Quantized charge pumping in superconducting double barrier structure: Non-trivial correlations due to proximity effect

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PACS 73.23.-b – Electronic transport in mesoscopic systems
PACS 74.45.+c – Proximity effects; Andreev effect; SN and SNS junctions
PACS 71.10.Pm – Fermions in reduced dimensions (anyons, composite fermions, Luttinger liquid, etc.)

Abstract. - We consider quantum charge pumping of electrons across a superconducting double barrier structure in the adiabatic limit. The superconducting barriers are assumed to be reflection-less so that an incident electron on the barrier can either tunnel through it or Andreev reflect from it. In this structure, quantum charge pumping can be achieved (a) by modulating the amplitudes, \(\Delta_1\) and \(\Delta_2\) of the two superconducting gaps or alternatively, (b) by a periodic modulation of the order parameter phases, \(\phi_1\) and \(\phi_2\) of the superconducting barriers. In the former case, we show that the superconducting gap gives rise to a very sharp resonance in the transmission resulting in quantization of pumped charge, when the pumping contour encloses the resonance. On the other hand, we find that quantization is hard to achieve in the latter case. We show that inclusion of weak electron-electron interaction in the quantum wire leads to renormalisation group evolution of the transmission amplitude towards the perfectly transmitting limit due to proximity effects. Hence as we approach the zero temperature limit, we get destruction of quantized pumped charge. This is in sharp contrast to the case of charge pumping in a double barrier in a Luttinger liquid where quantized charge pumping is actually achieved in the zero temperature limit. We also propose an experimental set-up to study these effects.

Introduction. – The phenomena of quantum charge pumping corresponds to a net flow of DC current between different electron reservoirs (at zero bias) connected via a quantum system whose (at least two) system parameters are periodically modulated in time [1]. The zero bias current is obtained in response to the time variation of the parameters of the quantum system which explicitly break the time reversal symmetry. When the frequency of modulation is smaller than any characteristic scale associated with the electron, adiabatic limit is reached. In this limit, the pumped charge in a unit cycle becomes independent of the pumping frequency. This is referred to as “adiabatic charge pumping”. It is worth mentioning that breaking of the time-reversal symmetry is necessary but not a sufficient condition in order to get a net pumped charge in a unit cycle. For obtaining a net pumped charge, parity must also be broken. In this letter, we explore possibilities of obtaining quantized pumped charge in unit of e for the double superconducting barrier system in one-dimensional (1–D) quantum wire (QW). We show that electron–electron (e–e) interactions in the wire lead to destruction of the quantized pumped charge as one goes to low temperatures. This is in sharp contrast to the case of normal double barrier system in a Luttinger liquid [2–4]. Pumping of free electrons across 1–D quantum well and multi-wire junction was studied earlier in Ref. [5–7], where using Brouwer’s formula [8], it was shown that the pumped charge can be expressed as a sum of two contributions, viz., a dissipative part and a quantized topological part, the latter being independent of the details of the pumping contour [9]. The dissipative part was found to be proportional to the conductance through the system on the pumping contour while the topological part was non-zero only if the pumping contour enclosed a resonance. Hence in order to obtain quantized pumped charge, one needs to reduce the dissipative part as much as possible. This is very naturally achieved if one considers pumping through a quantum well in a 1–D interacting electron gas [3] (Lut-
Luttinger liquid) as the interaction correlations make the resonance very sharp thereby reducing the conductance on the contour to zero in the zero temperature limit. This leaves behind a quantized topological part. The power law dependence of the pumped charge was shown to converge to a quantized value asymptotically. This was obtained using a perturbative approach in case of a weakly interacting electron gas followed by “poor-man’s scaling approach”. In this letter, we show that proximity effects give rise to a power law reduction from the quantized value (as opposed to enhancement) of the pumped charge as the temperature is varied.

**Proposed device and its theoretical modelling.** We consider pumping of electrons (in the adiabatic limit) across a superconducting double barrier (SDB) structure, as depicted in Fig. 1. Experimentally such a SDB structure can be realised by depositing thin strips of superconducting material on top of a single ballistic QW (like carbon nanotubes) at two places, which can induce a finite superconducting gap in the barrier regions of the QW as a result of proximity of the superconducting strips. We consider the superconducting barrier (SB) to be reflection-less so that an incident electron on the barrier can either tunnel through it or Andreev reflect from it. In the proposed geometry, we explore two scenarios to achieve quantization of pumped charge — (a) by periodic modulation of amplitudes $\Delta_1$ and $\Delta_2$ of the gap at the two SB or alternatively, (b) by periodic modulation of the order parameter phases $\phi_1$ and $\phi_2$ associated with the two SB. The latter can be realised experimentally by exploiting a Josephson junction [10] set-up. In the former case, it should be possible to simulate the situation experimentally by putting a top gate at the region of each SB, which can modulate the local chemical potential of the electrons in the barrier region and hence modulate the relative height of the gap with respect to the Fermi energy of the electron in the QW (see Fig. 1). We show that in the $\Delta_1 - \Delta_2$ plane, there is an isolated sharp resonance. However, the transmission across the double barrier has a line of sharp resonances in the $\phi_1 - \phi_2$ plane. As mentioned earlier, in order to obtain quantized pumped charge, the transmission on the pumping contour should be as small as possible. When we consider $\Delta_1$ and $\Delta_2$ as the pumping parameters, we can always choose a pumping contour which completely encloses the isolated resonance and hence it is possible to achieve quantization of charge. However, in the $\phi_1 - \phi_2$ plane, we have a line of resonances. Any closed contour enclosing the resonances will surely cross the resonance line at least twice thereby increasing the dissipative part and consequently resulting in destruction of quantization of pumped charge. Interestingly enough, inclusion of weak $e-e$ interaction results in a RG flow of the transmission amplitude towards perfectly transmitting limit due to proximity effects as we go to zero temperature limit. Hence the sharpness of resonance is lost due to RG enhancement of transmission through individual SB resulting in complete destruction of quantized charge pumping as we go down in temperature. It is worth noticing that the consequence of inclusion of correlations due to $e-e$ interaction is just opposite here with respect to the case of double barrier in a Luttinger liquid [3].

**Superconducting barrier.** Quantum transport in SB structure was considered in past in Ref. [11]. Here we consider a very similar set-up comprising of a ballistic 1–D QW with two short, but finite superconducting patches as shown in Fig. 1. Here, $\Delta^{(i)}$ is the pair potential on the two patches ($i$ refers to the index of the strip). In addition, we have extra top gates at the barrier region which allow for modulation of the local Fermi energy of the electrons in the patch region. Following [11], we use the Bogoliubov–de Gennes (BdG) equation [12,13] to calculate the transmission probability, $T^{(i)}_{ee}$ and the AR probability, $R^{(i)}_{eh}$, where $i$ is the barrier index. The space dependence of the order parameter (which also acts as the scattering potential) for the incident electron can be expressed as

$$V(x) = \Delta^{(i)} e^{i\phi_1} \Theta(x) \Theta(-x + a) + \Delta^{(i)} e^{i\phi_2} \Theta(x - (a + L)) \Theta(-x + (2a + L)) \quad (1)$$

where, $a$ is the width of the SB and $L$ is the distance between two barriers. Solving the BdG equation in the normal and the superconducting regions and matching the solution at $x = 0$ and $x = a$, we get

$$t^{(i)}_{ee} = \frac{e^{i\phi_1} a(u^1_+ - u^2_+)}{u^2_+ - u^2_+ e^{i(k^+ - k^-)a}}; \quad t^{(i)}_{hh} = \frac{e^{-i\phi_2} a(u^2_+ - u^2_+)}{u^2_+ - u^2_+ e^{i(k^+ - k^-)a}}$$

$$r^{(i)}_{eh} = \frac{u^2_+ e^{-i\phi_1} (1 - e^{i(k^+ - k^-)a})}{u^2_+ - u^2_+ e^{i(k^+ - k^-)a}}; \quad r^{(i)}_{he} = \frac{u^2_+ e^{i\phi_1} (1 - e^{i(k^+ - k^-)a})}{u^2_+ - u^2_+ e^{i(k^+ - k^-)a}} \quad (2)$$

![Fig. 1: A 1–D QW (e.g., carbon nanotube) connected to reservoirs, 1 and 2. The two thin strips on the QW depict layers of superconducting material deposited on top of the QW. The superconducting strips are connected to contacts, 3 and 4. The profile of the scattering potential is also shown. $V_g$ is the difference between the Fermi energy of electrons in the QW away from the barrier and at the barrier.](image-url)
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where \( \hbar k \pm = \sqrt{2m(E_F \pm (E^2 - \Delta^2 i))/2} \), \( u_{\pm} = \frac{1}{2}[(1 \pm (1 - (\Delta^2/E^2)^{1/2}))^{1/2}] \). Here, \( t_{ee}^{(1)}, t_{ee}^{(2)}, r_{eh}^{(1)}, r_{eh}^{(2)} \) are the transmission and AR amplitudes. \( m \) is the effective mass of the electron in the wire, \( E_F \) is the Fermi energy for the electrons in the superconducting region, and \( E \) is the Fermi energy of electrons in the normal region of the QW, measured with respect to \( E_F \). Hence the scattering matrix for the single SB problem for an incident electron or hole is given by

\[
S_e = \begin{bmatrix}
    r_{ee}^{(1)} & t_{ee}^{(1)} \\
    t_{ee}^{(1)} & r_{ee}^{(1)}
\end{bmatrix} \\
S_h = \begin{bmatrix}
    r_{eh}^{(1)} & t_{eh}^{(1)} \\
    t_{eh}^{(1)} & r_{eh}^{(1)}
\end{bmatrix}
\]  

Using the \( S \)-matrix given by Eq. 3, it is straightforward to obtain the effective \( S \)-matrix for the double barrier system [14]. The net transmission and net AR through the double barrier are

\[
T_{cc} = \frac{t_{ee}^{(1)} t_{ee}^{(2)} e^{i\phi} - L}{1 - r_{eh}^{(1)} r_{he}^{(1)} e^{i(\phi_x - \phi_y)} L} \\
R_{eh} = r_{eh}^{(1)} + t_{eh}^{(1)} \frac{r_{he}^{(1)} e^{i(\phi_x - \phi_y)} L}{1 - r_{eh}^{(1)} r_{he}^{(1)} e^{i(\phi_x - \phi_y)} L}
\]  

Note that the zero bias limit is the limit of perfect AR, and hence there will be no transmission. In our set-up there is top gate at each barrier resulting in Fermi-energy at the barrier region, different than Fermi energy in the wire. This leads to finite transmission even at zero bias. In order to obtain quantization of pumped charge, we choose to operate in the sub-gap regime i.e., \( E < \Delta \) (see Fig. 1). In this regime, \( |T_{cc}|^2 \) has sharp resonances at discrete values of \( E/\Delta \) for a given value of \( \phi_1 - \phi_2 \) [11]. These resonances result from multiple AR of electron to hole and vice-versa inside the double barrier.

Renormalisation group scheme and pumping formula. We include the proximity induced effects due to SB using an RG approach developed very recently [15] for the case of 1-D normal metal—superconductor—normal metal (NSN) junction. As we are only interested in the coherent regime (\( L_T >> L \), \( L_T \) is the thermal length), we can effectively treat the superconducting double barrier system (NSNSN junction) as a single barrier (NSN junction) as far as RG is concerned (see Fig. 2). The effective two-channel

\[
Q = \frac{e}{2\pi} \int_0^\tau dt \text{Im} \left[ -\frac{\partial S_{11}}{\partial V_1} S_{11}^* \dot{V}_1 + \frac{\partial S_{12}}{\partial V_1} S_{12}^* \dot{V}_1 - \frac{\partial S_{11}}{\partial V_2} S_{11}^* \dot{V}_2 + \frac{\partial S_{12}}{\partial V_2} S_{12}^* \dot{V}_2 \right] 
\]

where \( \tau \) is the pumping period and \( S_{ij} \) are the elements of the \( S \)-matrix. Note the negative sign in the above expression, which results from the fact that \( S_{11}^* \) and \( S_{22}^* \) correspond to conversion of an electron into a hole. Thus, the pumped charge is directly related to the amplitudes and phases that appear in the \( S \)-matrix. Inserting Eq. 5 in Eq. 6,

\[
Q = \frac{e}{2\pi} \int_0^\tau \left[ \dot{\theta} - G(t)(\dot{\theta} + \dot{\phi}) \right] dt 
\]

Here \( G(t) = |t_{ee}(t)|^2 \) is the time frozen two terminal linear conductance between the two terminals (labelled by 1, 2 in Fig. 1), in units of \( 2e^2/h \). The first term on the RHS in Eq. 7 is clearly quantized since \( e^{i\theta} \) returns to itself at the end of one cycle. So the only possible change in \( \theta \) in a period can be in integral multiples of \( 2\pi \) i.e., \( \theta(\tau) \rightarrow \theta(0) + 2\pi n \) (\( n \) integer). The second term is the ‘dissipative’ term which prevents the perfect quantization. The second term is directly proportional to the two terminal Landauer—Buttiker conductance for the system on the pumping contour. The relative sign between \( \dot{\theta} \) and \( \dot{\phi} \) in the expression for pumped charge in Eq. 7 originates from the AR process, which converts an electron to a hole. This is in contrast to what we have for the normal double barrier problem [3]. We include the effects of inter-electron interactions in the \( S \)-matrix using the RG method introduced recently for the case of NSN junction [15] in the context of 1-D QW. For a reflection-less junction, the basic idea of the method is as follows. The presence of a superconductor induces a finite yet weak pair potential in the QW resulting in scattering of incoming electrons to outgoing holes (Andreev processes) in the wire, away from the junction. Hence by calculating the total AR amplitude, due to scattering from the NSN junction and the (weak) pair potential in the wire perturbatively

\[
S = \begin{bmatrix}
    r_{ee} e^{i\theta} & t_{ee} e^{i\phi} \\
    t_{ee} e^{i\phi} & r_{ee} e^{i\theta'}
\end{bmatrix}
\]  

Fig. 2: Illustration of the fact that a double barrier can be replaced by a single barrier in the limit, \( L_T > L \) for the purpose of RG flow.
Fig. 3: Contours of transmission probability, $T_{\text{ee}}$, in the $\Delta_1 - \Delta_2$ plane at two different values of length scale, $L_P = 1$ and $L_P = 10$, at which the RG is cut-off for values of $V(0) = 0.1$ and $V(2k_F) = 0.1$. The red ellipse $C_1$ represents the pumping contour.

Fig. 4: Pumped charge $Q$, for pumping in $\Delta_1 - \Delta_2$ plane is shown in the figure as a function of length scale $L_P$ at which the RG is cut-off for three values of $V(0)$ and $V(2k_F)$.

In interaction strength and followed by “poor-man’s scaling” approach, we obtain the RG equation for the elements of the effective $S$-matrix of the SDB structure in the coherent regime ($L_T > L$). To experimentally observe the effect of the RG flow of $S$-matrix elements, one can start to cool the system and then compare the measured pumped charge obtained at different temperatures. Without loss of generality, we can choose the point of operation for cooling to be anywhere on the pumping contour. Hence, so as to avoid unnecessary complications arising due to the RG flow of phases associated with $S$-matrix elements ($\theta, \theta', \phi$), we choose the barriers to be symmetric at the operating point. This symmetry leads to vanishing of the RG flow of the phase hence making the observation of effects due to proximity on the pumped charge unambiguous and clearly visible. The RG flow of the normal transmission (and AR) amplitudes and phases are [15]

$$\frac{dt_{ee}}{dl} = \alpha'|t_{ee}|(1 - |t_{ee}|^2) \quad \text{and} \quad \frac{d\phi}{dl} = 0$$

$$\frac{dr_{eh}}{dl} = -\frac{\alpha'}{2}|r_{eh}|(|1 - |r_{eh}|^2|t_{ee}|^2\cos 2(\phi - \theta))\frac{d\phi}{dl} = \frac{\alpha'}{2}|t_{ee}|^2\sin 2(\phi - \theta)$$

(8)

Here we have considered the fully symmetric case, i.e. $\theta = \theta'$. Unitarity of the $S$-matrix in Eq. 5 implies that $\phi - \theta = \pi/2 + 2n\pi (n \rightarrow \text{integer})$. This simplifies the equations for RG flow for the reflection amplitude and phase,

$$\frac{dr_{eh}}{dl} = -\alpha'|r_{eh}|(1 - |r_{eh}|^2) \quad \text{and} \quad \frac{d\theta}{dl} = 0$$

(9)

Here, $\alpha' = (g_2 + g_1)/2\pi \hbar v_F$ where the $g_1, g_2$ are the running coupling constants whose bare values are set by $g_1(L_P = d) = V(2k_F)$ and $g_2(L_P = d) = V(0)$; $V(x)$ being the inter-electron interaction potential. The entries of $S$-matrix therefore become functions of the length scale $l$, or equivalently the physical length scale $L_P$, where $l = \ln(L_P/d)$ and $d$ is the short-distance cut-off (usually taken to be the average inter-particle spacing). The RG flow can also be considered to be a flow in the temperature since the length scale $L_P$ can be converted to a temperature using the thermal length $L_T = \hbar v_F/(k_B T)$. Hence, the RG flow has to be cut-off by either $L_T$, or the system size $L_S$, whichever is smaller [16]. We now integrate the RG equation for $t_{ee}$ complimented by the RG flow
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Fig. 5: The plot shows the variation of the AR phase $\phi$ with time $t$, along the pumping contour $C_1$ in the plane of $\Delta_1 - \Delta_2$ and the inset shows the variation of the same along the pumping contour $C_2$ in the plane of $\phi_1 - \phi_2$.

of $g_1$ and $g_2$ [15] to obtain the $L_P$ dependence of $T_{ee}$ as

$$T_{ee}(L_P) = \frac{T_{ee}^0}{R_{ch}^0 + T_{ee}^0} \left[ \left( 1 + 2\alpha_1 \ln \frac{L_P}{d} \right)^2 \left( \frac{d}{L_P} \right)^{(2\alpha_2 - \alpha_1)} \right]$$

(10)

Here $T_{ee}^0$ and $R_{ch}^0$ are the values of $T_{ee}$ and $R_{ch}$ at $L_P = L$ and $\alpha_1 = V(0)/2\pi \hbar v_F$ and $\alpha_2 = V(2k_F)/2\pi \hbar v_F$. There are two points worth mentioning here: (a) the transmission increases with increasing $L_P$ which is a consequence of the fact that the proximity effect due to superconductor induces an effective attractive interaction between the electrons, hence rendering the (Andreev) back-scattering an irrelevant operator, and (b) the expression for $T_{ee}(L_P)$ is not in the form of a pure power law even at $T_{ee}^0 \to 0$ limit, as is expected from Luttinger Liquid physics because of the RG flow of the $g_1, g_2$ parameters. Also, it is important to note that we take the short-distance cut-off $d$ to be the distance between the two barriers ($L$) since this is the length scale at which we glued the two barriers to a single barrier as far as RG is concerned. Using this, we can obtain the scaling behavior of the pumped charge ($Q$) as a function of the length scale $L_P$ (or the temperature $T$). In terms of the Landauer–Buttiker conductance, $G_0 = (2e^2/h)|t_{ee}|^2$, using Eq. 10, we obtain the pumped charge as

$$Q = Q_{int} - \frac{1}{4\pi e} \left( \frac{T}{v_F} \right)^{(2\alpha_2 - \alpha_1)} \int_0^\infty dt \ I(t)$$

(11)

Here $\delta = \theta + \phi$ and as earlier, $G_0$ is expressed in unit of $(2e^2/h)$. $Q_{int}$ is the integer contribution of the first term in Eq. 7.

Results and Discussions –

1. Pumping in the $\Delta_1 - \Delta_2$ plane : Here the pumped charge is obtained by periodically varying the top gate voltage which controls the Fermi energy of the electrons in the superconducting region. Hence it amounts to varying $E/\Delta$ for the two barriers periodically. Just like the double barrier problem, in this case too we observe resonant transmission of electrons at discrete values of $E/\Delta$ for fixed values of $\phi_1$ and $\phi_2$. These discrete values correspond to the existence of quasi-bound states formed inside the double barrier which are quite different from their double barrier counterpart as they are produced due to superposition of both electron and hole states and not just any one of them. In Fig. 3 (left panel), we see sharp resonance in transmission ($T_{ee}$) in the plane of $\Delta_1 - \Delta_2$ for $L = 1$. We employ the solutions to the RG equations (Eq. 8) to obtain the renormalized surface of transmission in the plane of $\Delta_1 - \Delta_2$ for a value of $L = 10$, this is shown in Fig. 3 (right panel). Note that the RG flow is such that the transmission increases in the entire $\Delta_1 - \Delta_2$ plane, hence reducing the sharpness of resonance and resulting in an increase of transmission (conductance) on the pumping contour $C_1$ giving rise to reduction in the pumped charge from its quantized value (see Fig. 4). From Fig. 5, we notice that the AR phase $\theta$ shows a total drop in its value by a factor of $2\pi$ during its time evolution along the contour $C_1$. This drop corresponds to the quantization of the topological part in the expression for pumped charge $Q$ (Eq. 7) to the value of e.

2. Pumping in the $\phi_1 - \phi_2$ plane : In contrast to the previous case, here we obtain two sharp lines of resonances for the transmission function in the $\phi_1 - \phi_2$ plane. Again we observe in Fig. 6 that the RG flow (Eq. 8) results in reduction of the sharpness of the resonance. We consider a pumping contour $C_2$ which encloses parts of both the resonance lines in the $\phi_1 - \phi_2$ plane. The intersection of the pumping contour $C_2$ with the lines of resonance results in vanishing of the topological part. This can be seen by observing the time-evolution of the AR phase along contour $C_2$ as shown in the inset of Fig. 5. In this case the drops are exactly compensated by corresponding rises in phase $\phi$ by same amount, leading to a net zero topological contribution to the pumped charge. Hence for small values of $L$ (see Fig. 7), the pumped charge is almost zero. This is because the topological part is identically zero while the dissipative part is non-zero but vanishingly small (due to the resonance being very sharp) as the conductance on most part of the contour is negligible. As we go to the larger $L$ values, the
pumped charge shows an interesting non-monotonic behavior, purely coming due to the variation of the dissipative part.

In conclusion, we show that pumping in the $\Delta_1 - \Delta_2$ plane is much more efficient as opposed to that in $\phi_1 - \phi_2$ plane. We also demonstrate that the quantization of the pumped charge in $\Delta_1 - \Delta_2$ plane is lost if we include correlations due to proximity effects in the 1-D QW.

Acknowledgements. – We thank Sumathi Rao for many stimulating and useful discussions and encouragement. We acknowledge use of the Bewoulf cluster at HRI. The work of SD was supported by the Feinberg Fellowship Programme at Weizmann, Israel.

Fig. 6: Contours of transmission probability, $T_{ee}$ in the $\phi_1 - \phi_2$ plane at two different values of length scale, $L_P = 1$ and $L_P = 10$ at which the RG is cut-off for values of $V(0) = 0.1$ and $V(2k_F) = 0.1$. The red circle $C_2$ represents the pumping contour.

Fig. 7: Pumped charge, $Q$ for pumping in $\phi_1 - \phi_2$ plane is shown in the figure as a function of length scale $L_P$ at which the RG is cut-off for three values of $V(0)$ and $V(2k_F)$.

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