Tunnelling in Organic Superconductors

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We discuss the possibility of deciding on the symmetry of the superconducting phase in the organic superconductors (Bechgaard salts), using tunnelling experiments. We first briefly review the properties of organic superconductors, and the possibility to have unconventional (triplet) superconductivity in these systems. We then present a simple scheme for computing the full current-voltage characteristics for tunnelling experiments within the framework of the non-equilibrium Keldysh Green function formalism. This formalism is flexible enough to address different pairing symmetries combined with magnetic fields and finite temperatures at arbitrary bias voltages. We then discuss extensively how to apply these results to probe for the symmetry of the superconducting order parameter.

§1. Introduction

Organic materials have provided physicists with an extraordinary laboratory to study the effects of interactions in solids. These materials indeed show remarkable effects due to the interplay of interactions and dimensionality.\textsuperscript{1} Indeed such compounds made of coupled chains allow to realize practically the physics of Luttinger liquids,\textsuperscript{2} one of the very few controlled examples of non-Fermi liquids\textsuperscript{3} due to interactions. But contrarily to other realizations of one-dimensional systems such as nanotubes,\textsuperscript{4,5} quantum wires\textsuperscript{7,8} or edge states in the quantum Hall effect,\textsuperscript{10,11} the organics offer unique challenges. Indeed, because of their very three-dimensional nature, they provide not a single one-dimensional electron gas, but a very large number of such one-dimensional systems coupled together. This allows thus for a unique new physics to emerge where the system is able to cross over from a one-dimensional behavior to a more conventional three-dimensional one.

As a consequence of this dimensional crossover, at low temperature these materials undergo instabilities towards three-dimensionally ordered states, such as spin-Peierls, antiferromagnetic and even superconducting states. Needless to say, the presence of superconductivity in these compounds is a tantalizing and challenging question. Despite a period of quarter of a century since the discovery of superconductivity in these materials,\textsuperscript{13} the mechanism and even the symmetry of this superconducting phase have remained elusive, and many efforts have been devoted to this subject. Recently, the case for triplet nature of this superconducting phase has been made,\textsuperscript{14} mostly by measurements of the upper critical field and by NMR measurements, but the subject is far from being settled. If true, this would put the organics in the relatively small club of condensed matter materials where interactions can lead to unusual triplet superconductivity,\textsuperscript{15} other examples being $^3$He (a superfluid) and Sr$_2$RuO$_4$,\textsuperscript{16,17} and can certainly shed light on the pairing mechanism.

In this note we explain how point contact tunnelling experiments can be used to probe for the symmetry of a superconducting state. In the recent years, the
possibilities to perform point contact tunnelling have been drastically enhanced with the development of scanning tunnelling microscopy (STM). Correspondingly, theories to interpret tunnelling experiments in superconductors have evolved from simple semiconducting band models to more systematic approaches based on the tunnelling Hamiltonian. We show here how one can extend and simplify the formalism to make it more versatile and easy to implement. This allows to study the case of superconducting (singlet or triplet) leads, as well as the effects of magnetic fields on the junction. The resulting theory can be thus directly used as a probe of the symmetry of the leads. We thus discuss how such a probe can be used for the case of the organics.

The plan of this note is as follows. In §2 we briefly recall the salient points of the physics of organic compounds. In §3 we give an outline of our formalism for the tunnelling between triplet superconductors. Section 4 discusses how to test for triplet superconductivity in the organics. Finally some conclusions and perspectives can be found in §5.

§2. A brief word on organics

Let us briefly recall some of the properties of the organic materials. We will focus here on the Bechgaard salts TMTSF$_2$X which were the first organic compounds to exhibit superconductivity, and have thus been the focus of intense experimental and theoretical studies. A general review on these, and parent compounds (so-called Fabre salts TMTTF$_2$X), can be found in Ref. 1.

In addition to the superconducting phase itself, these materials have a notoriously rich phase diagram (cf. Fig. 1) and exhibit a host of remarkable properties (non-FL metallic behavior, quantized Hall conductance, Fröhlich conductivity), many of which are still poorly understood. These materials exhibit a quasi-one-dimensional structure due to the nature of the molecular orbitals protruding from the flat molecule TMTSF (tetramethyltetraselenafulvalene) which is the basic building block of the Bechgaard salts: the overlap of the $\pi$-orbitals of the selenium or sulfur atoms leads to a high mobility of electrons along the stacking direction. The hopping integrals in the perpendicular directions are smaller by more than one order of magnitude. Estimated values of the hopping integrals along the stack direction ($a$-axis) and the two perpendicular axes pointing towards neighboring stacks ($b$-axis) and towards the anions ($c$-axis) respectively are: $t_a : t_b : t_c = 1000$ K : $100$ K : $30$ K. Therefore one can think of these materials as one-dimensional chains coupled by small inter-chain hoppings. Given the hierarchy of transverse coupling the system is first expected to become two dimensional and then three dimensional at low temperatures. The normal phase of these materials presents unusual transport and optical conductivity properties that have been interpreted as the signature of Luttinger liquid physics. The high temperature phase of these materials thus clearly shows effects of strong correlations. As the temperature is lowered, the material recovers features that are more and more reminiscent of a normal Fermi-liquid material, and finally undergoes instabilities towards various ordered phases (spin-Peierls (SP), antiferromagnetic (AF), spin-density wave (SDW) and supercon-
Fig. 1. Unified experimental phase diagram for the TM compounds (from Ref. 27). Either pressure or chemical changes (increasing pressure corresponds to either going from the TMTTF to the TMTSF family or changing the anions) yields the same phases [MI: Mott insulator, LL: Luttinger liquid metal, FL: Fermi liquid metal, SP: spin-Peierls, AF: antiferromagnetic spin-density wave, SC: superconducting]. Note that the pressure axis does not change the doping of the material that remains a quarter filled system. The TMTTF family is insulating at ambient pressure whereas the TMTSF family shows good metallic behavior at room temperature. The superconducting phase, stemming out of the antiferromagnetic one is still very poorly understood.

The nature of the molecule (TMTTF vs TMTSF) or the ions, as well as the application of pressure allows to control the phase diagram, by changing the hopping integrals. This modifies the relative importance of the kinetic energy and Coulomb interaction. The chemical and pressure changes have similar effects, which can be summarized by the unified phase diagram of Fig. 1. More issues on the effect of interactions and the properties of the normal phases can be found in Ref. 34) and references therein.

Among the ordered phases, the superconducting one\textsuperscript{13}) is by far the most mysterious. The mechanism behind it is still unknown, and considerable debate takes place on that point (see e.g. Refs. 35)–37) and references therein). Since at the temperature when superconductivity occurs, the compound is not in the one-dimensional limit anymore, but well in the three-dimensional regime, one could in principle think of a conventional BCS mechanism. However the proximity of the AF phase as well as the fact, clear from the high temperature phase, that interactions are particularly strong in these compounds make such a simple pairing mechanism very unlikely. Besides the mechanism itself, the symmetry of its superconducting order parameter seems to be spin-triplet $p$-wave\textsuperscript{38}) but claims of $f$-wave orbital symmetry were also put forward;\textsuperscript{39),40}) the microscopic origin for the pairing remains largely a matter of debate. One of the most striking features is the existence of a common boundary between an antiferromagnetic phase (SDW) and the superconducting phase that is
stabilized as pressure is increased. Remarkably, the largest values of $T_c$ are actually reached near this boundary, at which the system is expected to display an enhanced $SO(4)$ symmetry.\textsuperscript{41} Intriguingly, recent experiments have identified a pressure window around the boundary at which the SDW and the superconducting regions seem to segregate.\textsuperscript{42} Among the properties of the superconducting phase that would be in favor of triplet superconductivity are the sensitivity to non-magnetic impurities,\textsuperscript{37} the absence of Knight shift signals at the superconducting transitions\textsuperscript{43} and, perhaps the most striking, the anomalously large values of the measured upper critical fields.\textsuperscript{44} Given the nature of these systems, it is difficult to perform phase sensitive experiments such as the ones done for the cuprates, so other probes are needed to unambiguously decide on the symmetry of the order parameter. We now examine tunnelling as such a probe.

§3. Tunnelling

The earliest successful theoretical models to study superconducting tunneling junctions were based on a scattering picture and semiconducting-like bands.\textsuperscript{20,21,45,46} It was later shown that those results can be recovered using a tunnelling Hamiltonian as the starting point.\textsuperscript{22–24} A large series of experiments on atomic-size contacts\textsuperscript{47–52} showed impressive agreement with the theory by considering a small number of independent conduction channels, each of them well described by a point-contact model. Tunnelling is thus a very efficient probe of the properties of the junctions, and one can expect to use it to determine the symmetry of the superconducting order parameter in the junction leads. However the techniques used so far are either semi-phenomenological or very heavy, and one thus need to find a formulation that is both simple enough to be extended to the case of unusual superconductors and to finite magnetic fields and finite temperature, and yet accurate enough to compute the full current-voltage curve. We describe here such a technique, based on the Keldysh formalism. More details can be found in Refs. 25 and 26).

Let us use a tunnelling Hamiltonian formalism to calculate the full current-voltage characteristics of different types of tunnel junctions where each side of the junction can be either a normal metal (N), a singlet (S) or a triplet (T) superconductor. We model the system with a (one-dimensional) Hamiltonian that includes the two leads and a tunnelling term: $H = H_1 + H_2 + H_{\text{tun}}$. Each lead is described by

$$K = \xi_{ck\sigma} \psi_{ck\sigma}^\dagger \psi_{ck\sigma} - \left\{ \Delta_a \left[ \psi_{Rk\beta}^\dagger \sigma^a_{\beta\alpha} \alpha \psi_{Lk\alpha}^\dagger \right] + \text{h.c.} \right\},$$

where $K = H - \mu N$ and $\mu$ is the corresponding electrochemical potential. All the indices are summed over, in particular $k$ is the lattice momentum, Greek indices correspond to the spin, and $c \in (L, R) \equiv (-1, +1)$ sums over the two possible chiralities. $\xi_{ck\sigma} = c v_F k - \mu - \sigma h$ are the corresponding linear dispersions, shifted by the inclusion of chemical potential and magnetic field along the $\hat{z}$-axis (for convenience we will take $v_F = 1$). This is the extension of the pairing-approximation Hamiltonian found in BCS theory to the triplet case (for more details see Refs. 25 and 26)).
third term in the Hamiltonian describes the tunnelling:

\[ H_{\text{tun}} = \sum_{\ell,\ell',\sigma} t_{\ell\ell'} \psi_{\ell\sigma}^\dagger(0) \psi_{\ell'\sigma}(0). \]  

(3.2)

According to this term an electron with spin \( \sigma \) can hop from lead \( \ell' \) into lead \( \ell \) with a tunnelling matrix amplitude given by

\[ t_{\ell\ell'} = \begin{pmatrix} 0 & t^* \\ t & 0 \end{pmatrix}. \]  

(3.3)

Since the number of particles in each lead is a conserved quantity in the absence of tunnelling (pair fluctuations are to be regarded as an artifact of the mean field approach), we can define the current as given by the rate of change in the relative particle number caused by tunnelling. One writes

\[ I = \frac{e}{2i} \langle [H_{\text{tun}}, N_1 - N_2] \rangle. \]  

(3.4)

Having in mind that very simple models of the leads suffice to achieve even quantitative agreement with the experiment when calculations are carried out to capture the main features of point-contact transport on conventional superconductors, with dimensionality playing little or no role, it is justified to use one-dimensional leads to carry out all the standard calculations. For unconventional superconductors the situation is more complex, because the anisotropic nature of the pair wave function has to be taken into account when modelling the leads. But the organic superconductors that we are interested in are supposed to have \( p \)-wave symmetry and be highly anisotropic. We can therefore, as a first approximation, conveniently restrict ourselves to a one-dimensional model and adopt a formalism that encompasses both \( s \)- and \( p \)-wave symmetries, as well as the normal state.

In order to deal with an out of equilibrium situation, we use the so-called Keldysh formalism.\(^{53}\) We treat the tunnelling term to all orders to calculate the full I-V line and give a quantitative account of the subgap structure. Notice that in this framework the current is given by

\[ I = \frac{e}{2i} \sum_{\sigma} \int \frac{d\omega}{2\pi} \left\langle \psi_{2,\sigma}^\dagger(0) \psi_{1,\sigma}(0) - \psi_{1,\sigma}^\dagger(0) \psi_{2,\sigma}(0) \right\rangle_{\text{kel}}, \]  

(3.5)

where ‘kel’ denotes the Keldysh component of the correlation function. Since the current depends only on the fields at \( x = 0 \) one can integrate the \( x \) dependence in the leads to obtain from \( H \) a local quadratic action for the contacts:

\[ S_{\text{jun}} = S_1 + S_2 + S_{\text{tun}}. \]

With \( S_{\text{tun}} \) obtained directly from \( H_{\text{tun}} \) and \( S_\ell = \int \frac{d\omega}{2\pi} \Psi_\ell^\dagger(\omega) \hat{g}^{-1} \Psi_\ell(\omega) \). Here \( \Psi_\ell \) is an 8 component spinor and \( \hat{g}^{-1} \) is a matrix whose inverse \( (\hat{g}) \) is given by the standard advanced, retarded and Keldysh components of the local Green functions of the lead. For instance, for the case where \( \Delta_1 = \Delta_2 = 0 \),

\[ g_{c\sigma,c\sigma}^{\text{[ret,adv]}} = \frac{-\left(\omega - \mu_\ell + c\sigma h \mp i0^+\right)}{2\sqrt{\left|\Delta_0 + c\sigma \Delta_3\right|^2 \left(\omega - \mu_\ell + c\sigma h \pm i0^+\right)^2}}. \]
Fig. 2. Representation of a set of frequencies involved in a high-order tunnelling process (the vertical axis corresponds to frequencies). The chemical potentials, thermal distributions and superconducting quasiparticle densities of states in each side of the junction are all indicated. The figure illustrates the lowest order contribution to tunnelling at low voltages for a case that corresponds to a process involving one electron plus three Cooper pairs.

\[ g_{\text{ret}, \text{adv}}^{[r, \ell]} = \frac{\left(\Delta_0 + c\sigma\Delta_3\right)^{[r]}_{\ell=L}}{2\sqrt{\left|\Delta_0 + c\sigma\Delta_3\right|^2 - (\omega - \mu_\ell + c\sigma\hbar \pm i0^+)^2}}. \]

And the Keldysh component is 

\[ g_{\text{kel}} = (g_{\text{ret}} - g_{\text{adv}})^{\text{tanh}} \left((\omega - \mu_\ell)/2T\right). \]

Attention must be paid to the fact that frequencies have different reference levels in each side of the junction when a bias is present. While, in each side of the junction, states of frequencies with equal positive and negative shifts from the Fermi level are related by the pairing fluctuations, across the junction, same frequency states are related by tunnelling (see Fig. 2). Thus, the full action for the junction is not diagonal. To each value \( \omega_0 \) in the frequency window (of size \( eV = \mu_1 - \mu_2 \)) defined by the chemical potentials in the two leads, one infinite set of related frequencies can be assigned \( (p > 0) \):

\[
\begin{align*}
\omega_p &= 2\mu_{2-p \mod 2} - \omega_{p-1}, \\
\omega_{-p} &= 2\mu_{1+p \mod 2} - \omega_{1-p}.
\end{align*}
\]

The action is block diagonal between these sets. Discretizing the frequencies in the window defined by the voltage difference defines in turn a discretization of the whole frequency axis. We shall deal with one set of frequencies at a time, and since the sets are infinite, we will truncate their hierarchy. This amounts to introducing a limit in the number of Andreev reflections. One can then numerically invert the action block by block (for more details see Refs. 25 and 26). The off-diagonal Green function matrix elements thus obtained allow to compute the current using Eq. (3.5).
§4. Test for triplet

As it was stated above, in the Bechgaard salts, both the orbital and the spin symmetries of the superconducting phase are not clearly known, not to speak of the microscopic mechanism responsible for the pairing. Since the uncertainties are such, we think the best way to proceed would be if the experiments try to resolve these issues one at a time and to our opinion the most approachable one, from the point of view of tunnelling, is the issue of the symmetry in spin space of the superconducting order parameter. Thus, the question to be answered is whether the electrons in a Cooper pair form a singlet or a triplet.

To discuss and compare the I-V characteristics for different types of junctions, we choose some convenient set of parameters that we will use in all the figures. For the tunnelling overlap integral we choose the values $t = 0.2$ and $t = 0.5$ (that correspond, in the notation of Ref. 20), to $\alpha \simeq 0.15$ or $Z = 2.4$ and to $\alpha = 0.64$ or $Z = 0.75$, respectively), and when there is a magnetic field we fix its value to $h = 0.2$ in units of $\Delta$ (by $\Delta$ we mean the magnitude of the singlet gap, $\Delta_0$, or of the triplet vector order parameter depending on the case). We show curves for the dc response in the limit of vanishing temperatures. The geometry of the junction we consider corresponds to tunnelling perpendicular to the chains of the quasi-one-dimensional compound, with an orbital order parameter aligned along them. In this situation no mid-gap states are expected and, therefore, no midgap features either. We concentrate on what happens at the conduction edge, in particular the effects of applied fields. We refer the reader to the literature for some recent studies that look at the effect of fields on zero bias anomalies.

Let us start with the case of normal-metal–superconductor junctions, that would correspond to standard STM experiments. We show in Fig. 3(a) typical curves for a point-contact junction between a normal metal and a conventional singlet-paring superconductor. The solid lines correspond to the N-S junction in zero field and the dashed line is for one of the junctions in the presence of a magnetic field that produces what would be seen as a Zeeman splitting of the differential conductance peak. The second part of the figure corresponds to a junction between a normal metal and an unconventional triplet-pairing superconductor. The solid lines correspond to the N-T junction in zero field and the dashed line is for the $t = 0.2$ junction when in the presence of a magnetic field that is aligned with the vector order parameter $\vec{\Delta}$. If one considers a magnetic field perpendicular to the order parameter ($\vec{h} \perp \vec{\Delta}$), it has no effect on the I-V characteristic.

In Fig. 4, we display typical curves for junctions in which both sides are conventional spin-singlet superconductors and junctions connecting a spin-singlet and spin-triplet superconductor. These curves address the situation of a potential STM experiment carried out using a superconducting tip made out of a conventional superconductor used to probe a superconducting phase of unknown symmetry. We remind the reader that, in the case of conventional superconductors (dotted lines) and when orbital effects can be ignored, the I-V is not sensitive to applied magnetic fields. On the other hand, the solid lines correspond to singlet-triplet junctions, that are insensitive to the orientation of the vector order parameter on the triple-pairing side.
Fig. 3. I-V characteristics of normal-superconductor junctions. (a) N-S junctions for $t = 0.2$ (lower curves: with and without applied magnetic field, dashed and solid line respectively, $h = 0.2$) and $t = 0.5$ (upper curve, solid line only); the curves are vertically displaced for clarity. (b) N-T junctions for $t = 0.2$ (upper curves) and $t = 0.5$ (lower curve).

Fig. 4. I-V characteristics of S-T junctions with and without magnetic field (dashed and solid lines respectively, $h = 0.2$). The lower solid and dashed curves are for $t = 0.2$ (without and with magnetic field) and the upper solid curve is for $t = 0.5$. The dotted lines are (i) the straight unitary slope line given for reference and (ii) the S-S characteristics for similar-parameters junctions (upper $t = 0.5$ and lower $t = 0.2$) given for comparison purposes. Insert: I-V characteristics of S-S junctions with different values of the left and right gap amplitudes, at finite temperature. The dotted line is the reference N-N characteristics and the dashed line is the curve for $\Delta_2 = \Delta_1$. The solid lines correspond to $\Delta_2 \neq \Delta_1$ ($\Delta_2 = 3\Delta_1$ for the line closest to the dashed one and $\Delta_2 = 37/3\Delta_1$ for the other one). The rounded cusp at $eV \approx \Delta_2 - \Delta_1$ is a thermally activated feature that appears when the upper gap edges at both sides of the junctions are aligned. Note that the position of the thermal cusp is not exactly $\Delta_2 - \Delta_1$, but it is shifted towards lower voltages. This shift is due to higher order Andreev processes.
of the junctions, and their current amplitude is found to be systematically smaller than in the case of the respective singlet-singlet junctions. The ‘sub-gap’ structure shows only two steps and the current is zero when $eV < \Delta_{\text{Triplet}}$ (the vector order parameter on the spin-triplet side of the junction). Concerning the effects of an applied magnetic field, the curves remain unchanged if the field is applied parallel to the direction of the vector-order-parameter, but show instead a Zeeman effect if the field is perpendicular to it (dashed line). To illustrate the effects of finite temperature on the tunnelling characteristics, we also show in the insert of Fig. 4 an I-V characteristics for a S-S junction at finite temperature. More details on finite temperature effects can be found in Ref. 26).

Regarding the organic superconductors,\(^1\) the experiments show that for magnetic fields along the direction of the conducting chains (\(a\)) the upper critical field is possibly paramagnetically limited for small fields (before crossing the upper critical field along \(b'\) that is never paramagnetically limited). In that range we could assume that the direction of the order parameter is fixed respect to the lattice and does not follow the applied field.\(^{38}\) With this geometry, a Zeeman splitting of the differential conductance peak should be observed in a normal-tip STM experiment. As the field is rotated the splitting would be suppressed and for a magnetic field oriented parallel to the \(b'\) crystalline-axis there should be no Zeeman effect. The disappearance of splitting, even as the field is possibly being increased, would constitute a clear signature of spin-triplet superconductivity. Similarly, an \(s\)-wave-tip STM would also be a direct probe for spin-triplet order. When a field is applied along the \(b'\) crystalline-axis, a Zeeman splitting would occur. This would constitute a clear sign of unconventional superconductivity since such an effect does not take place for standard BCS superconductors. The \(b'\) direction is the one on which the upper critical field is not paramagnetically limited, so relatively large fields could be applied in order to obtain a clear signal, and as the field alignment changes the splitting should disappear. No successful attempts of this kind of experiments were as yet made in the case of the quasi-one-dimensional organic salts, but efforts in this direction are on their way (recently, preliminary experiments involving junctions between two Bechgaard salts were performed, and they showed a number of puzzling features including a zero-bias conductance peak and zero excess current).\(^{56}\) We certainly hope that further tunnelling experiments along the lines described in this note can be performed.

§5. Conclusions and perspectives

In this note we have presented a new formalism, based on Keldysh technique, to compute the full current-voltage characteristics of point contact junctions between normal or superconducting (singlet or triplet) leads. This formalism is flexible enough to easily incorporate the effects of magnetic fields or temperature. As a possible application of this formalism we have shown how it can be used to devise a tunnelling experiment allowing to probe for the triplet nature of the superconducting phase in the Bechgaard salts.

Clearly, this formalism can be extended in many ways. One of the most inter-
esting extensions, in view of recent tunnelling experiments,\textsuperscript{56} is to take into account a finite region for the tunnelling contact, an intermediate between the point contact and the planar junction. Other extension, needed for example in the case of ruthenates, would be to incorporate the modifications of the superconducting order parameter close to the surface. Such a difficulty could be avoided for the organics given the large mass anisotropy, but is more crucial for an isotropic two-dimensional system, since surfaces are pair breaking for triplet superconductors when the momentum is perpendicular to the surface. Finally, extensions to other pairing symmetries such as \textit{d}-wave are also possible; such extensions are, of course, exciting questions to be covered in future studies.

Acknowledgements

We would like to thank Ø. Fischer, M. Eskildsen, M. Kugler, and G. Levy for discussions about tunneling and STM. We also thank Y. Maeno and M. Sigrist for discussions about tunneling into triplet superconductors and ruthenates in particular. Part of this work has been supported by the Swiss National Science Foundation under MaNEP and Division II.

References

1) For recent reviews on organics, see Chem. Rev. \textbf{104} (2004).
2) T. Giamarchi, \textit{Quantum Physics in One Dimension} (Oxford University Press, Oxford, 2004).
3) L. D. Landau, Sov. Phys. JETP \textbf{3} (1957), 920.
4) M. S. Dresselhaus, G. Dresselhaus and P. C. Eklund, \textit{Science of Fullerenes and Carbon Nanotubes} (Academic Press, San Diego, CA, 1995).
5) M. Bockrath, D. H. Cobden, J. Lu, A. G. Rinzler, R. E. Smalley, L. Balents and P. L. McEuen, Nature \textbf{397} (1999), 598.
6) Z. Yao, H. W. C. Postma, L. Balents and C. Dekker, Nature \textbf{402} (1999), 273.
7) S. Tarucha, T. Honda and L. Saku, Solid State Commun. \textbf{94} (1995), 413.
8) O. M. Auslaender, A. Yacoby, R. de Picciotto, K. W. Baldwin, L. N. Pfeiffer and K. W. West, Science \textbf{295} (2002), 825.
9) Y. Tserkovnyak, B. I. Halperin, O. M. Auslaender and A. Yacoby, Phys. Rev. Lett. \textbf{89} (2002), 136805.
10) M. P. A. Fisher and L. I. Glazman, in \textit{Mesoscopic Electron Transport}, ed. L. Kowenhoven et al. (Kluwer Academic Publishers, Dordrecht, 1997); cond-mat/9610037.
11) C. Glattli, in \textit{High Magnetic Fields: Applications in Condensed Matter Physics and Spectroscopy}, ed. C. Berthier et al. (Springer-Verlag, Berlin, 2002), p. 1.
12) F. P. Milliken, C. P. Umbach and R. A. Webb, Solid State Commun. \textbf{97} (1996), 309.
13) D. Jérome, A. Mazaud, M. Ribault and K. Bechgaard, J. of Phys. \textbf{41} (1980), L95.
14) T. Ishiguro, in \textit{Lecture Notes in Phys. 595}, ed. C. Berthier, L. P. Lévy and G. Martinez (Springer-Verlag, Heidelberg, 2002), p. 301.
15) M. Sigrist and K. Ueda, Rev. Mod. Phys. \textbf{63} (1991), 239.
16) T. M. Rice and M. Sigrist, J. of Phys.: Cond. Mat. \textbf{7} (1995), L643.
17) A. P. Mackenzie and Y. Maeno, Rev. Mod. Phys. \textbf{75} (2003), 657.
18) E. L. Wolf, “Principles of Electron Tunneling Spectroscopy”, \textit{International Series of Monographs on Physics}, second edition (Oxford Science Publications, Clarendon Press, Oxford, 1989).
19) G. Binnig and H. Rohrer, Rev. Mod. Phys. \textbf{71} (1999), S324.
20) G. E. Blonder, M. Tinkham and T. M. Klapwijk, Phys. Rev. B \textbf{25} (1982), 4515.
21) M. Octavio, M. Tinkham, G. E. Blonder and T. M. Klapwijk, Phys. Rev. B \textbf{27} (1983), 6739.
22) M. H. Cohen, L. M. Falicov and J. C. Phillips, Phys. Rev. Lett. 8 (1962), 316.
23) J. W. Wilkins, “Tunneling Phenomena in Solids”, Multiparticle Tunneling, first edition (Plenum Press, New York, 1969), chapter 24, p. 333.
24) J. C. Cuevas, A. Martín-Rodero and A. Levy Yeyati, Phys. Rev. B 54 (1996), 7366.
25) C. J. Bolech and T. Giamarchi, Phys. Rev. Lett. 92 (2004), 127001.
26) C. J. Bolech and T. Giamarchi, Phys. Rev. B 71 (2005), 024517.
27) C. Bourbonnais and D. Jérome, in Advances in Synthetic Metals, Twenty years of Progress in Science and Technology, ed. P. Bernier, S. Lefrant and G. Bidan (Elsevier, New York, 1999), p. 206; cond-mat/9903101.
28) D. Jérome, Organic superconductors: From (TMTSF)$_2$PF$_6$ to fullerenes (Marcel Dekker, New York, 1994), p. 405.
29) D. Jérome, Chem. Rev. 104 (2004), 5565.
30) M. Dressel, K. Petukhov, B. Salameh, P. Zornoza and T. Giamarchi, Phys. Rev. B 71 (2005), 075104.
31) M. Dressel, A. Schwartz, G. Grüner and L. Degiorgi, Phys. Rev. Lett. 77 (1996), 398.
32) A. Schwartz, M. Dressel, G. Grüner, V. Vescoli, L. Degiorgi and T. Giamarchi, Phys. Rev. B 58 (1998), 1261.
33) T. Giamarchi, Physica B 230-232 (1997), 975.
34) T. Giamarchi, Chem. Rev. 104 (2004), 5565.
35) The Fifth International Symposium on Crystalline Organic Metals, Superconductors and Ferromagnets, Port-Bourgenay, September 21st-26th, 2003.
36) J. I. Oh and M. J. Naughton, Phys. Rev. Lett. 96 (2004), 067001.
37) N. Joo, P. Auban-Senzier, C. R. Pasquier, P. Monod, D. Jérome and K. Bechgaard, Eur. Phys. J. B 40 (2004), 43.
38) A. G. Lebed, K. Machida and M. Ozaki, Phys. Rev. B 62 (2000), R795.
39) R. W. Cherng and C. A. R. Sá de Melo, Phys. Rev. B 67 (2003), 212505.
40) Y. Tanaka and K. Kuroki, Phys. Rev. B 70 (2004), 060502(R).
41) D. Podolsky, E. Altman, T. Rostunov and E. Demler, Phys. Rev. Lett. 93 (2004), 246402.
42) T. Vuletić, P. Auban-Senzier, C. Pasquier, S. Tomić, D. Jérome, M. Hérmand and K. Bechgaard, Eur. Phys. J. B 25 (2002), 319.
43) I. J. Lee, S. E. Brown, W. G. Clark, M. J. Strouse, M. J. Naughton, W. Kang and P. M. Chaikin, Phys. Rev. Lett. 88 (2002), 017004.
44) I. J. Lee, P. M. Chaikin and M. J. Naughton, Phys. Rev. B 65 (2002), 180502(R).
45) J. Nicol, S. Shapiro and P. H. Smith, Phys. Rev. Lett. 5 (1960), 461.
46) T. M. Klapwijk, G. E. Blonder and M. Tinkham, Physica B & C 109-110 (1982), 1657.
47) E. Scheer, P. Joyez, D. Esteve, C. Urbina and M. H. Devoret, Phys. Rev. Lett. 78 (1997), 3535.
48) E. Scheer, N. Agraït, J. C. Cuevas, A. Levy Yeyati, B. Ludoph, A. Martín-Rodero, G. Rubio-Bollinger, J. M. van Ruitenbeek and C. Urbina, Nature 394 (1998), 154.
49) B. Ludoph, N. van der Post, E. N. Bratus’’, E. V. Bezuglyi, V. S. Shumeiko, G. Wendin and J. M. van Ruitenbeek, Phys. Rev. B 61 (2000), 8561.
50) E. Scheer, W. Belzig, Y. Naveh, M. H. Devoret, D. Esteve and C. Urbina, Phys. Rev. Lett. 86 (2001), 284.
51) G. Rubio-Bollinger, C. de las Heras, E. Bascones, N. Agraït, F. Guinea and S. Vieira, Phys. Rev. B 67 (2003), 121407R.
52) M. Hafner, P. Konrad, F. Pauly, J. C. Cuevas and E. Scheer, cond-mat/0407207.
53) L. V. Keldysh, Sov. Phys. - JETP 24 (1965), 1018.
54) K. Sengupta, I. Žutić, H.-J. Kwok V. M. Yakovenko and S. Das Sarma, Phys. Rev. B 63 (2001), 144531.
55) Y. Tanuma, K. Kuroki, Y. Tanaka, R. Arita, S. Kashiwaya and H. Aoki, Phys. Rev. B 66 (2002), 094507.
56) H. I. Ha, J. I. Oh, J. Moser and M. J. Naughton, Synth. Metal 137 (2003), 1215.