presence of disorder-induced pair breaking for the carriers. The theory shows that quantum antiferromagnetic fluctuations, which are at the core of unconventional Cooper pairing and inelastic scattering of T -linear resistivity in these systems, are weakly affected by disorder while the phase coherence responsible for a finite T_c can be completely suppressed.

DOI: [10.1103/PhysRevB.98.115111](https://doi.org/10.1103/PhysRevB.98.115111)**I. INTRODUCTION**

Quantum criticality that becomes unstable against the emergence of superconductivity is a common feature of most recent unconventional superconductors. It happens when the ordering temperature of an antiferromagnetic (AFM) phase, which can be metal-like, semimetallic, or insulating, ends towards zero temperature, as a control parameter doping or pressure is varied [1]. Such is the situation encountered in the vast family of pnictide superconductors [2,3]. That may also be the situation met in hole-doped cuprates, although there the quantum critical point (QCP) corresponding to the termination at zero temperature of the pseudogap phase is beneath the superconducting dome of their phase diagram [4]. This is also found in electron-doped cuprates and heavy-fermion superconductors, which both exhibit an AFM phase terminating toward zero temperature when the electron doping level or pressure is varied respectively in the former [5] or latter compounds [6,7]. The proximity of an AFM Mott-insulating phase to superconductivity is also present in layered organic superconductors bearing much resemblance with the phase diagram of cuprates [8–10].

Most experimental investigations performed near QCP conditions have revealed that several physical properties deviate significantly from the canonical Fermi liquid behavior, in particular the resistivity which instead of the expected T^2 dependence exhibits a linear- T dependence [11], a puzzling feature shared by essentially all the aforementioned unconventional superconductors close to a magnetic QCP.

Moreover, there exists another class of materials displaying both antiferromagnetism and superconductivity that has not been as much highlighted, namely, the quasi-one-dimensional

(Q1D) Bechgaard organic superconductors [12,13]. Although superconductivity in heavy fermions [14] was discovered before that in Q1D organics, it is in these latter materials that superconductivity was first revealed following the suppression of an insulating phase under pressure which was subsequently identified as a spin-density-wave (SDW) state [15–17]. A strong argument in favor of Q1D superconductors is the relative simplicity of their electronic structure, comprising a single nearly flat Fermi surface as well as a temperature-pressure phase diagram which can be considered as a textbook example of magnetic quantum criticality that becomes unstable to the formation of superconductivity [18,19].

An unusual behavior of metallic transport in $(\text{TMTSF})_2\text{PF}_6$ under pressure at low temperature has been observed in the first studies [15] related to a pronounced T -linear tendency between 0.1 and 10 K and even a downward curvature of the resistivity in the pressure range corresponding to the vicinity of the QCP [20]. While superconducting fluctuations were first considered as a possible interpretation in the three-dimensional (3D) temperature regime very close to T_c where the transverse coherence length is larger than the interstack distance, paraconductivity from superconducting origin failed to account for the correct temperature dependence observed up to 10 K or so [20]. Thirty years have been required to clarify the problem when similar behaviors of transport have been detected in other superconductors close to a QCP, in particular in cuprates and pnictides [4,11].

As far as organic superconductors are concerned, detailed investigations of transport at low temperature were conducted first on $(\text{TMTSF})_2\text{PF}_6$ under pressure [21]. In this case the temperature dependence of the inelastic scattering in

conditions very close to the QCP is a clear-cut behavior. $\rho(T)$ follows a T -linear law over a decade in temperatures above T_c , already visible on a log-log plot of the resistivity. The behavior is becoming quadratic as expected in a Fermi liquid above 10 K [21]. However, as the location of the compound in the T - P phase diagram is moved away from the QCP conditions the resistivity acquires a quadratic component even at low temperature and becomes fully quadratic when T_c is suppressed. The log-log plot of the resistivity reveals a power law $\Delta\rho \propto T^\beta$ with an exponent β evolving from 1 to 2 between 11.8 and 20.8 kbars [21]. Subsequently, another procedure has been followed to analyze the inelastic scattering of $(\text{TMTSF})_2\text{PF}_6$, namely, a sliding fitting procedure [22] with a second-order polynomial form such as $\rho(T) = \rho_0 + A(T)T + B(T)T^2$. Here, the fitting procedure enables the detection of a possible temperature dependence of the A and B prefactors, whereas the value of ρ_0 depends on pressure only [22]. This procedure has shown that the T -linear contribution is dominant below 10 K compared to the Fermi liquid contribution close to the QCP, but otherwise becomes of the order of the quadratic contribution.

The same log-log analysis failed to provide a clear-cut picture in $(\text{TMTSF})_2\text{ClO}_4$, but the polynomial fitting procedure conducted on the resistivity along c^* led to conclusions that A and B are only weakly temperature dependent in this compound between T_c and 15 K (though still pressure dependent) [23]. The detailed analysis of the temperature-dependent resistivity over an extended range of pressures has shown that $\rho_{c^*}(T)$ can thus be fitted to the polynomial expression $\rho_{c^*0} + AT + BT^2$ where all parameters are evolving under pressure but the experimental feature, $A \rightarrow 0$ as $T_c \rightarrow 0$, is preserved [23], as found in other materials like the cuprates for instance [24].

Although a precise functional relation between A and T_c is uncertain, the salient result of these pressure studies is the existence of a T -linear contribution to the resistivity being finite only when T_c is finite, suggesting in turn the existence of a common origin for pairing and the T -linear inelastic scattering [22].

While there is no general consensus about the microscopic origin of T -linear resistivity appearing near a QCP [25], theoretical efforts have been displayed to shed some light on its origin in the context of the Q1D Bechgaard salts [19,26,27]. The electron-electron scattering rate, as extracted from the imaginary part of the one-particle self-energy, has been derived by a renormalization group procedure. The calculation that takes into account the quantum interference between electron-electron and electron-hole scattering channels was able to show in the first place the emergence of unconventional (d -wave) superconductivity from quantum SDW fluctuations [28–30], in a way that goes beyond the standard spin fluctuation scenarios in which both scattering channels are not dynamically coupled [31–33]. In the second place, it allowed the possibility to link the strength of the linear resistivity term to the one of pairing as a function of pressure [19]. Another strategy has been to connect the solution of the transport Boltzmann equation to the same renormalization group approach to nonconserving momentum scattering amplitudes. This procedure has confirmed the existence of a T -linear term in resistivity and the correlation between its strength

and the size of superconductivity pairing mediated by SDW fluctuations [27].

These findings have motivated the present work devoted both experimentally and theoretically to the influence of a finite elastic lifetime on the T -linear resistivity in the Bechgaard salts. On the theoretical side impurity scattering, assumed to be absent so far in the renormalization group calculations of the scattering amplitudes, is a source of pair breaking for both unconventional superconductivity and density-wave correlations. We show that this affects to a certain extent the constructive interference between both types of correlations at low energy. Although this can be sufficient for the complete suppression of phase coherence responsible for a finite T_c , it turns out that SDW correlations that build up at higher energy are less affected by disorder-induced pair breaking and remain relatively strong in amplitude. As the source of umklapp scattering for carriers in the Boltzmann equation, these correlations sustain to a great extent the non-Fermi-liquid character of resistivity in spite of the scaling breakdown between the T -linear term and T_c .

On experimental side, previous investigations have shown that T_c in the $(\text{TMTSF})_2X$ series is strongly affected by disorder of nonmagnetic origin [34–37], providing in turn some hint for the existence of a nonconventional superconducting gap displaying both signs over the Fermi surface in this spin singlet superconductor [38], pointing for instance to a d -wave (or g -wave) type of superconductivity. As the linear term of transport at low temperature was found to be related to T_c , we feel it is legitimate to test the robustness of the established connection between A and T_c when T_c can be modified by other means but pressure, namely, by the influence of nonmagnetic disorder [39]. It is worth mentioning that in Sr_2RuO_4 , an equal sensitivity to nonmagnetic disorder has been observed [40], suggesting that spin triplet pairing is a possible candidate for the superconductivity (SC) ground state in this material [41].

Fortunately, $(\text{TMTSF})_2\text{ClO}_4$ is the unique system among other superconductors in the $(\text{TMTSF})_2X$ series in which structural disorder can control the elastic scattering time. This property is based on the peculiarity of the tetrahedral symmetry of the anion ClO_4^- .

The structure of $(\text{TM})_2X$ salts belongs to the triclinic space group $P\bar{1}$ with every anion site located on an inversion center. Hence, as long as the local symmetry of anions like PF_6^- is centrosymmetric, their actual orientation always fits the overwhole crystal symmetry. Thus, no additional disorder is introduced. However, noncentrosymmetric anions such as tetrahedral ClO_4^- or ReO_4^- may have their oxygen atoms pointing towards methyl groups or Se atoms of one or the other neighboring donor molecules [42,43]. These (at least two) possible directions lead in turn to a potential source of disorder. While there is no disorder on average at elevated temperature since these tetrahedral anions are rotating as shown by NMR data [44], the minimization of entropy due to the reduction of degrees of freedom in thermodynamical equilibrium triggers an anion ordering at low temperature with the concomitant occurrence of a superstructure. In addition, the relative orientation at low temperature of neighboring anions depends on the nature of the anions. For $(\text{TMTSF})_2\text{ReO}_4$,

the orientation of ReO_4^- alternates along the three directions providing a periodic lattice distortion of wave vector $(1/2, 1/2, 1/2)$ (in units of reciprocal lattice vector in each direction) with a concomitant metal insulator transition below $T_{\text{AO}} = 176$ K [45], whereas the alternation of ClO_4^- occurs solely along the b axis in $(\text{TMTSF})_2\text{ClO}_4$ below $T_{\text{AO}} = 24$ K, leading in turn to the $(0, 1/2, 0)$ order [46,47]. Such a doubling of periodicity along b is thus responsible for the doubling of the Fermi surface folded along b^* in slowly cooled samples. These samples retain in turn metallic properties at low temperature.

A possible means of introducing controllable disorder in an otherwise well-ordered $(\text{TMTSF})_2\text{ClO}_4$ crystal is the substitution of ClO_4^- by ReO_4^- anions. The consequences of the isostructural anion alloying on the electronic properties such as superconductivity were revealed long ago [35], even though the actual structural reasons were not elucidated. For instance, ReO_4^- replacing ClO_4^- in $(\text{TMTSF})_2\text{ClO}_4$ affects T_c so severely that a 1% ReO_4^- concentration is large enough to shift T_c from 1.3 to 0.9 K and 10% suppresses T_c totally [35]. Additional investigations of structural and electronic properties of $(\text{TMTSF})_2(\text{ClO}_4)_{(1-x)}(\text{ReO}_4)_x$ have been conducted subsequently [36,48]. What structural studies have revealed in the ClO_4^- alloys at low concentration of ReO_4^- is the existence of domains in which the $(0, 1/2, 0)$ ClO_4^- order persists, the size of them controlling the carrier elastic mean free path [49].

Another way to create defects in $(\text{TMTSF})_2\text{ClO}_4$ is to beat the slow anion orientation process by rapid cooling of the sample, as first detected by NMR, specific heat, EPR, and transport measurements [50–52]. A recent reinvestigation of magnetic and transport properties of $(\text{TMTSF})_2\text{ClO}_4$ under a well-controlled cooling procedure in the vicinity of the anion ordering temperature has revealed at increasing cooling speed above 1 K/min a crossover from a rather homogenous localized disorder to a granular situation in which anion-ordered puddles are embedded in an anion-disordered background [53]. This picture is reminiscent of the present situation in alloys where the carrier mean free path of the well-ordered grains is also limited by their own size.

The fact that the elastic mean free path of $(\text{TMTSF})_2\text{ClO}_4$ can be controlled by disorder is a remarkable property of the material and makes $(\text{TMTSF})_2\text{ClO}_4$ a unique system in which one can check simultaneously the effect of a change of the mean free path on T_c and on quantum criticality behavior possibly related to the onset of superconductivity [21].

The present report attends to fulfill this goal with a quantitative study of both superconducting and metallic state properties in two situations where disorder can be introduced with a concomitant control of the mean free path: first, via studying lightly ReO_4^- substituted $(\text{TMTSF})_2\text{ClO}_4$ samples, and second, studying the fast cooling of pure $(\text{TMTSF})_2\text{ClO}_4$. Theoretically, we broaden the renormalization group (RG) approach of the Q1D electron gas model to include impurity pair-breaking effects on the renormalization of umklapp scattering that enters the linearized Boltzmann equation of electrical transport.

Our study will show that the correlation between the non-Fermi-liquid resistivity and T_c , which is well established under pressure [21], breaks down when T_c in this d -wave superconductor is suppressed by nonmagnetic disorder. This

result emphasizes the strong pair-breaking role of a limited elastic mean free path on T_c as opposed to the much weaker influence on antiferromagnetic fluctuations, which are believed to be a major ingredient for pairing in this superconductor. The positive connection between theory and experiment widens the consistency of the model of magnetically driven mechanisms for quantum criticality and superconductivity in the Bechgaard salts.

In Sec. II, the low-temperature transverse resistivity measurements of $(\text{TMTSF})_2(\text{ClO}_4)_{(1-x)}(\text{ReO}_4)_x$ alloys are presented and their superconducting and non-Fermi-liquid properties described. Section III is devoted to the results of the numerical solution of the linearized Boltzmann equation of resistivity for the Q1D electron gas model with umklapp scattering and pair breaking, as yielded by the renormalization group method detailed in the Appendix. In Sec. IV, we discuss the results and their connection with theory. In Sec. V, we summarize and conclude this work.

II. EXPERIMENT

The present study is based on transport measurements of $(\text{TMTSF})_2(\text{ClO}_4)_{(1-x)}(\text{ReO}_4)_x$ single crystals grown electrochemically by Bechgaard in Copenhagen. All data have been taken at Orsay except for one $(\text{TMTSF})_2\text{ClO}_4$ sample already measured in a previous study at Kyoto [53] for its superconducting properties.

The measurement of the transverse resistivity ρ_{c^*} along the least conducting c^* direction has been privileged because of the requested quantitative comparison between samples with different ReO_4^- concentration. The frequent cracks occurring on cooling for the resistivity along a or b axes prevent any comparison between samples whereas ρ_{c^*} is known to be less influenced by such cracks even after multiple cooling processes [54]. In addition, magnetoresistance measurements have shown that the band theory should apply below 10 K as also supported by the existence of a Drude edge along the c^* axis in the same temperature range [55]. A previous comparative study of transport along a and c^* has shown that in spite of a very large anisotropy, electron scattering times measured along a or c^* exhibit a similar temperature dependence at least up to 30 K [23].

For $(\text{TMTSF})_2\text{ClO}_4$ samples studied in Kyoto, the resistivity was measured with the resistivity option of a PPMS with a dc current of 10 μA reversed to cancel thermoelectric voltages. For $(\text{TMTSF})_2(\text{ClO}_4)_{(1-x)}(\text{ReO}_4)_x$ studied in Orsay, a lock-in amplifier technique was used with an ac current of 10 μA .

One of the major difficulties in interpreting low-temperature transport in most $(\text{TMTSF})_2X$ superconductors is the existence of a negative curvature in the $R(T)$ curve approaching T_c below 4 K [15]. This negative curvature seems to be fairly robust as it can be observed even when SC is suppressed either by a magnetic field or in alloyed samples [56]. While precursors due to SC could contribute to a possible collective paraconductive contribution up to about $1.2T_c$ [56], there is still much uncertainty as to the origin of paraconductivity up to 3.5–4 K remaining a pending problem. Since precursor effects are not discussed in the present article the analysis of the $\rho_{c^*}(T)$ curves has been limited to the range

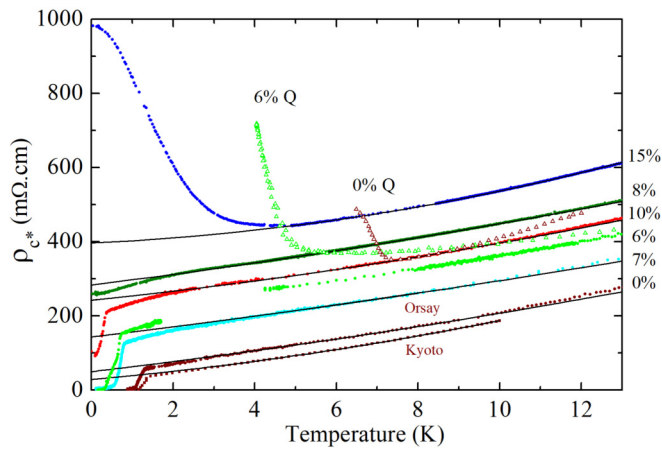


FIG. 1. Interlayer resistivity data for the $(\text{TMTSF})_2(\text{ClO}_4)_{(1-x)}(\text{ReO}_4)_x$ solid solutions all measured in the relaxed state and normalized to $28 \Omega \text{ cm}$ at ambient temperature. Samples have been labeled according to their nominal ReO_4^- concentration which may deviate in some cases from the actual defects concentration given by an electron microprobe analysis which in turn agrees with the value of the residual resistivity. The clean sample measured at Kyoto displayed a slightly smaller residual resistivity ($30 \text{ m}\Omega \text{ cm}$) than the Orsay one ($50 \text{ m}\Omega \text{ cm}$). This is in line with $T_c = 1.35 \text{ K}$ at Kyoto while T_c is only 1.18 K in the Orsay sample. We also show the data for the quenched 6% sample. The data above 8 K or so may be affected by some annealing occurring on warming up [57]. The clean but quenched $(\text{TMTSF})_2\text{ClO}_4$ data have been taken from Ref. [58].

4–10 K since our purpose is a study of the single-particle non-Fermi-liquid properties only.

The study of transport has been conducted in the alloy series $(\text{TMTSF})_2(\text{ClO}_4)_{(1-x)}(\text{ReO}_4)_x$ where the size of the ordered regions is controlled by the amount of the ReO_4^- substituent. The resistivity data are displayed in Fig. 1. We also emphasize that the samples in Fig. 1 have been labeled according to the nominal concentration of ReO_4^- which may differ significantly from the results obtained in an electron microprobe analysis [59].

The resistivity of $(\text{TMTSF})_2\text{ClO}_4$ is *at variance* with that of $(\text{TMTSF})_2\text{PF}_6$ since no clear T -linear resistivity behavior is emerging from the raw data of $(\text{TMTSF})_2\text{ClO}_4$ [60].

Note that every sample in Fig. 1 differs by the behavior of its resistivity at low temperature. Both 0% samples exhibit a complete superconducting transition although their T_c and their residual resistivity ρ_{c^*0} are slightly different. The SC transition is still complete in the 7% sample. We also notice a complete transition for the 6% sample (green curve in Fig. 1) but the broadening at low temperature is a signature of a proximity effect between superconducting puddles [53]. A further depression of T_c is observed in the 10% sample, but there the finite zero temperature value of the resistivity suggests the existence of superconducting islands far from each other precluding global SC coherence through the proximity effect. The 8% indicates the absence of any transition at finite temperature. The value of its ρ_{c^*0} appears to be reliable, although a tiny break in the $\rho_{c^*}(T)$ data (not visible in Fig. 1) occurring around 5 K made the polynomial analysis of the

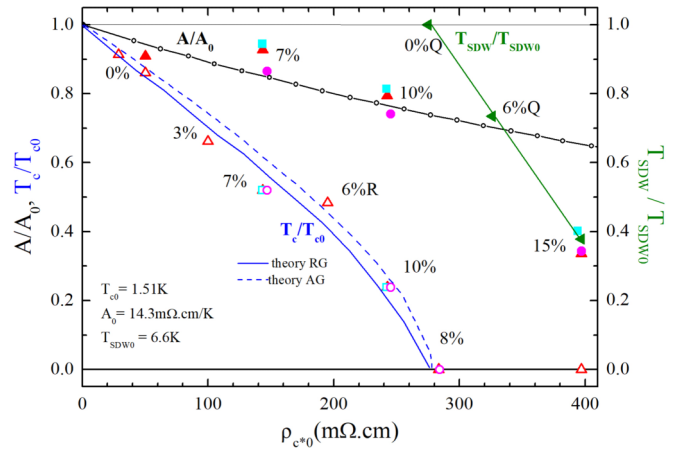


FIG. 2. Normalized value of A from the two conductor model (open circles) versus the residual resistivity and from the fits with different fitting ranges (red triangles, 4.5–10 K; blue squares, 5–9 K; magenta circles, 5–10 K). A has been normalized to the value linearly extrapolated to $\rho_{c^*0} = 0$. T_c for the onset of superconductivity from experiment and theories and T_{SDW} versus ρ_{c^*0} . Data for the 0%Q are provided by Ref. [58].

temperature dependence less reliable. In the 15% sample a metal-insulator transition is observed at 2.5 K which can be related to a SDW state occurring in the anion-disordered sample after fast cooling [35]. Similarly, the sample with a nominal concentration of 17% (not shown in Fig. 1) displays the onset of a SDW at 3.9 K .

The residual resistivity ρ_{c^*0} can be extracted relatively easily at different ReO_4^- concentrations from the data in Fig. 1, as it is only weakly dependent on the determination of the other prefactors in the polynomial fit (*vide infra*), but the correct analysis of the inelastic temperature dependence is a very delicate operation even in pure samples as already discussed in several previous publications [21,23].

The present study utilizes a polynomial fitting procedure keeping A and B fixed within the fitting window for each sample in Fig. 1. As it is difficult to put error bars on such fits we have privileged the data coming from three different fitting windows between 4.5 and 10 K leading to values for ρ_{c^*0} , A , and B from the temperature dependence in Fig. 1 which are displayed in Fig. 2. The fitting window has been adequately chosen in order to avoid the onset of the as yet uncharacterized paraconductive contribution at low temperature and the proximity of the anion ordering at high temperature [61]. Transition temperatures defined by the onset temperatures have been normalized to our highest observed transition for superconductivity, namely, 1.51 K , and to 6 K for the SDW transition as observed in the case of strong anion disorder [62]. All samples in Fig. 2 refer to the very slowly cooled regime except for the 6% which has been measured both under relaxed and quenched conditions (see Sec. IV for a further discussion). The data of A for the 15% sample are departing from the slow decrease with x observed at lower concentrations. This may actually be the consequence of the resistivity upturn due to the SDW phase leading in turn to an underestimation of A using the present fitting window.

The x axis of Fig. 2 refers to the residual resistivity as shown in the fits of Fig. 1 and not to the actual scattering rate. We do not expect this resistivity to be directly related to the elastic scattering rate whenever the conducting medium departs from homogeneity. This is likely true for the present data at high ReO_4^- concentrations when the superconducting transition becomes incomplete at, say, 10%.

In the solid solution $(\text{TMTSF})_2(\text{ClO}_4)_{(1-x)}(\text{ReO}_4)_x$ two different anion orders are competing according to x-ray diffuse scattering experiments [48,49]. The $(1/2, 1/2, 1/2)$ ReO_4^- ordering is responsible for a metal-insulator transition at 176 K in $(\text{TMTSF})_2\text{ReO}_4$ [45] while the $(0, 1/2, 0)$ ClO_4^- ordering at 24 K folding the Q1D Fermi surface leads to a metallic phase at low temperature with a reduced value of the residual resistivity [53]. At concentrations for x between zero and 1, both anion orders are competing, but on the ClO_4^- rich side of the solid solution $(0, 1/2, 0)$ long-range order of ClO_4^- persists up to about $x = 3\%$ [48,49]. Increasing the ReO_4^- content further, the ClO_4^- ordering becomes short ranged above 7% or so and a SDW ground state is stabilized instead of superconductivity [35]. The picture which can be inferred from the x-ray analysis is thus a homogenous medium with local impurities at very low ReO_4^- concentration only becoming heterogenous above 3% where each impurity is surrounded by a disordered volume. The mean free path has been shown to be very large, of the order of 1600 nm in a pristine sample [53,63]. As the size of the ordered domains in alloys is likely to be much smaller than 1600 nm, we may assume that the mean free path in alloys is determined by the typical domain size. The predominant role of the domain size limiting the mean free path is also corroborated by a similar study performed on the solid solution $(\text{TMTSF})_2(\text{AsF}_6)_{(1-x)}(\text{SbF}_6)_x$ when the symmetry of the anions is central, which has shown that even a 15% concentration of SbF_6 has no detectable effect on T_c [64], being *at variance* with the present alloying effect in $(\text{TMTSF})_2\text{ClO}_4$.

The Abrikosov-Gor'kov theory [65] of a conventional superconducting T_c in the presence of pair breaking has also been reported in Fig. 2 together with the results of the renormalization group theory (Sec. III and the Appendix). We observe a satisfactory agreement between the onset temperatures and theories when the residual resistivity is considered as a representative of the scattering rate [39]. Such an agreement, however, may be surprising as the residual resistivity is the macroscopic resistivity of the entire sample which we found to become heterogenous at high ReO_4^- concentration, while T_c is the onset transition temperature of well-ordered domains. Consequently, this agreement between experiment and theories infers in turn the existence of some connection, $\rho_{c=0} \propto 1/\lambda$, between the mean free path λ determined by domain size entering the theory and the macroscopic resistivity.

III. THEORY: RENORMALIZED LINEARIZED BOLTZMANN EQUATION FOR RESISTIVITY

In this section we sketch out the Boltzmann-RG approach to the calculation of electric resistivity in the presence of pair breaking. The calculation combines the linearized Boltzmann equation with the renormalization group method in the framework of the Q1D electron gas model for the Bechgaard

salts [27]. Prior to doing so, it is important to draw attention on the fact that on experimental grounds, both longitudinal ($\Delta\rho_d$) and transverse ($\Delta\rho_{c^*}$) inelastic contribution to resistivity, though differing by a factor of $\sim 10^4$ in amplitude for a compound like $(\text{TMTSF})_2\text{ClO}_4$, exhibit essentially the same temperature dependence below 30 K or so [23]. Such a concurrence indicates that in this temperature range, they are apparently governed by the same scattering mechanism. In the following we then focus on the calculation of longitudinal resistivity for an array of weakly coupled chains in the *Fig4* plane perpendicular to the c direction.

We first consider the Boltzmann equation of the Fermi distribution function f_k ,

$$\left[\frac{\partial f_k}{\partial t} \right]_{\text{coll}} = e\mathcal{E} \cdot \nabla_{\hbar\mathbf{k}} f_k, \quad (1)$$

which stands for coherent carriers of charge e coupled to an external static electric field $\mathcal{E} = \mathcal{E}\hat{a}$. The collision term on the left depends on the interparticle (umklapp) scattering and collisions to impurities or defects. The linearization of the equation standardly proceeds by looking at small dimensionless deviations ϕ_k to the equilibrium Fermi distribution function $f_k^0 = 1/(e^{\beta\epsilon_k^p} + 1)$, which in leading order yields the small variation $\delta f_k \simeq f_k^0(1 - f_k^0)\phi_k$. Here

$$\epsilon_k^p = \hbar v_F(p\mathbf{k} - k_F) - 2t_{\perp} \cos k_{\perp} d_{\perp} - 2t'_{\perp} \cos 2k_{\perp} d_{\perp} \quad (2)$$

is the electron spectrum of the Q1D electron gas model, where v_F is the Fermi velocity for right- ($p = +$) and left-moving ($p = -$) carriers along the chain, $\pm k_F$ are the one-dimensional (1D) Fermi points in the absence of interchain hopping t_{\perp} , the second nearest-neighbor interchain hopping term of amplitude t'_{\perp} acts as an antineesting term that simulates the effect of pressure in the model, and d_{\perp} is the interchain distance along the b direction.

The linearization of (1) then yields

$$-\left[\frac{\partial \phi_k}{\partial t} \right]_{\text{coll}} = \sum_{k'} \mathcal{L}_{kk'} \phi_{k'} = e\beta\mathcal{E} \cdot \mathbf{v}_k \quad (3)$$

where \mathbf{v}_k is the carrier velocity in the $\hat{\mathbf{k}}$ direction. The collision operator is given by

$$\begin{aligned} \mathcal{L}_{kk'} &= \mathcal{L}_{kk'}^u + \mathcal{L}_{kk'}^{\text{imp}} \\ &= \frac{(\pi \hbar v_F)^2}{(LN_P)^2} \sum_{k_2, k_3, k_4} \frac{1}{2} \left| g_3(\mathbf{k}_F^p, \mathbf{k}_{F,2}^{p_2}, \mathbf{k}_{F,3}^{-p_3}, \mathbf{k}_{F,4}^{-p_4}) \right. \\ &\quad \left. - g_3(\mathbf{k}_F^p, \mathbf{k}_{F,2}^{p_2}, \mathbf{k}_{F,4}^{-p_4}, \mathbf{k}_{F,3}^{-p_3}) \right|^2 \frac{2\pi}{\hbar} \delta_{\mathbf{k}+\mathbf{k}_2, \mathbf{k}_3+\mathbf{k}_4+p\mathbf{G}} \\ &\quad \times \delta(\epsilon_k^p + \epsilon_{k_2}^{p_2} - \epsilon_{k_3}^{p_3} - \epsilon_{k_4}^{p_4}) \\ &\quad \times \frac{f_{k_2}^0 [1 - f_{k_3}^0] [1 - f_{k_4}^0]}{[1 - f_k^0]} \\ &\quad \times (\delta_{k,k'} + \delta_{k_2,k'} - \delta_{k_3,k'} - \delta_{k_4,k'}) \\ &\quad + \frac{\pi \hbar v_F}{LN_P} \frac{2\pi}{\hbar} g_{\text{imp}}^2 \delta(\epsilon_k^p - \epsilon_{k'}^{p'}) (1 - \delta_{k,k'}), \end{aligned} \quad (4)$$

where g_{imp}^2 is the square of the impurity scattering matrix element (normalized by $\pi \hbar v_F$) times the impurity concentration and N_P is the number of transverse momentum wave

vectors. In the framework of the Q1D electron gas model, the electron-electron umklapp scattering amplitude is given by $g_3(\mathbf{k}_F^p, \mathbf{k}_{F,2}^{p_2}; \mathbf{k}_{F,3}^{-p_3}, \mathbf{k}_{F,4}^{-p_4})$ (normalized by $\pi\hbar v_F$). For an array of quarter-filled but weakly dimerized chains, this momentum dissipative process has its origin in the scattering of two carriers from one side of the Fermi surface to the other, which is made possible if momentum conservation in (4) involves the longitudinal reciprocal lattice vector $\mathbf{G} = (4k_F, 0)$ at half filling ($k_F = \pi/2a$) [66]. Umklapp scattering amplitude is evaluated on the Fermi surface sheets $\mathbf{k}_F^{p=\pm} = (k_F^p(k_\perp), k_\perp)$, which is determined by the equation $\varepsilon_{\mathbf{k}_F}^p = 0$ and parametrized by the transverse wave vector k_\perp . As we will see below, its momentum-dependent renormalization at temperature T will be obtained by the RG technique at the one-loop level (see also the Appendix).

The longitudinal electric current is given in leading order by the expression

$$j_a \simeq \frac{2e}{LNpd_\perp} \sum_{\mathbf{k}} v_F f_{\mathbf{k}}^0 (1 - f_{\mathbf{k}}^0) \phi_{k_\perp}, \quad (5)$$

where $\phi_{k_F} \equiv \phi_{k_\perp}$ can be taken on the Fermi surface at low temperature. The conductivity σ_a or the inverse of the resistivity along the chains can then be written in the form

$$\sigma_a = \rho_a^{-1} = \frac{e^2}{c\hbar} \langle \bar{\phi}_{k_\perp} \rangle_{\text{FS}}, \quad (6)$$

where we have rescaled the deviation $\bar{\phi}_{k_\perp} = \phi_{k_\perp}/(\beta e \mathcal{E} d_\perp)$ and inserted the lattice constant c in the c^* direction. Here $\langle \dots \rangle_{\text{FS}}$ corresponds to an average over the Fermi surface. Therefore the numerical solution of (3) for $\bar{\phi}_{k_\perp}$ using the momentum and temperature dependence of g_3 provided by the RG method leads to the resistivity as a function of the temperature [27].

We can complete the standard description of interactions in the Q1D electron gas model entering in the RG calculations of the Appendix. Thus besides the spectrum (2) and the aforementioned umklapp term, we have the backward and forward scattering amplitudes $g_1(\mathbf{k}_{F,1}^-, \mathbf{k}_{F,2}^+; \mathbf{k}_{F,3}^-, \mathbf{k}_{F,4}^+)$ and $g_2(\mathbf{k}_{F,1}^+, \mathbf{k}_{F,2}^-; \mathbf{k}_{F,3}^-, \mathbf{k}_{F,4}^+)$, also defined on the Fermi surface. For a superconducting compound like (TMTSF)₂ClO₄, which at ambient pressure is close to a magnetic QCP, typical values of the bare parameters of the model have been assessed in previous works [19,27,67]. For example the Fermi temperature $T_F = \hbar v_F k_F / k_B \simeq 3000$ K and interchain hopping $t_\perp / k_B \simeq 200$ K are representative band parameters for the Bechgaard salts [68,69]. The bare couplings (normalized by $\hbar\pi v_F$) we use are $g_1 = g_2/2 \simeq 0.32$ and $g_3 \simeq 0.033$, consistent with uniform susceptibility measurements and weak dimerization of the organic stacks. In these conditions the RG calculations for relatively weak antineesting parameter t'_\perp lead to SDW ordering temperatures that are compatible with those found in the antiferromagnetic Bechgaard salts at ambient pressure.

To suppress the SDW state and bring the system close to the QCP, t'_\perp / k_B is raised to 42 K or so, yielding a critical temperature $T_c^0 \sim 1$ K for (*d*-wave) superconductivity [27], which is congruent to the situation that prevails in (TMTSF)₂ClO₄ in slow cooling. This set of figures will fix the initial conditions with zero pair-breaking effects ($\tau_0^{-1} = 0$) in the RG calculations [Eqs. (A1)] and the numerical solution

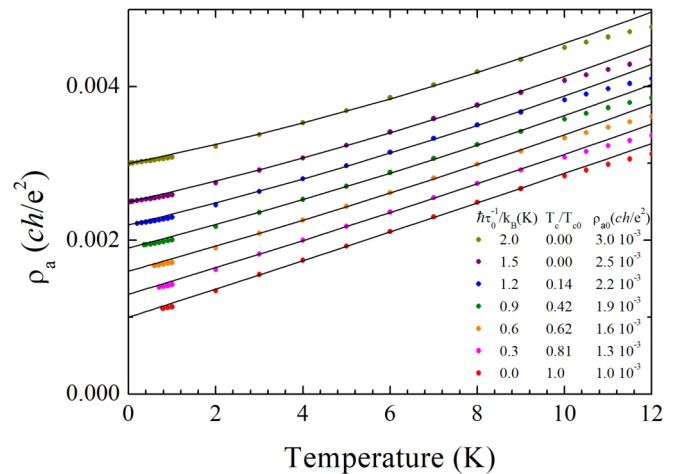


FIG. 3. Calculated temperature-dependent resistivity as a function of pair-breaking parameter $\hbar\tau_0^{-1}/k_B$. The continuous lines are fits to the polynomial expression $\rho_a(T) = \rho_{a0} + AT + BT^2$.

of the Boltzmann equation (3). Finally a normalized impurity scattering matrix element $g_{\text{imp}}^2 = 0.001$ has been used in order to simulate a small residual resistivity term from (4); its value smoothly increases with τ_0^{-1} . It is worth noticing that the interplay between the impurity and inelastic parts in (4) has been found to be negligible. In Fig. 3 we show the resistivity calculated from Eqs. (6) and (A1) in the low-temperature range down to the onset of critical superconducting domain occurring close to T_c where the renormalization flow goes to strong coupling.

Thus for zero pair breaking ($\tau_0^{-1} = 0$), the system is very close to the QCP and the resistivity becomes essentially T linear up to 10 K or so where a crossover to a sublinear temperature dependence takes place (see Ref. [70]). The origin of T linearity stems from the anomalous temperature dependence of umklapp scattering along the Fermi surface. Actually, despite the absence of SDW long-range order, spin correlations remain important and keep growing as temperature is lowered [28]. This derives from Cooper pairing terms in the RG flows [Eq. (A1)] which interfere constructively with the electron-hole (Peierls) pairing ones and then SDW fluctuations. This quantum fluctuation effect results in an anomalous temperature dependence of umklapp scattering that transforms the Fermi liquid quadratic T^2 dependence of resistivity into a T -linear behavior. Roughly speaking, the resistivity goes like $\rho_a \sim T^2 \langle g_3^2 \rangle_{\text{FS}}$ at low temperature where the mean-square value of umklapp scattering on the Fermi surface, which is representative of spin fluctuations, goes as $1/T$ at the critical t'_\perp [27].

When pair-breaking τ_0^{-1} departs from zero, T_c gradually decreases and follows relatively closely the Abrikosov-Gor'kov mean-field result down to the critical $\hbar\tau_0^{-1}/k_B \simeq 1.25$ K where $T_c \rightarrow 0$ (see Fig. 2). Concomitant with the fall of T_c , the temperature dependence of resistivity shows comparatively slow, albeit noticeable, alterations. We have proceeded to the fit of calculated resistivity curves of Fig. 3 to the polynomial form, $\rho_a = \rho_0 + AT + BT^2$, in the interval from 10 K down to the lowest temperature. According to Fig. 4, the linear coefficient A of resistivity weakens with

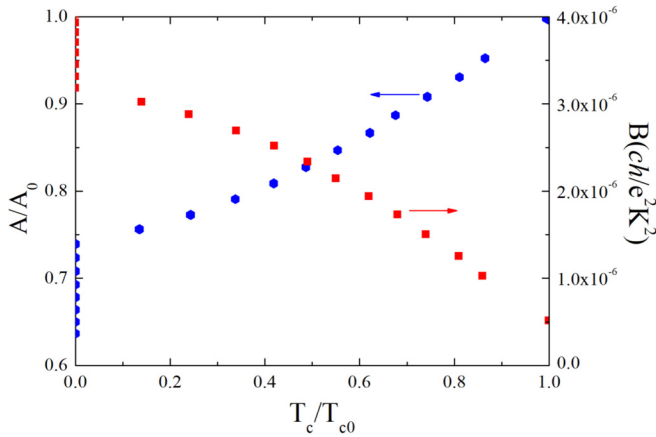


FIG. 4. A (left scale) and B (right scale) coefficients of the polynomial fit $\rho_a(T) = \rho_{a0} + AT + BT^2$ of the calculated resistivity (Fig. 3) as a function of T_c .

the size of τ_0^{-1} or the reduction of T_c . When $T_c \rightarrow 0$ as we approach the threshold τ_0^{*-1} , the reduction in the ratio A/A_0 reaches about 25%; it continues a steady decrease beyond that point. As for the Fermi liquid BT^2 term of resistivity, it becomes visible and weakly grows, but regularly for $T_c/T_c^0 < 1$.

It follows that when pair breaking is present the scaling of A with T_c no longer holds. It must be stressed, however, that the threshold energy scale for pair-breaking $\hbar\tau_0^{*-1}$, which is of the order of $k_B T_c^0$, only affects electronic states within a narrow energy shell around the Fermi surface. Although the corresponding decline of the Cooper loop (\mathcal{L}_C) entering the RG flow [Eq. (A1)] is sufficient to suppress the instability against d -wave superconductivity, scattering processes taking place at higher energy distance from the Fermi surface remain virtually unaffected. These still contribute to the mixing of Cooper pairing and SDW correlations. The latter alongside umklapp scattering continue to be strong and are responsible for sustaining the amplitude of the T -linear coefficient of resistivity despite the absence of a finite T_c .

Before closing this section, it is worth noticing that the suppression of T_c as x approaches the critical concentration x_c in Fig. 2 gives in principle rise to a quantum critical point. According to Figs. 3 and 4, however, this point does not give rise to any irregular change in the variation of linear coefficient for resistivity. This comes from the fact that the electronic degrees of freedom involved in the suppression of T_c are confined to a small energy shell, that is of the order of T_c itself, near the Fermi surface. These have essentially no additional influence whatsoever on the strength of spin fluctuations responsible for anomalous resistivity.

IV. DISCUSSION

Before we proceed with a comparison between theory and experiments in these alloys, it is important to notice that the present study is not dealing with a random distribution of strongly localized defects. Instead, we are facing a situation where the carrier lifetime in ordered regions is governed by the size of extended disordered domains. The transport data

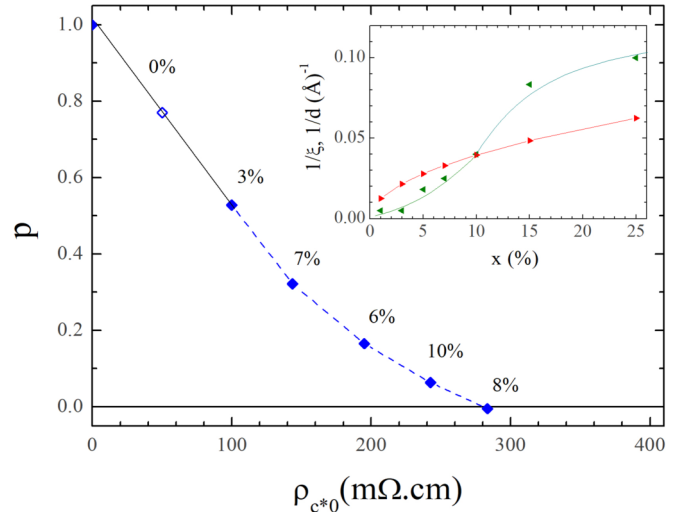


FIG. 5. The anion-ordered volume fraction p derived from the residual resistivity and the effective medium model. The inset shows the x dependence of the inverse correlation length for the $(0, 1/2, 0)$ ordering [48,49] and the inverse length between ReO_4^- sites located at random. We note that around $x = 10\%$ the $(0, 1/2, 0)$ anion-ordered coherence length becomes shorter than the average distance between centers. It is above such a concentration that a granular behavior prevails. The superconducting coherence spreads over the whole volume in the 6% sample with $p = 0.18$ whereas the percolation threshold for ordered regions should arise around $p = 0.3$, i.e., close to the 7% sample. This is a consequence of the proximity effect for superconductivity.

showing the remanence of nonsuperconducting domains at high ReO_4^- concentration suggest we are facing a problem of granular materials. Hence, the value of the residual resistivity cannot be simply extracted from the knowledge of the ReO_4^- concentration. The situation in the solid solution has similarities with that encountered in pure $(\text{TMTSF})_2\text{ClO}_4$ when the disorder is caused by rapid cooling of the sample [53]. Therefore, we use a similar procedure to extract the fraction of the respective ordered and disordered regions from the resistivity data according to the effective-medium theory for a two-component conductor [71,72],

$$1/\rho_{\text{eff}} = (p(1/\rho_{\text{Min}})^{1/3} + (1-p)(1/\rho_{\text{Max}})^{1/3})^3, \quad (7)$$

where p is the volume fraction of the anion-ordered domains. We take $\rho_{\text{Min}} = 50 \text{ m}\Omega \text{ cm}$ according to the data in Fig. 1 and ρ_{Max} is given by $\rho_{\text{Max}} = \rho_{\text{Min}} + \Delta\rho_{c^*} = 280 \text{ m}\Omega \text{ cm}$, where $\Delta\rho_{c^*}$ is the drop in residual resistivity coming from the anion ordering at 24 K which is known from the pure $(\text{TMTSF})_2\text{ClO}_4$ data ($\Delta\rho_{c^*} = 230 \text{ m}\Omega \text{ cm}$) [53]. The derived values for p are displayed in Fig. 5. Using the value for the residual resistivity of sample 0% in Eq. (7) would provide $p = 1$ for that sample in which the residual defects are not under control. As x-ray data have shown that full anion ordering does not exist even in pure and the best relaxed samples [73,74], we feel entitled to proceed differently for the determination of p in sample 0%. The derivation of p from Eq. (7) is expected to become meaningful only in samples where the additional disorder provided by alloying is large compared to the residual disorder in a pristine sample.

Therefore we have proceeded for p of the 0% sample with a linear interpolation between $\rho_{c=0} = 0$ and $100 \text{ m}\Omega \text{ cm}$ ($x = 3\%$) in Fig. 5 leading to the determination of p ($\simeq 0.78$) in the sample 0% of residual resistivity $\rho_{c=0} = 50 \text{ m}\Omega \text{ cm}$.

From Fig. 1, the 10% sample exhibits a superconducting transition ending at a finite value for the resistivity at zero temperature, which means that ordered regions are not percolant. Furthermore, in the 6% corresponding to $p = 0.18$ according to Fig. 5, the transition becomes complete, meaning that pairing coherence becomes infinite. The volume fraction for ordered domains in 6% is smaller than the actual value expected for the percolation threshold, namely, $p = 0.3$ at which the first cluster of 3D ordered anions with infinite length develops [75]. Such a situation can be understood taking into consideration the proximity effect which enables superconducting coherence over the entire volume even though ordered regions are still separated by small gaps of anion-disordered matter.

Regarding the SDW transition in Fig. 2, the onset of such a ground state in fast cooled samples has first been reported by Tomic *et al.* [52]. This phenomenon has been ascribed to the presence of Fermi surface folding whenever ClO_4^- anions do not order. The behavior of fast cooled samples has been revisited recently leading in turn to the highest SDW transition temperature at 6.6 K in a clean although fully anion-disordered sample [58]. Such a SDW transition temperature is actually lower than the transition occurring around 12 K in $(\text{TMTSF})_2\text{PF}_6$ or even 15 K in salts with tetrahedral anions such as $(\text{TMTSF})_2\text{ReO}_4$ when these latter anions remain ordered at low temperature [76]. As far as the clean but fully disordered $(\text{TMTSF})_2\text{ClO}_4$ is concerned, the decrease of the mean free path is likely responsible for the lowering of the transition temperature. In the quenched-6% sample more scattering coming from the disordered ClO_4^- anions adds to the original scattering in the relaxed 6% ReO_4^- substituted sample leading in turn to a further lowering of the transition temperature. As emphasized in Sec. III, nonmagnetic defects provide alongside antineesting some additional pair-breaking reducing density-wave correlations and then T_{SDW} , in the same way as nonmagnetic defects act in d -wave superconductors [77–79].

We can now plot the dependence of A on T_c , as shown in Fig. 6 comparing experimental data provided by the analysis of the resistivity of the metallic phase above T_c to the Boltzmann-RG calculation. We must be cautious because A is related to a homogenous medium in the two extreme cases of high and low T_c only where the sample is anion ordered or disordered, respectively, but in between (say, 7%) we face a more complex situation where A is related to the mixture of two metallic phases whereas T_c is given by the onset of superconductivity in the ordered phase only.

Experiments and theory of Fig. 6 agree on the fact that the suppression of A is no more than 25% once superconductivity in alloys is fully suppressed. In these conditions a sustained strength of a T -linear component for resistivity indicates that in the presence of nonmagnetic pair breaking the system moves away only slowly from the magnetic QCP. The amplitude of SDW correlations and with it the enhancement of umklapp scattering persist to a great extent in the anomalous temperature dependence of resistivity. The comparisons between theory and experiments for the growth of the Fermi

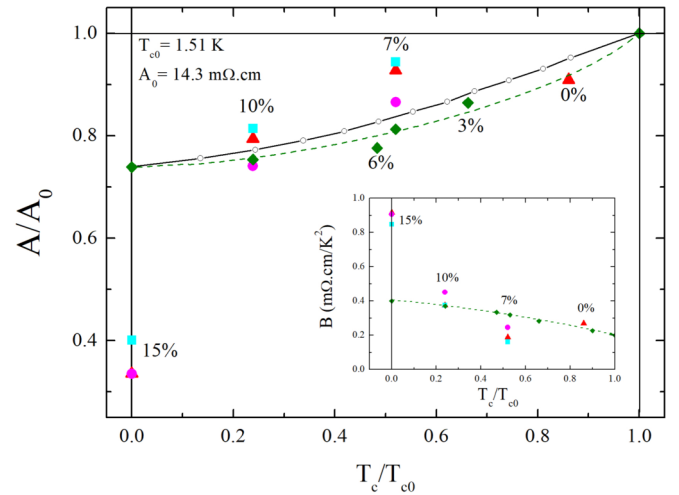


FIG. 6. Dependence of A/A_0 on T_c , from experiment and theory (solid line). The colored square, triangle, and circle symbols correspond to experimental data according to the different fitting ranges described in Fig. 2. The green diamonds are estimated A/A_0 derived from the two-conductor model in Eq. (7) for the inelastic contribution. $A/A_0 = 0.75$ at $T_c = 0 \text{ K}$ and the volume fractions p according to Fig. 5 have been used. The dashed green line is a guide for the eye. The inset displays the dependence of the unnormalized B on T_c . The green diamonds are provided by the two-conductor model and the dashed green line is a guide for the eye.

liquid B coefficient of resistivity are displayed in Fig. 4 for the theory and in the inset of Fig. 6 for the experiment. The rise of B as T_c decreases is steady and compatible with calculations; its amplitude remains small.

Similarly, in the case of pure $(\text{TMTSF})_2\text{ClO}_4$, the evolution of the T -linear term A as a function of the cooling rate can be extracted from a polynomial fit of the raw data of previous works in Ref. [53], which we reproduce in Fig. 7 for representative cooling rates. The fit is in the window between 4.5 and 10 K, namely, outside the paraconductive downturn of resistivity whose amplitude increases with the cooling rate. The suppression of A is at most 25% when the volume fraction of anion-ordered regions is monitored by the cooling rate leading in turn to a complete suppression of T_c , as displayed in Fig. 7.

Finally, we comment on the behavior of $\rho_c(T)$ in pristine $(\text{TMTSF})_2\text{ClO}_4$ at various cooling rates in Fig. 7. These data come from the extension to finite temperature of a previous work investigating residual resistivity and superconductivity in $(\text{TMTSF})_2\text{ClO}_4$ as the cooling rate is varied following a procedure already presented in Ref. [53]. The resistivity has been normalized to the value of $28 \Omega \text{ cm}$ at ambient temperature following reports of previous $\rho_c(T)$ measurements [37,53]. What is clear at first sight in Fig. 7 is the marked contrast between the strong effect of cooling rate on the residual resistivity and the much weaker effect on the inelastic part, a feature captured by the calculations of Sec. III.

The pure $(\text{TMTSF})_2\text{ClO}_4$ study [53] has shown that the material acquires a granular texture above a critical cooling rate of 1 K/min where disordered regions separate anion-ordered domains in which bulk superconductivity arises and possibly spread over the whole sample via proximity effect. The volume fraction p related to ordered domains has been

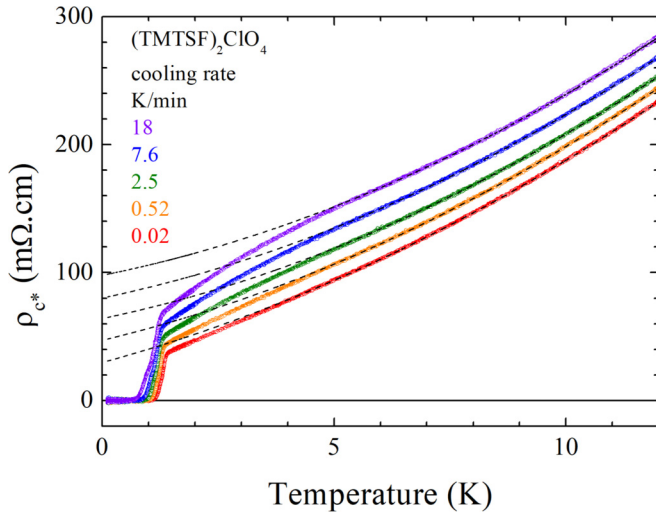


FIG. 7. $\rho_{c^*}(T)$ measured on $(\text{TMTSF})_2\text{ClO}_4$ and displayed in the temperature range between 12 K and the superconducting region for 0.020, 0.52, 2.5, 7.6, and 18 K/min cooling rates across the anion ordering transition. The same polynomial fitting procedure as for the data in Fig. 1 has been used to extract the A , B , and ρ_{c^*0} parameters at each cooling rate within the fitting window of 5–10 K. The results are plotted as broken lines.

estimated at each cooling rate using the model of a two-component conductor [53]. The data for the fitted values of A are displayed on Fig. 8 versus the anion-disordered volume fraction $1 - p$. A has been normalized to its value at the smallest cooling rate. We also show in the same figure the prediction of the effective medium model using Mathiessen's law applied to the inelastic contribution, taking $A/A_0 = 1$ and 0.75 at $1 - p = 0$ and 1, respectively, following the results from the alloys series. The solid squares with the dotted line are theoretical Boltzmann-RG points corrected for the evolution of the mean free path within each puddle at concentration p . These values provide an improvement for the agreement

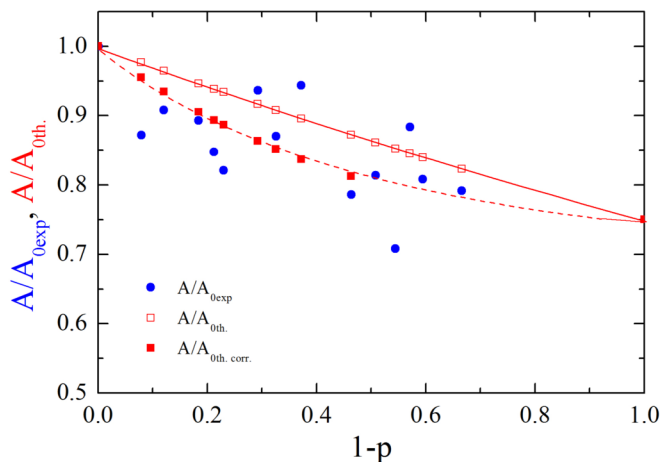


FIG. 8. Dependence of A/A_0 on the disordered volume fraction $1 - p$ in pure $(\text{TMTSF})_2\text{ClO}_4$ at various cooling rates (blue dots). The model calculations for an effective conductor with $A/A_0 = 0.75$ at $p = 0$ are shown as red open squares. The model calculations taking into account the evolution of the mean free path in anion-ordered regions are the red solid squares.

with the experimental data. The variation of the T -linear resistivity of the metallic phase above T_c at various cooling rates appears to be controlled nearly equally by lifetime and phase mixture effects.

V. CONCLUSION

In this work, we have examined the influence of non-magnetic disorder generated by finite domains of anion ordering on the temperature dependence of resistivity in both $(\text{TMTSF})_2(\text{ClO}_4)_{1-x}(\text{ReO}_4)_x$ alloys and pristine $(\text{TMTSF})_2\text{ClO}_4$ at different cooling rates. These domains, presumably of different electronic structure, were found to have a strong impact on the onset of superconducting order. T_c is suppressed pointing to the existence of pair-breaking effects congruent with unconventional Cooper pairing. At variance with T_c , the pronounced T -linear resistivity in $(\text{TMTSF})_2\text{ClO}_4$ observed in slow cooling conditions does not reveal as a strong decline in alloying. A polynomial decomposition of resistivity temperature dependence has revealed that the linear term is not strongly weakened with x , indicating that quantum critical features remain strong despite the suppression of superconducting long-range order. The scaling between the T -linear term strength and T_c , which is seen under pressure in pure samples, was found to no longer hold as a function of disorder tuned by x or to a certain extent the cooling rate.

These results have been compared to a theoretical calculation based on the combination of the linearized Boltzmann equation of transport and the renormalization group approach to the umklapp vertex function. In the framework of the Q1D electron gas model, we have shown that a range of relatively small energy scales for pair breaking is sufficient to instill a characteristic fall in the metallic instability against d -wave superconductivity. The calculated drop in T_c agrees with the one observed with x . Although small pair-breaking energy is enough to suppress a scale like T_c , its impact on the strength of SDW fluctuations that feeds the amplitude of interparticle umklapp scattering remains relatively weak. The anomalous growth of umklapp at low temperature was found to be weakly altered along with the linear resistivity term whose amplitude remains sizable, in agreement with experiment. Although the correlation between T_c and linear resistivity apparently breaks down, the correlation of the latter with spin fluctuations, which are the core of d -wave pairing, prevails.

ACKNOWLEDGMENTS

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APPENDIX: RENORMALIZATION GROUP EQUATIONS FOR THE ELECTRON GAS WITH PAIR-BREAKING EFFECT

In this Appendix we give the one-loop RG flow equations for the momentum scattering amplitudes $g_{i=1,2,3}$ in the presence of pair-breaking effects due to a finite lifetime τ_0 of the carriers. The RG approach to the quasi-1D electron gas is detailed in Refs. [19,28,80]. Each energy shell of thickness $\frac{1}{2}E_0(\ell)d\ell$, located at $\pm\frac{1}{2}E_0(\ell)$ from either sides of the Fermi sheets, is segmented into N_p patches, each centered at a particular value of the transverse momentum k_\perp , where $E_0(\ell) = E_0 e^{-\ell}$ is the scaled bandwidth at step ℓ and $E_0 \equiv 2E_F$ is the initial bandwidth. The successive partial trace integration of electron degrees of freedom in the partition function leads to the renormalization or flow of the scattering amplitudes g_i as a function of ℓ . At the one-loop level, the flow combines corrections from the electron-hole (Peierls) and electron-electron (Cooper) interfering channels of scattering. Impurity (back) scattering introduces pair breaking in the form of a carrier lifetime τ_0 in both $2k_F$ electron-hole (Peierls) [81] and (*d*-wave) Cooper [82] channels. This leads to the flow equations for the $g_i(\mathbf{k}_{F1}^{p1}, \mathbf{k}_{F2}^{p2}; \mathbf{k}_{F3}^{p3}, \mathbf{k}_{F4}^{p4}) \rightarrow g_i(k_{\perp 1}, k_{\perp 2}, k_{\perp 3}, k_{\perp 4})$ on the Fermi surface, which can be written in the compact form

$$\begin{aligned} & \partial_\ell g_1(k_{\perp 1}, k_{\perp 2}, k_{\perp 3}, k_{\perp 4}) \\ &= (-2g_1 \circ g_1 + g_1 \circ g_2 + g_2 \circ g_1) \partial_\ell \mathcal{L}_P \\ & \quad - (g_1 \circ g_2 + g_2 \circ g_1) \partial_\ell \mathcal{L}_C, \\ & \partial_\ell g_2(k_{\perp 1}, k_{\perp 2}, k_{\perp 3}, k_{\perp 4}) \\ &= -(g_1 \circ g_1 + g_2 \circ g_2) \partial_\ell \mathcal{L}_C \\ & \quad + (g_2 \circ g_2 + g_3 \circ g_3) \partial_\ell \mathcal{L}_P, \\ & \partial_\ell g_3(k_{\perp 1}, k_{\perp 2}, k_{\perp 3}, k_{\perp 4}) \\ &= (-g_1 \circ g_3 - g_3 \circ g_1 + g_2 \circ g_3 + g_3 \circ g_2) \partial_\ell \mathcal{L}_P \\ & \quad + 2g_2 \bullet g_3 \partial_\ell \mathcal{L}_P, \end{aligned} \quad (\text{A1})$$

where $\partial_\ell = \partial/\partial\ell$. $\mathcal{L}_{v=P,C}$ are the Peierls and Cooper loops whose derivative at finite temperature comprises an integration over the patch. These take the form

$$\begin{aligned} \partial_\ell \mathcal{L}_v(k_\perp, q_{\perp v}^{(v)}) &= \frac{E_0(\ell)}{4} \text{Re} \left\{ \sum_{\mu=\pm 1} \int_{k_\perp - \frac{\pi}{N_p}}^{k_\perp + \frac{\pi}{N_p}} \frac{dk_\perp}{2\pi} \right. \\ & \quad \times \frac{\theta(|E_0(\ell)/2 + \mu A_v| - E_0(\ell)/2)}{E_0(\ell) + \mu A_v + i\hbar\tau_0^{-1}} \end{aligned}$$

$$\begin{aligned} & \times [\tanh[\beta E_0(\ell)/4] \\ & \quad + \tanh[\beta(E_0(\ell)/4 + \mu A_v/2)]] \Big\}, \end{aligned} \quad (\text{A2})$$

where

$$\begin{aligned} A_v(k_\perp, q_{\perp v}^{(v)}) &= -\epsilon_\perp(k_\perp) - \eta_v \epsilon_\perp(\eta_v k_\perp + q_{\perp v}^{(v)}) \\ & \quad + \eta_v \epsilon_\perp(\eta_v k_{\perp 2(4)} + q_{\perp v}^{(v)}) + \epsilon_\perp(k_{\perp 2(4)}) \end{aligned} \quad (\text{A3})$$

with

$$\epsilon_\perp(k_\perp) = -2t_\perp \cos k_\perp d_\perp - 2t'_\perp \cos 2k_\perp d_\perp$$

and $q_{\perp P}^{(v)} = k_{\perp 3} - k_{\perp 2} = k_{\perp 1} - k_{\perp 4}$ ($k_{\perp 3} - k_{\perp 1} = k_{\perp 2} - k_{\perp 4}$), $q_{\perp C} = k_{\perp 1} + k_{\perp 2} = k_{\perp 3} + k_{\perp 4}$; $\eta_P = 1$ and $\eta_C = -1$. $\theta(x)$ is the Heaviside function [$\theta(0) \equiv \frac{1}{2}$].

The momentum dependence of couplings in the discrete convolution products “ \circ ” over the internal k_\perp loop variable on the right-hand side of Eqs. (A1) are in order $g(k_\perp, k_{\perp 4}, k_{\perp 1}, k_\perp - q_{\perp P})g(k_\perp, k_{\perp 2}, k_{\perp 3}, k_\perp - q_{\perp P})$ for the Peierls channel, $g(k_{\perp 1}, k_{\perp 2}, k_\perp, q_{\perp C} - k_\perp)g(k_{\perp 3}, k_{\perp 4}, k_\perp, q_{\perp C} - k_\perp)$ for the Cooper channel, and $g(k_\perp, k_{\perp 4}, k_{\perp 2}, k_\perp - q'_{\perp P})g(k_{\perp 1}, k_\perp, k_{\perp 3}, k_{\perp 1} - q'_{\perp P})$ for the “ \bullet ” product of the off-diagonal Peierls channel.

The integration of Eqs. (A1) up to $\ell \rightarrow \infty$ gives the values of the moment-dependent scattering amplitudes g_i at temperature T . A singularity in the scattering amplitudes signals an instability of the electron gas against the formation of an ordering state at a critical temperature. For the band parameters and repulsive interactions fixed in Sec. III, the most probable instabilities are against either SDW or *d*-wave SC ordering. At low antinesting t'_\perp and zero pair breaking $\tau_0^{-1} = 0$, an SDW instability occurs with a singularity in both g_2 and g_3 amplitudes at the modulation wave vector $\mathbf{q}_0 = (2k_F, q_{\perp P} = \pi)$ and $T_{\text{SDW}} \sim 10\text{--}20$ K. By tuning the antinesting at the QCP near $t'_\perp \simeq 42$ K, the SDW is suppressed and the instability occurs in the *d*-wave superconducting channel at $T_c^0 \sim 1$ K, which is associated to a singularity for both g_2 and g_1 couplings at zero Cooper pair momentum $\mathbf{q}_0 = (0, q_{\perp C} = 0)$, but with a cosine modulation in k_\perp space. T_c^0 stands as the maximum T_c at $\tau_0^{-1} = 0$ which we liken to the situation of the pure and slowly cooled (TMTSF)₂ClO₄ compound.

By increasing τ_0^{-1} it is the weakening of the Cooper loop (\mathcal{L}_C) singularity in Eq. (A2) that affects the instability against *d*-wave superconductivity. T_c/T_c^0 goes down and follows relatively closely the mean-field Abrikosov-Gor'kov result [82] as shown in Fig. 2. The concomitant impact of τ_0^{-1} on the momentum and temperature profile of umklapp scattering and in turn on resistivity as obtained from the Boltzmann equation (4) is discussed in Secs. III and IV.

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