Anomalous Conductances in an Ultracold Quantum Wire

The Harvard community has made this article openly available. Please share how this access benefits you. Your story matters

| Citation       | Kanász-Nagy, M., L. Glazman, T. Esslinger, and E. A. Demler. 2016. “Anomalous Conductances in an Ultracold Quantum Wire.” Physical Review Letters 117 (25). https://doi.org/10.1103/physrevlett.117.255302. |
|----------------|-------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------|
| Citable link   | http://nrs.harvard.edu/urn-3:HUL.InstRepos:41412219                                                                                                                                               |
| Terms of Use   | This article was downloaded from Harvard University’s DASH repository, and is made available under the terms and conditions applicable to Open Access Policy Articles, as set forth at http://nrs.harvard.edu/urn-3:HUL.InstRepos:dash.current.terms-of-use#OAP |
Anomalous conductances in an ultracold quantum wire

M. Kanász-Nagy$^1$, L. Glazman$^2$, T. Esslinger$^3$, and E. A. Demler$^1$

$^1$Department of Physics, Harvard University, Cambridge, MA 02138, U.S.A.
$^2$Department of Physics, Yale University, New Haven, CT 06520, U.S.A.
$^3$Department of Physics, ETH Zurich, 8093 Zurich, Switzerland

We analyze the recently measured anomalous transport properties of an ultracold gas through a ballistic constriction [S. Krinner et al., PNAS 201601812 (2016)]. The quantized conductance observed at weak interactions increases several-fold as the gas is made strongly interacting, which cannot be explained by the Landauer theory of single-channel transport. We show that this phenomenon is due to the multichannel Andreev reflections at the edges of the constriction, where the interaction and confinement result in a superconducting state. Andreev processes convert atoms of otherwise reflecting channels into the condensate propagating through the constriction, leading to a significant excess conductance. Furthermore, we find the spin conductance being suppressed by superconductivity; the agreement with experiment provides an additional support for our model.

PACS numbers: 67.10.Jn, 67.85.De, 68.65.La, 74.25.F-
est plateau of a single transverse mode. This is in apparent contradiction with the simple Blonder–Tinkham–Klapwijk (BTK) model of transport through a single ballistic channel [33, 34]: although interactions can make the channel superconducting (SC), this can at most lead to a factor of 2 increase in the conductance, since in AR each incident atom drags along at most another atom through the constriction, as a Cooper pair.

We resolve the puzzle of anomalous conductance by associating it with multichannel AR processes at the normal-superconductor interfaces at the two ends of the constriction (see Fig. 1). Confinement significantly renormalizes the interactions within the central part of the constriction, leading to strong SC pairing. This pairing field penetrates into the normal leads, with several channels below the Fermi energy. Atoms in higher transverse modes, that would otherwise be reflected by the constriction, can go through AR processes within this thin superconducting interface. As they become part of the condensate they propagate through the junction as Cooper pairs [35, 36]: the resistance of the channel is entirely determined by the interface [37]. Furthermore, as the interaction increases, current is increasingly carried by Cooper pairs, the spin current approaches zero. This agrees with the experimental observations of Ref. 22.

The experimental geometry is shown in Fig. 1(a). The central part of the gas is squeezed into two dimensions using lithographic imprinting, whereas a narrower perpendicular laser beam pinches the middle of this region to form a short one-dimensional ballistic quantum wire [20, 22]. The conductance of the wire is tunable either by tuning the confinement frequencies $\nu_x$, $\nu_z$, or using a gate potential $V_g$, created by an additional, wide laser beam along the $z$ axis (see the caption of Fig. 2). By creating a density or spin imbalance between the two sides of the junction, the conductance $G_a$ and spin conductance $G_s$ can be determined by monitoring the relaxation of the population imbalance in time, and making use of the equation of states of the gas within the leads [20, 21].

We determine the superconducting profile in the constriction within the local density approximation (LDA), whereby we consider a small part of the system of length $L_y$, where the parameters of the gas are assumed to be constant. We also take into account the renormalization of interactions due to confinement effects. The constriction is described by a harmonic Hamiltonian of trapping frequencies $\omega = (\omega_x, \omega_z) = (\nu_x, \nu_z)/2\pi$, local gate potential $V_g$ and chemical potential $\mu$,

$$H_{\text{kin}} = \sum_{n,\sigma,q} \xi_{n,\sigma,q} a_{n,\sigma,q}^\dagger a_{n,\sigma,q},$$

where $\xi_{n,q} = \hbar^2 q^2/(2m) - V_g - \mu + (n_x + \frac{1}{2})\hbar \omega_x + (n_z + \frac{1}{2})\hbar \omega_z$ denotes the channel energies, and $a_{n,\sigma,q}$ annihilates an atom in channel $n = (n_x, n_z)$, spin $\sigma = \uparrow, \downarrow$ and momentum $q$ along the $y$ axis. The interaction between the $^6$Li atoms is given by the standard point interaction $g \delta^{(3)}(r)$, where $g$ is the bare interaction strength [38]. In order to simplify the treatment of the interaction term, it is worth going into the center of mass (COM) and relative frame of the colliding atoms along the trapped directions, $(x_1, x_2) \rightarrow (x_1 + x_2, x_1 - x_2)$, and similarly for $z$, with the coordinates of the atoms denoted by $(x_1, z_1)$ and $(x_2, z_2)$. One can thus transform the interaction Hamiltonian according to the unitary transformation $(\mathbf{N}, \nu) \rightarrow (\mathbf{N}, \nu)$, where $\mathbf{N} = (N_x, N_z)$ and $\nu = (\nu_x, \nu_z)$ denote the COM and relative harmonic oscillator quantum states, and $\mathbf{n}_1, \mathbf{n}_2$ stand for those in the laboratory frame. These matrix elements are non-zero only for $\mathbf{n}_1 + \mathbf{n}_2 = \mathbf{N} + \nu$ combinations, due to energy conservation. Since harmonic trapping and interactions both conserve $\mathbf{N}$ and the COM momentum $Q$, the interaction Hamiltonian can be decoupled exactly as [39]

$$H_{\text{int}} = \frac{1}{\tilde{g}} \sum_{N,Q} \hat{\Delta}^\dagger_{N,Q} \hat{\Delta}_{N,Q},$$

$$\hat{\Delta}_{N,Q} \equiv \tilde{g} \sum_{\mathbf{n}_1, \mathbf{n}_2, k} V^\dagger_{\mathbf{N}} a_{\mathbf{n}_1, \lambda, Q - k} a_{\mathbf{n}_2, \lambda, k},$$

where the interaction strength $\tilde{g} = g/(l_x L_y l_z)$ of energy dimension is defined using oscillator lengths $l_x(z) = \sqrt{\hbar / \omega_x(z)}$. The matrix elements $V^\dagger_{\mathbf{N}} = \varphi_{\nu_x}(0) \varphi_{\nu_z}(0) |(\mathbf{N}, \nu)\rangle \langle \mathbf{n}_1, \mathbf{n}_2|$, with $\nu = \mathbf{n}_1 + \mathbf{n}_2 - \mathbf{N}$, arise from the matrix elements of the point-like interaction.

FIG. 2. (Color online.) (a) SC pairing amplitude $\Delta_0$ (lines with symbols) across the constriction at different interaction strengths shown in the inset of (b). Solid lines indicate the energies of the transverse modes of different $n_z$ quantum numbers, with only $n_z = 0, 1$ and 2 modes shown. At the edges of the constriction, the gas is quasi-two dimensional, and the $n_z$ modes are almost completely degenerate. They split near the middle, where the gas becomes quasi-one dimensional. (b) Bound state energies along the constriction, renormalized by the confinement. [Parameters: $\nu_x(y), \nu_z(y)$ and $V_g(y)$ are approximated as Gaussians of HWHM $(d_x, d_y, d_v) = (4.7, 17.7, 15.1) \mu$m, as in Ref. 22 and heights $(V_{g0}, V_{g0}, V_{g0}) = (23.2 \text{ kHz}, 9.2 \text{ kHz}, 0.625 \mu\text{K})$. Cut-off in the channel number: $n_x, n_z \leq 8$. $\lambda_F$ and $E_F$ denote the Fermi wavelength and the Fermi energy within the reservoirs.]
potential [39] (see Supplementary Material). In the previous expression, the value of the relative harmonic oscillator wave function $\varphi_\nu$ at the origin is given by $\varphi_\nu(0) = \frac{(-1)^{\nu/2}}{(2\pi)^{1/4}} \sqrt{\frac{1}{r^2}} \, \varphi_\nu$ for $\nu$ even and $\varphi_\nu(0) = 0$ for $\nu$ odd.

We decouple Eq. (2) in a standard BCS approximation $\Delta_N = (\Delta_{N,q=0})$. Although in general it could be possible to have SC ordering in many COM modes, we verified that for the experimental parameters of Ref. [22] considered here, only the $N = 0$ mode gains non-zero pairing amplitude. Thus, in the following, we focus on this case and leave the general discussion to the Supplementary Material. The resulting Bogoliubov–de Gennes Hamiltonian for quasi-particle excitations reads

$$H_{\text{BdG}} = \sum_q \left( a_{\uparrow,q}^\dagger a_{\downarrow,-q} - a_{\downarrow,q}^\dagger a_{\uparrow,-q} \right) \left( \xi_{\uparrow,q} \Delta - \xi_{\downarrow,q} \right),$$

in vectorial notation for the band indices. Here, $(a_{\sigma,q})_n = \delta_{n,q}$ denotes the vector of annihilation operators, the SC matrix is given by $\Delta_{n,n'} = \Delta_0 \delta_{n,n'}$, and the matrix $\xi_{\sigma,q}$ contains the band energies on its diagonal. Using a Bogoliubov transformation, one can now determine the quasi-particle energies $E_{n,q}$. Then, in order to determine the pairing amplitude $\Delta_0$, one needs to minimize the finite temperature BCS free energy $F_{\text{MF}} = E_{\text{MF}} - T \sum_q \log \left( 1 + e^{-E_{n,q}/T} \right)$ at a fixed chemical potential, as set by the leads. The mean-field condensation energy $E_{\text{MF}} = \sum_q \left( \xi_{q,n} - E_{n,q} \right) - \frac{\Delta_0^2}{g}$ however still contains the bare interaction term $g$, and a divergent sum over excitation energies. In order to regularize this term, we make use of the vacuum Bethe–Salpeter equations [40–41], and express $g$ in terms of physical quantities: the scattering length $a$, or, equivalently, the vacuum bound state energy $E_B$ (see Supplementary Material),

$$\frac{1}{g} = \frac{m}{4\pi \hbar^2 a} - \int \frac{d^3 q}{(2\pi)^3} \frac{m}{\hbar^2 q^2 + i0^+}$$

$$= \frac{1}{l_x l_z} \int \frac{dq}{2\pi} \sum_{n_1 n_2} \frac{|V_{n_1 n_2}|^2}{E_B - \left( \frac{\hbar^2 q^2}{m} + h(n_1 + n_2) \right)}.$$  

In contrast to three-dimensional systems, Eq. (5) always has a bound state solution $E_B > 0$ in quasi-one dimensional gases, even on the attractive side of the Feshbach resonance [38–42, 43]. As we show in Fig. 2 (b), $E_B$ strongly depends on the confining frequencies as well as on the scattering length, and incorporates the confinement-induced renormalization of the interaction. Making use of Eq. (5), we can now express the condensation energy in terms of $E_B$, and, as we show in the Supplementary Material, the resulting expression is regular,

$$E_{\text{MF}} = \sum_{n,q} (\xi_{q,n} - E_{n,q}) - \sum_{n_1 n_2, q} \frac{|\Delta_0|^2 |V_{n_1 n_2}|^2}{E_B - \left( \frac{\hbar^2 q^2}{m} + h(n_1 + n_2) \right)}.$$  

FIG. 3. (Color online.) Conductance (a) and spin conductance (b) as a function of the gate potential at $1/(k_F r_e a) = 9.2, 23.2$ kHz. The geometry is identical to the one in Fig. 2. [Parame-

ters: $T = 62$ nK, $\mu = 8.5$ kHz, $(\nu_{\text{vib}}, \nu_{\text{opt}}) = (9.2, 23.2)$ kHz, the geometry is identical to the one in Fig. 2.

Fig. 2 (a) shows typical profiles of the SC order parameter $\Delta_0(y)$ at various interaction strengths. Due to strong confinement towards the middle of the constriction, the bound state becomes significantly deeper in energy favoring superconductivity in Eq. (6). Although in the middle there is only one channel below the Fermi energy that can contribute to pairing, higher transverse modes are also coupled to the condensate in the SC-normal interface through Eq. (1). At the largest interaction strengths, the SC gap becomes comparable to the Fermi energy [43]. This strong pairing also extends around the central potential hill of the constriction, providing a thin superconducting layer that is responsible for the excess conductance seen in the experiment [22], due to multichannel Andreev processes. The length scale over which these processes happen are of the order of the SC healing length $\xi_s$. Even though the width of this region is just a few times the Fermi wavelength $\lambda_F$, the strong pairing within the constriction leads to $\xi_s \sim \lambda_F$, and the AR probabilities become non-negligible.

We determine the conductance and spin conductance of the waveguide in a Landauer picture [33], by calculating the reflection and AR coefficients $(r_{pp/n' n})$ and $(r_{hp})$, respectively, describing reflections from channel $n$ to $n'$, with the $p$ and $h$ indices denoting particle and hole states. To do this, we determine the eigenmodes of the Bogoliubov–de Gennes Hamiltonian Eq. (4) at all incoming energies $\epsilon$, (see Supplementary Material). The zero bias conductance and spin conductance are given by a thermal average over these contributions [33].

$$G_{n/s} = -\int d\epsilon \left( \epsilon - \mu \right) \text{Tr} \left( 1 - r_{pp}^\dagger r_{pp} \pm r_{hp}^\dagger r_{hp} \right)$$

where $\hat{1}$ denotes the unit matrix, $n_F$ stands for the Fermi function, and the energy arguments of $r_{pp}(\epsilon)$ and $r_{hp}(\epsilon)$ are neglected for brevity. As can be seen from Eq. (7), AR processes contribute to the conductance, but they decrease the spin conductance. The definition of the spin conductance in Eq. (7) differs from that of Ref. [22] by a
Furthermore, in agreement with experiment [22], we find the conduc-
tance plateau at weak interactions. In agreement with the experiment [22], the conduc-
tance plateau is still somewhat visible at larger interaction strengths, but pushed to a much larger value due to superconduc-
tivity (see the curves $1/(k_{F,0}a) = -0.70$ and $-0.75$). However, we also find an interesting non-monotonicity of the conduc-
tance curves at strong confinement, that has not been observed experimentally. This behavior is due to the confinement-induced renormalization of the interaction, that leads to the onset of SC at tighter con-
finements. This is accompanied by a sudden decrease in the spin conductance (see Fig. 4 (b)). This non-
monotonicity does not appear at higher temperature as the confinement-induced onset of pairing is killed by tem-
perature fluctuations, see the inset of Fig. 4 (a). This effect thus may be observable by further cooling the gas in the experiment.

The comparison of Fig. 4 (a) and the inset also demonstrates the sensitivity of the conductance curves to exper-
imental parameters, as also seen in Ref. [21]. As we show in the Supplementary Material, conductance at strong
interactions, $1/(k_{F,0}a) \sim 0.5$, can change as large as a factor of 5 just by changing the temperature and chemi-
potential within their $\sim 15\%$ experimental error bars.

The reason is that $\Delta_0$ depends very sensitively on these parameters near the onset of superconductivity, and its value has a significant influence on conductance. Further
important uncertainties arise from experimental aber-
rations of the laser fields that form the constriction. Since
the transport is largely governed by an interface effect at the edge of the constriction, these geometric factors become important [21].

As an experimental test of our theory, we propose to
investigate the channel’s conductance at large, equal spin
imbalances in both leads, leading to the suppression of the constriction’s superconductivity due to Fermi surface
mismatch. At large imbalances, the SC-normal transition could thus be measured using the drop of anomalous con-
ductances, and from the increase of spin conductance, to
their respective values in the normal state [15, 17].

The above analysis of quantum transport assumes a static order parameter in the superconducting region. Its
finite size may constrain the fluctuations of the number of atoms in the region. The constrained particle number
fluctuations enhances the fluctuations of phase of the order parameter. These effects were studied exten-
sively in the context of Coulomb blockade in a super-
conducting island coupled to a normal-metal lead, see, e.g. Refs. 18 and 19. The overall conclusion is that at large conductance of the interface the effects of Coulomb
blockade (i.e., constraints on the particle number) are negligible. The corresponding energy scale turns out to
scale as exponent of $-G/G_0$ if the large conductance of a junction is achieved by increasing the number of con-
ducting channels [15, 19], and as a product of reflection
amplitudes in each of the channels, in case of an arbi-

Factor of two, due to the ambiguity in defining the chemi-
cal potential difference in case of the spin current. Using
the definition above, the spin and charge conductances are identical in the normal phase, and their deviation
indicates the onset of superconductivity.

As shown in Fig. 3, both $G_s$ and $G_{\alpha}$ show the usual Landauer quantization as a function of the gate poten-
tial $V_{g0}$ at weak interactions, as has been observed experimen-
tally [20]. At increasing interaction strengths, the constriction becomes superconducting, leading to in-
creased conductance and suppressed spin conductance. As $V_{g0}$ is tuned, SC order appears first in the middle of the constriction (see Fig. 2), thus only the otherwise transmitting channels can participate in Andreev
processes. This is the regime of the BTK theory, and we
observe well defined plateaus, within a factor of two in-
crease in conductance. At larger gate potentials, how-
ever, the number of channels in the superconducting
interface increases, leading to a strong increase in conduc-
tance. Since the SC layer at the end of the wire is thin, most channels cannot go through perfect ARs and they
only contribute a small fraction of a conductance quantum to $G_s$. The plateaus thus become less well-defined.
Furthermore, in agreement with experiment [22], we find
that $G_s$ depends non-monotonically on the gate poten-
tial in Fig. 3 (b). The reason for this is that as $V_{g0}$ increases, additional channels are pulled down below the Fermi energy, and the system gains additional condensa-
tion energy by forming Cooper pairs in these channels.
As a result, SC pairing increases, and a larger fraction
of the current is carried by Cooper pairs, leading to a
sudden drop in $G_s$.

Fig. 4 shows $G_s$ and $G_{\alpha}$ as a function of the horizontal conefinement $\nu_{z0}$, exhibiting a broad conductance plateau

The comparison of Fig. 4 (a) and the inset also demon-
strates the sensitivity of the conductance curves to exper-
imental parameters, as also seen in Ref. [21]. As we show in the Supplementary Material, conductance at strong
interactions, $1/(k_{F,0}a) \sim 0.5$, can change as large as a factor of 5 just by changing the temperature and chemi-
potential within their $\sim 15\%$ experimental error bars.

The reason is that $\Delta_0$ depends very sensitively on these parameters near the onset of superconductivity, and its value has a significant influence on conductance. Further
important uncertainties arise from experimental aber-
rations of the laser fields that form the constriction. Since
the transport is largely governed by an interface effect at the edge of the constriction, these geometric factors become important [21].

As an experimental test of our theory, we propose to
investigate the channel’s conductance at large, equal spin
imbalances in both leads, leading to the suppression of the constriction’s superconductivity due to Fermi surface
mismatch. At large imbalances, the SC-normal transition could thus be measured using the drop of anomalous con-
ductances, and from the increase of spin conductance, to
their respective values in the normal state [15, 17].

The above analysis of quantum transport assumes a static order parameter in the superconducting region. Its
finite size may constrain the fluctuations of the number of atoms in the region. The constrained particle number
fluctuations enhances the fluctuations of phase of the order parameter. These effects were studied exten-
sively in the context of Coulomb blockade in a super-
conducting island coupled to a normal-metal lead, see, e.g. Refs. 18 and 19. The overall conclusion is that at large conductance of the interface the effects of Coulomb
blockade (i.e., constraints on the particle number) are negligible. The corresponding energy scale turns out to
scale as exponent of $-G/G_0$ if the large conductance of a junction is achieved by increasing the number of con-
ducting channels [15, 19], and as a product of reflection
amplitudes in each of the channels, in case of an arbi-
terary (even small) number of highly-transparent channels \[29\, 50\]. The phase fluctuations are small, and their estimate in the Gaussian approximation is provided in Section 6 of the Supplementary.

**Conclusion** – We demonstrated that the recently observed anomalous transport measured in Ref. 22 is the result of a subtle interface effect at the ends of the ballistic wire, that becomes superconducting due to confinement-induced renormalization of interactions. Since SC penetrates in the quasi-two dimensional part of the lead, channels that would otherwise be reflected by the constriction can participate in Andreev processes, thus delivering Cooper pairs to the condensate which propagate through the interior part of the channel as a spinless superfluid. We could also explain non-monotonicities in the spin-conductance curve as the gate potential was changed, and predict additional non-monotonicities of the conductance curve as the confinement frequency at low temperatures.

**Acknowledgements:** Enlightening discussions with F. Pientka, S. Gopalakrishnan, J.-P. Brantut, M. Lebrat, S. Ninomiya, and D. Husmann are gratefully acknowledged. The work of E. D. and M. K.-N. was supported by the Simons Foundation and the Humboldt Foundation. E. D. also acknowledges support from Dr. Max Rössler, AFOSR Quantum Simulation MURI, the ARO-MURI Quism program. E. D. acknowledges the Staatssekretariat für Bildung, Forschung und Innovation SBFI for the support of the Horizon2020 project Quantum simulations of insulators and conductors QUIC.

---

[1] B. J. van Wees et al., Phys. Rev. Lett. 60, 848 (1988).
[2] M. König et al., Science 318, 766 (2007).
[3] K. Schwab, E. A. Henriksen, J. M. Worlock, and M. L. Roukes, Nature 404, 974 (2000).
[4] A. F. Andreev, Sov. Phys. JETP 19, 1228 (1964).
[5] H. Meisner, Phys. Rev. Lett. 2, 458 (1959).
[6] M. Tinkham, "Introduction to Superconductivity" (Courier Dover, 2012).
[7] C. W. J. Beenakker, Rev. Mod. Phys. 80, 1337 (2008).
[8] H. Shiba, Prog. theor. Phys. 40, 435 (1968).
[9] T. M. Eiles, J. M. Martinis, and M. H. Devoret, Phys. Rev. Lett. 70, 1862 (1993).
[10] F. W. J. Hekking, L. I. Glazman, K. A. Matveev, and R. I. Shekhter, Phys. Rev. Lett. 70, 4138 (1993).
[11] J. M. Hermenrother, M. T. Tuominen, and M. Tinkham, Phys. Rev. Lett. 72, 1742 (1994).
[12] G. Binnig et al., Phys. Rev. Lett. 55, 991 (1985).
[13] J. Linder and J. W. A. Robinson, Nat. Phys. 11, 307 (2015).
[14] L. Hofstetter, S. Csonka, J. Nygard and C. Schnenerberger Nature 461, 960 (2009).
[15] L. Fu and C. L. Kane, Phys. Rev. Lett. 100, 096407 (2008).
[16] V. Mourik, K. Zuo, S. M. Frolov, S. R. Plissard, E. P. A. M. Bakkers and L. P. Kouwenhoven, Science 336, 1003 (2012).
[17] M. T. Deng, C. L. Yu, G. Y. Huang, M. Larsson, P. Caroff and H. Q. Xu, Nano Letters 12, 6414 (2012).
[18] A. Das, Y. Ronen, Y. Most, Y. Oreg, M. Heiblum, H. Shtrikman, Nature Physics 8, 887 (2012).
[19] S. Nadji-Perge et al., Science 346 602 (2014).
[20] S. Knirrer et al., Nature 517, 64 (2015).
[21] D. Husmann et al., Science 357, 1498 (2015).
[22] S. Knirrer et al., Proc. Natl. Acad. Sci. U.S.A., 201601812 (2016).
[23] A. Bezryadin, C. N. Lau and M. Tinkham, Nature 404, 971 (2000).
[24] P. Cadden-Zimansky and V. Chandrasekhar, Phys. Rev. Lett. 97, 237003 (2006).
[25] T. L. Hylton et al., Appl. Phys. Lett. 53, 1343 (1988).
[26] C. J. Lobb, D. W. Abraham, and M. Tinkham, Phys. Rev. B 27, 150 (1983).
[27] S. Chakravarty, S. Kivelson, G. T. Zimanyi and B. I. Halperin, Phys. Rev. B 35, 7256 (1987).
[28] D. Spivak, A. Zvyuzin and M. Hruška, Phys. Rev. B 64, 132502 (2001).
[29] Feigel’Man, M. V., A. I. Larkin, and M. A. Skvortsov, “Quantum superconductormetal transition in a proximity array,” Physics-Uspekhi 44.10S (2001): 99.
[30] H. Hilgenkamp and J. Mannhart, Rev. Mod. Phys. 74, 485 (2002).
[31] Yonatan Dubi, Yigal Meir and Yshai Avishai, Nature 449, 876 (2007).
[32] G. R. Boogaard, A. H. Verbruggen, W. Belzig, and T. M. Klapwijk Phys. Rev. B 69, 220503(R).
[33] Y. V. Nazarov, Y. M. Blanter, Quantum transport: Introduction to Nanoscience (Cambridge University Press, 2009).
[34] G. E. Blonder, M. Tinkham, and T. M. Klapwijk, Phys. Rev. B 25, 4515 (1982).
[35] R. H. Parmenter, Phys. Rev. 118, 1173 (1960).
[36] V. Mourik et al., Science 336, 1003 (2012).
[37] C. W. J. Beenakker "Why Does a Metal/Superconductor Junction Have a Resistance?" in Quantum Mesoscopic Phenomena and Mesoscopic Devices in Microelectronics 51-60 (Springer Netherlands, 2000).
[38] I. Bloch, J. Dalibard, and W. Zwerger, Rev. Mod. Phys. 80, 885 (2008).
[39] A. M. Fischer and M. M. Parish, Phys. Rev. A 88, 023612 (2013).
[40] V. Pietilä, D. Pekker, Y. Nishida, and E. A. Demler, Phys. Rev. A 85, 023621 (2012).
[41] M. Kanász-Nagy, E. A. Demler, and G. Zaránd, Phys. Rev. A 91, 032704 (2015).
[42] M. Oshshani, Phys. Rev. Lett. 81, 938 (1998).
[43] E. Haller et al., Science 325, 1224 (2009).
[44] D. K. Efetov, et al. "Specular interband Andreev reflections at van der Waals interfaces between graphene and NbSe₂. ", Nat. Phys. (2015).
[45] D. S. Sarma, J. Phys. Chem. Solids 24, 1029 (1963).
[46] A. M. Clogston, J. Phys. Chem. Solids 24, 1029 (1963).
[47] Y. Shim, H. C. Schunck, A. Schirotzek and W. Ketterle, Nature 451, 689 (2008).
[48] G. Schön and A. D. Zaikin, Phys. Rep. 198, 237 (1990).
[49] M. V. Feigelman, A. Kamenev, A. I. Larkin, M. A. Ivanov, and V. M. Vinokur, Phys. Rev. B 42, 9385 (1990).
Skvortsov, Phys. Rev. B 66, 054502 (2002).
[50] Y. V. Nazarov, Phys. Rev. Lett. 82, 1245 (1999).