Adiabatic Quantum Pumping at the Josephson Frequency

S. Russo,¹ J. Tobiska,² T. M. Klapwijk,¹ and A. F. Morpurgo¹

¹Kavli Institute of Nanoscience, Faculty of Applied Science, Delft University of Technology, Lorentzweg 1, 2628 CJ Delft, The Netherlands
²NTT Basic Research Laboratories, NTT Corporation, Atsugi, Japan

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We analyze theoretically adiabatic quantum pumping through a normal conductor that couples the normal regions of two superconductor–normal-metal–superconductor Josephson junctions. By using the phases of the superconducting order parameter in the superconducting contacts as pumping parameters, we demonstrate that a nonzero pumped charge can flow through the device. The device exploits the evolution of the superconducting phases due to the ac Josephson effect, and can therefore be operated at very high frequency, resulting in a pumped current as large as a few nanoamperes. The experimental relevance of our calculations is discussed.

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In a mesoscopic conductor in which electrons move phase coherently, a direct current can flow in response to a slowly varying periodic perturbation and in the absence of any applied bias. This phenomenon, known as quantum pumping, was first noticed by Thouless [1], who analyzed theoretically the response of an electron system to a “traveling” periodic potential $U(x - vt)$. The occurrence of quantum pumping requires that the periodic perturbation consists of at least two independent oscillating parameters $X_1(t)$ and $X_2(t)$, and that the trajectory representing the perturbation in the parameter space $(X_1, X_2)$ encloses a finite area [2,3]. When the cyclic perturbation is slower than the electron dwell time in the conductor, adiabatic pumping occurs and the system remains in thermodynamic equilibrium. In this case, the pumped charge can be expressed as a function of the scattering matrix and of its derivatives with respect to the pumping parameters $X_1$ and $X_2$ [3].

Attempts to investigate experimentally adiabatic quantum pumping have been made using electrostatically defined quantum dots in GaAs-based heterostructures [4]. In such a system, pumping is induced by oscillating voltages applied to the gate electrodes defining the dot. Although signatures of pumping signals may have been observed, the experiments are hindered by rectification effects originating from parasitic coupling of the ac signal applied to the gates [5]. Many other proposals of devices have been put forward in the literature, in which different physical quantities have been used as pumping parameters such as a time-varying magnetic field, the height of a tunnel barrier, etc. [6–8]. Often, however, these proposals do not consider the difficulties involved in the experimental realization.

Here we demonstrate theoretically the occurrence of pumping in a system of electrons and holes in a metallic conductor coupled to superconductors, where the pumping parameters are the phases of the superconducting order parameters in two different superconducting contacts. This system can operate at very high frequency without the need of feeding microwave radiation, simply by exploiting the evolution of the superconducting phases due to the ac Josephson effect. As a consequence, measurable pumped currents as large as a few nanoamperes, can be expected, while avoiding many of the spurious effects which affected previous experiments.

Figure 1(a) shows a schematic representation of the circuit that we propose. Two superconducting–normal-metal–superconducting (SNS) Josephson junctions are connected in parallel via a superconducting ring, and their N regions are additionally coupled by a normal-metal bridge. Andreev reflection [9] of electrons and holes takes place at each NS interface, resulting in a phase shift of the particle wave function, which at the Fermi energy is given by $\pm \chi_{1,2}$, the phase of the superconductor order parameter at the two different superconducting contacts (the sign is for reflection from hole to electron; the sign + for the reverse process). Hence, the total scattering matrix $(S_{tot})$ of the normal-metal bridge connecting the left and right reservoirs depends on the quantities $X_1 = e^{i\chi_1}$ and $X_2 = e^{i\chi_2}$. Here we calculate the signal due to pumping when $X_1$ and $X_2$ are used as pumping parameters.

The appealing aspect of the proposed device is the way in which the pumping parameters can be driven at high frequency, and their relative phase controlled. Specifically, the pumping parameters become time dependent when a constant voltage $V_{dc}$ is present across the SNS junctions (e.g., by biasing the junctions with a current higher than their critical current), since then $X_1(t)$ and $X_2(t) \approx e^{i (2eV_d/\hbar)t}$ owing to the ac Josephson effect [10]. Similarly to what happens in superconducting quantum interference devices (SQUIDS), the phase difference $\varphi$ between $X_1$ and $X_2$ can be easily controlled by applying a magnetic flux $\Phi$ to the superconducting loop, so that $X_2 = X_1 + \varphi$, with $\varphi = 2\pi \Phi/\Phi_0$ ($\Phi_0 = \hbar/2e$).

To demonstrate the occurrence of pumping, we model the system in the simplest possible way. We confine ourselves to the case of a fully phase coherent system at $T = 0$ K. The normal conductor is taken to consist of one channel supporting ballistic motion, and the separation...
between the Josephson junctions is $L$. The NS interfaces are all supposed to be perfectly transparent, i.e., the probability of Andreev reflection is unity [11]. The pumped current is equal to the charge pumped per cycle, multiplied by the pumping frequency. The calculation of the charge pumped per cycle follows the approach developed by Brouwer [3], modified to take into account the presence of the superconducting electrodes [8]. The relation between the charge $Q_{P,m}$ pumped in one of the two reservoirs (labeled by $m = 1, 2$) and the scattering matrix reads

$$ Q_{P,m} = e \int_0^\tau dt \left( \frac{dn_m}{dX_1} \frac{dx_1}{dt} + \frac{dn_m}{dX_2} \frac{dx_2}{dt} \right). $$

(1)

in which

$$ \frac{dn_m}{dX_{1,2}} = \frac{1}{2\pi} \sum_{ij} \gamma_{ij} \text{Im} \frac{\partial (S_{tot})_{ij}}{\partial X_{1,2}} (S_{tot})^*_{ij}, $$

(2)

where $\tau$ is the period of one pumping cycle ($\tau = \frac{2\pi}{\omega_j}$, with $\omega_j = \frac{2eV}{\hbar}$. In Eq. (2), the sum over $i$ extends to the electron and hole ($e$, $h$) channels in both leads. The sum over $j$ is performed over the electron and hole channels only in the lead connected to the reservoir $m$ for which the pumped charge is calculated. The function $\gamma_{ij}$ is equal to $+1$ when the element $S_{ij}$ of the scattering matrix corresponds to a process in which a current is pumped from lead $m$, $\gamma_{ij} = -1$ when a current is pumped into lead $m$. This difference in sign accounts for the fact that electrons and holes contribute oppositely to the pumped charge. Since the electron and hole contributions to the pumped charge could exactly compensate each other, it is not obvious a priori whether a net charge can be pumped.

The problem of computing the pumped charge is then reduced to the calculation of the total scattering matrix $S_{tot}$ of electrons and holes in the normal conductor bridge [see Fig. 1(b)], as a function of the parameters $X_1$ and $X_2$. The calculation is lengthy but conceptually straightforward (calculations were done using Mathematica™). We consider a perfectly symmetric configuration with two identical SNS junctions, which are also identically coupled to the normal-metal bridge. For each junction, the coupling is described by a “beam splitter” [12], whose scattering matrix $(S_{1,2})$ is assumed to be energy independent (i.e., it is the same for electrons and holes). We have chosen the simplest expression compatible with unitarity and time reversal symmetry. The expression reads

$$ S_{1,2} = \begin{pmatrix} a & \sqrt{\varepsilon} & b & \sqrt{\varepsilon} \\ \sqrt{\varepsilon} & -a & \sqrt{\varepsilon} & -b \\ b & \sqrt{\varepsilon} & a & \sqrt{\varepsilon} \\ \sqrt{\varepsilon} & -b & \sqrt{\varepsilon} & -a \end{pmatrix}, $$

(3)

where $\varepsilon$ varies between 0 and $1/2$ ($\varepsilon/2$ is the probability for an incoming particle to be deflected towards one of the superconductors). The amplitudes for backscattering $a$ and direct transmission $b$ across the beam splitter satisfy the relations $a^2 + b^2 + \varepsilon = 1$ and $\varepsilon/2ab = 1$, imposed by unitarity. For every fixed value of $\varepsilon$ two solutions, with $a > b$ and $b > a$, are possible and we considered both cases (for $\varepsilon = 1/2$ the two solutions coincide and $a = b = 1/2$).

We model the pair potential at the NS interface with a step function $\Delta = 0$ in $N$ and $\Delta = \Delta e^{i\chi}$ in $S$. Having assumed perfect transparency at the NS interfaces, the matrix describing Andreev reflection in the “vertical” branches of the circuit (see Fig. 1) depends only on the phase $\chi$ of the superconducting order parameter. It reads

$$ S_{AR} = \begin{pmatrix} 0 & r_{he} & 0 & 0 \\ r_{eh} & 0 & 0 & 0 \end{pmatrix} = \begin{pmatrix} 0 & -i e^{-i\chi} & 0 & 0 \\ -i e^{i\chi} & 0 & 0 & 0 \end{pmatrix}. $$

(4)

To calculate the total scattering matrix of the device we first calculate the scattering matrix associated to transport across only one SNS junction. We then consider the two SNS junction connected in series; i.e., we consider all the multiple reflection processes in the normal-metal bridge, taking into account the corresponding dynamical phases acquired by electrons and holes. The result is the scattering matrix $S_{tot}(X_1, X_2)$ that mixes the electron and hole channels in reflection and transmission.

Having determined $S_{tot}$, we obtain the pumped charge from Eqs. (1) and (2). As shown in Fig. 2 we find that, unless $k_F L = n\pi$ (with $n$ integer), the pumped charge is a nonzero, antisymmetric, and $2\pi$-periodic function of $\varphi$, as expected. The $2\pi$ periodicity in conjunction with the anti-
symmetry imply that the pumped charge has to vanish when \( \varphi = \pm \pi \). This is the case since for \( \varphi = \pi \) the trajectory in the space of the pumping parameters \( (X_1, \ X_2) \) does not enclose a finite area. In addition, the anti-symmetry of \( Q_P \) with respect to \( \varphi \) also implies that the sign of the pumped current changes when reversing the sign of the relative phase of the two superconducting junctions.

Figures 3(a)–3(c) summarize the outcome of our calculations for the different cases \( a > b \), \( a = b \), and \( a < b \). We first discuss the features of the results that are common to all three cases. We always find that the pumped charge does not depend on the separation \( W \) between the beam splitters and the superconducting leads (see Fig. 3). This is due to the phase conjugation [13] of electrons and holes at the Fermi energy, since the dynamical phase acquired by an electron propagating from the beam splitter to the superconducting interface is exactly compensated by the phase acquired by the Andreev reflected hole. In all cases the dependence of \( Q_P \) on \( L \) is periodic for all values of \( \varphi \) and \( \varepsilon \), with period given by \( k_F L = \pi \) (\( k_F \) is the Fermi wave vector). This implies that the pumped charge is sensitive to the geometry of the device on the scale of the Fermi wavelength \( \lambda_F \), indicating that charge pumping in the device considered here is a sample specific phenomenon. That this should be so is not obvious \textit{a priori}: owing to phase conjugation, one may have expected the pumped charge to show a component independent of the precise geometry of the device [14]. The magnitude of the calculated pumped charge strongly depends on \( \varepsilon \). The maximum pumped charge is approximately 0.1 electron per cycle, for \( \varepsilon = 1/2 \) \( (a = b) \). For small \( \varepsilon \), the magnitude of \( Q_P \) decreases with decreasing \( \varepsilon \) (and eventually vanishes for \( \varepsilon = 0 \)) both when \( a > b \) and \( b > a \). The dependence of \( Q_P \) on \( \varphi \) and \( k_F L \), however, is different in the two cases.

When \( a > b \) and for small values of \( \varepsilon \) (i.e., \( a \sim 1 \)) the bridge connecting the two SNS junctions is only weakly coupled to the reservoirs, because backscattering at the beam splitters is the dominant process [see Fig. 3(a)]. In this regime, sharply defined resonances appear in the conductance of the system when \( k_F L = n \pi \) (with \( n \) integer), due to the presence of quasibound states in the bridge connecting the two SNS junctions. When \( k_F L = n \pi \) the energy of a quasibound state aligns with the Fermi levels in the reservoirs. Interestingly, the pumped charge is also significantly different from zero only when \( k_F L \) is close to being a multiple of \( \pi \). This suggests a close link between pumping and the presence of resonances due to quasibound states in the system, as already noted by others in different contexts [7]. This link is further supported by observing that increasing \( \varepsilon \) from 0 to 1/2—corresponding to increasing the broadening of the quasibound states—results in a broader range of values of \( L \) for which charge pumping is observed [Fig. 3(b)].

In the case \( b > a \), the behavior of the pumped charge for small values of \( \varepsilon \) is qualitatively different [see Fig. 3(c)]. In this regime, the dominant process at the beam splitters is direct transmission. Therefore electrons and holes have only a small probability to be deflected from the normal bridge to the NS interfaces. However, if they are deflected, they perform many Andreev reflections in one of the SNS junctions before they can escape again to the normal-metal bridge. As a consequence, along the dominant trajectories responsible for pumping, electrons and holes have a large probability to acquire a phase \( e^{i\pi\varepsilon} \) (with different, and

![FIG. 2](image_url)  
**FIG. 2.** Pumped charge per cycle \( Q_P \) as a function of phase difference \( \varphi \) between the pumping parameters, for the case \( a = b = \varepsilon = 1/2 \) and different values of \( k_F L \).
even large, integer values of \( N \), rather than simply \( e^{i\chi} \). This causes the phase dependence of the pumping signal to be richer in harmonics and, consequently, to exhibit very strong deviations from a simple sine dependence, as seen from Fig. 3(c).

Having established the occurrence of adiabatic quantum pumping, we briefly discuss some of the advantages of the proposed device. The use of the ac Josephson effect to generate the time dependence of the pumping parameters should allow operation at frequencies of several hundreds GHz. In fact, with superconductors such as Nb, NbN, or NbTiN, values for the superconducting gap \( \Delta \) corresponding to frequencies in excess of 1 THz are possible, so that our superconducting pump can operate at a few hundreds GHz when the voltages applied across the SNS junctions is still sufficiently lower than \( \Delta \). At a Josephson frequency of 100 GHz, the pumped current can exceed 1 nA, which is easily measurable. Note that, since the pumping parameters are coupled to the electron-hole wave functions via Andreev reflection, the coupling will remain good at these high frequencies. In addition, the fact that no external microwave signals are required to drive the circuit implies that only a negligible high-frequency power will be irradiated, thus minimizing rectification effects known to cause problems in other systems [5]. We note, however, that in our superconducting pump a dissipative current mediated by multiple Andreev reflections flows between the Josephson junctions during device operation (i.e., when a bias \( V_{dc} \) is applied). The electrons carrying this dissipative current can reach the terminals used to measure pumping, thereby generating a signal that mimics the pumping signal. It will be important to analyze the behavior of the signal due to the dissipative current (i.e., magnitude dependence on pumping parameters, etc.) to understand how this signal can be distinguished from the actual pumping signal. This analysis, however, goes beyond the scope of this Letter.

For the practical realization of the proposed superconducting pump we suggest the use of a ballistic InAs-based two-dimensional electron gas as normal conductor. Present technology enables the reduction of the number of conducting channels to \( \approx 10 \) [15], which is important since the predicted effect is of the order of 1 channel. The use of InAs also enables the realization of the needed highly transparent contacts to superconductors [16]. Furthermore, in ballistic devices in which the distance between the two SNS junctions is \( L = 1 \mu m \), the typical propagation time in the device will be of the order of \( L/v_F = 10^{-12} \) s \((v_F \approx 10^6 \text{ m/s}) \) is typically realized in InAs heterostructures). This is 10 times faster than the period of an ac pumping signal oscillating at 100 GHz, ensuring that the dwell time of electrons is much shorter than the period of the ac pumping signal, as it is needed for the device to operate in the adiabatic regime.

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Note added in proof.—It has been suggested by P. Brouwer and by Y. Nazarov [17] that the dissipative current due to multiple Andreev reflections gives no additional contribution to the pumped current that we have calculated here.