Noise rectification by a superconducting loop with two weak links

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We consider a superconducting loop with two weak links that encloses a magnetic flux. The weak links are unequal and are treated as Josephson junctions with non-sinusoidal phase dependence. We devise a model that takes into account the fluctuation of the critical currents, due to the fluctuations of the order parameter in the weak links. These fluctuations are important near the onset of superconductivity; in this regime they may significantly weaken and eventually disconnect the superconducting loop. As a consequence of these fluctuations and of the resistive noise in the junctions, the average dc voltage does not vanish. Our model can be easily extended to provide a qualitative description of a recent experiment.

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I. INTRODUCTION

This study is a cross of two well established phenomena: Brownian motors and fluctuations near a phase transition. The offspring of this combination will be a circuit that can rectify thermal noise and sustain a dc voltage.

Brownian motors are systems in which thermal noise leads to the diffusion of some coordinate (typically, the position of a particle), whereas a time-dependent applied potential (typically, periodic in some direction of space) confines this coordinate and controls how far and when it is allowed to diffuse. If the applied potential is asymmetric, then the fluctuating coordinate drifts in a preferred direction and Brownian motion is rectified. Many reviews\cite{1,2,3,4,5} and experiments\cite{6} on Brownian motors are available.

Brownian motors are also called “flashing ratchets”, alluding to the fact that the confining potential pulsates in time. A related phenomenon is a “rocked ratchet”; in this case the confining potential is static, but an additional time-dependent applied potential is superposed to it. This addition distorts the confining potential in such a way that any of its minima can temporarily disappear and the confined object is transferred to the following minimum. A rocked ratchet may also operate without thermal noise.

Phase transitions are usually characterized by an “order parameter” that vanishes for one of the phases and has a non zero value for the other phase. Near a second order phase transition, the equilibrium size of the order parameter is small and can be comparable to that of its thermal fluctuations. The system we will consider is a superconducting loop near its transition temperature. Having a loop automatically provides the periodicity which is encountered in Brownian motors and the presence of permanent currents in a preferred direction might be thought of as a starting point for the required asymmetric potential.

The feature that in my view turns a superconductor into an interesting rectifying system is the presence of two competing currents that are governed by independent fields: the supercurrent is governed by the “order parameter”, whereas the normal current is governed by the electromagnetic field. Both undergo thermal fluctuations. However, the strength of the electromagnetic fluctuations increases with temperature, whereas fluctuations of the order parameter are most important near the critical temperature. This difference gives us the freedom to assign independent sizes to each of these fluctuations.

In the system we will study, electromagnetic fluctuations will play the role of “Brownian motion”, whereas fluctuations of the order parameter will cause the time variation of the confining potential.

The simplest superconducting component we can consider is a Josephson junction. A Josephson junction is actually a zero-dimensional superconducting wire: its state is completely described by the values of the order parameter and of the electrochemical potential at its extremes. Therefore, in order to prevent the physical ideas from being obscured by mathematical complexity, we will model our loop by a uniform superconducting wire interrupted by Josephson junctions. Modeling weak links by Josephson junctions will enable us to reduce the dynamics of the system to standard textbook procedures. Ideas for using Josephson junctions as rectifying ratchets have been studied in Refs.\cite{10,11,12}. Rectification of the motion of vortices has been considered in Ref.\cite{13}.

In the following section we define the system to be considered and the rules that govern its evolution. We also provide some motivation for the choices involved in the model and some heuristics as to why rectification is expected to occur. Section III is a report of the numeric results. In Sec. IV our model will be compared with available experiments. Finally, there is a summary of the achieved results.
The equilibrium value of the critical current will be denoted by $\alpha I_{c0}$ and, in order to reduce the number of parameters in the model, the resistance will be taken as $R_2 = R_1/\alpha$ and the capacitance as $C_2 = \alpha C$. $I_{c0}$, $R_1$ and $\Phi_0/2\pi e = h/2e$ will be taken as units; accordingly, the units of voltage, time, capacity and inductance are $I_{c0}R_1$, $h/2eI_{c0}R_1$, $h/2eI_{c0}R_1^2$ and $h/2eI_{c0}$, respectively. In these units the ac Josephson relation is

$$d\gamma_i/dt = V_i.$$  \hspace{1cm} (1)

Additional equations are obtained by equating the current around the loop, which is related to the magnetic flux encircled by the loop, with the current flowing across each junction:

$$I_{j1} + V_i/R_i + C_i dV_i/dt + I_{N1} = -(-1)^i(\varphi_x + \gamma_2 - \gamma_1)/L.$$  \hspace{1cm} (2)

$I_{j1}$ is the current that flows across the “pure” junction, $V_i/R_i$ the current that ideally flows through the resistor according to Ohm’s law, $C_i dV_i/dt$ is the current that goes into the capacitor, and $I_{N1}$ is due to thermal fluctuations; $\varphi_x$ is the magnetic flux induced by external sources multiplied by $2\pi/\Phi_0$ and $L$ is the self inductance of the loop.

$I_{N1}$ is the usual Johnson–Nyquist current. Let us choose a period of time $\Delta t$ which is large compared to the autocorrelation time, but can be treated as infinitesimal in macroscopic processes. Averaged over $\Delta t$, $I_{N1}$ has the form

$$I_{N1} = \eta g_1\sqrt{\Delta t}; \quad I_{N2} = \eta g_2/\alpha\Delta t,$$  \hspace{1cm} (3)

with $\eta^2 = 2k_B T/R_1$, where $k_B$ is Boltzmann’s constant, $T$ is the temperature, and $g_{1,2}$ are random numbers with normal distribution, zero average and unit variance. (This follows since $\langle (I_{N1})^2 \rangle = \int_0^{\Delta t} dt' \int_0^{\Delta t} dt'' (I_{N1}(t')I_{N1}(t'')) \approx 2\Delta t \int_0^{\infty} (I_{N1}(0)I_{N1}(s)) ds = 2k_B T \Delta t/R_1$)

**B. Heuristic Considerations and Additional Choices**

If we ignore $I_{N1}$, Eqs. (1) and (2) predict an evolution for $\gamma_i$ that is equivalent to viscous motion down a potential $(\varphi_x + \gamma_2 - \gamma_1)/2L + Z_1 + Z_2$, where $dI_j/d\gamma_i = I_{j1}$. In order to rectify Brownian motion, there should be some direction in the $\gamma_{1,2}$-space such that this potential is periodic and is not symmetric under reflection.

This symmetry breaking cannot be achieved if $I_{j1}$ and $I_{j2}$ are both sinusoidal functions of the phase difference, as is the case for tunnel junctions. Fortunately, for superconducting wires the current-phase dependence becomes sinusoidal only as a limit. In general, $I_{j1}$ has its maximum away from $\gamma_1 = \pi/2$. The deviation from sinusoidality can be quite large and $I_{j1}$ can even be multivalued. Various types of
nonsinusoidal current-phase relation were also encountered in junctions such as superconductor-ferromagnet-superconductor point contacts.\(^\text{15}\)

We may gain a better understanding of the present model by reviewing Ref. 14 in some detail. This reference considers a long and thin wire \(A\) with a region of it, denoted by \(B\), replaced by a different material. \(B\) is the “weak link”. Both materials are almost identical, the only difference between them being that \(B\) has a shorter mean free path. The weak link is characterized by two parameters. One of them, denoted by \(\Gamma\), is the ratio between Gorkov’s universal function of the impurity parameter for both materials. (Denoting the coherence lengths by \(\xi_{A,B}\), \(\Gamma = (\xi_B/\xi_A)^2\).) The second parameter is the length of \(B\), which is written as \(2d\xi_A\). The article provides explicit expressions that exactly solve the Ginzburg–Landau equations and permit to evaluate the current-phase relation for the wire. We are interested in temperature close to critical, so that we should consider \(d \ll 1\); on the other hand, any value \(0 < \Gamma < 1\) is physically admissible. The authors conclude that only in the regime \(d^2 \ll \Gamma \ll d\) the current-phase relation becomes sinusoidal. In order to appreciate the deviation from sinusoidality, Fig. 2 shows the current as a function of the phase difference for \(d = 0.001\) and several values of \(\Gamma\).

For simplicity, we shall keep only two harmonics and take

\[
I_{c1}(\gamma_i) = I_{c1}(\sin \gamma_i + \beta_i \sin 2\gamma_i),
\]

where \(I_{c1}\) and \(\beta_i\) are constants that characterize the junction \(i\). In order to present results that are clearly visible, most of our figures use the values \(\beta_1 = -\beta_2 = 0.7\), but let me emphasize that the numerical investigation covered the range \(0 \leq |\beta_i| \leq 1\) and the qualitative conclusions of the following section remain valid for small values of \(|\beta_i|\).

In Ref. 10 the second harmonic is brought in by connecting two junctions in series. It should be noted, though, that this method involves the implicit assumption that the same current flows through both “pure” junctions in series; this assumption is not acceptable in the context of this article, since the currents through both resistors fluctuate independently.

Using Eq. 11, the “potential” that describes the evolution of \(\gamma_i\) becomes \((\varphi_x + \gamma_2 - \gamma_1)^2/2L - I_{c1}(\cos \gamma_1 + \frac{1}{2}\beta_1 \cos 2\gamma_1) - I_{c2}(\cos \gamma_2 + \frac{1}{2}\beta_2 \cos 2\gamma_2)\). A contour graph of this potential is shown in Fig. 3 for a given set of parameters. The dark areas are minima, and the light areas are barriers between them. If the barriers are high, the values of \(\gamma_i\) will be confined within a minimum, but, if the barriers are lowered, the noise will cause these values to diffuse, either to the right or to the left. Since the barrier at the left is closer than that at the right, the probability of crossing it is larger than that of crossing the barrier at the right. Therefore, if the barriers are raised again after a suitable time, the values of \(\gamma_i\) may be pushed towards the consecutive minimum at the left, with a greater probability than that of being pushed to the right. In this way we may expect a non zero average drift of \(\gamma_1 + \gamma_2\); \(\gamma_1 - \gamma_2\) remains always close to \(\varphi_x\).

We require a mechanism to vary the asymmetric potential of Fig. 3 in time. This is provided by fluctuations of \(I_{c1}\) and \(I_{c2}\). The maximum current through a junction is proportional to \(|\psi|^2\), where \(\psi\) is the order parameter. The evolution of \(\psi\) can be described by the time-dependent Ginzburg–Landau theory with the addition of a Langevin “force”.\(^\text{18}\) The “standard” terms drive \(\psi\) towards equilibrium, whereas the Langevin term causes stochastic deviations from it. Since the Langevin term increases with the absolute temperature and the “standard” terms increase with the distance from the critical temperature, the influence of fluctuations is largest close to the critical temperature.

Due to the evolution of \(\psi\), the evolution of \(I_{c1}\) should also have the form of a random walk, together with relaxation to the equilibrium value with some characteristic time \(\tau_i\). The explicit implementation we took was

\[
I_{c1}(t + \Delta t) = \frac{\tau_1 [I_{c1}(t) + (\delta_0 r_0 + \delta_1 r_1)\sqrt{\Delta t}] + \Delta t I_{c0}}{\tau_1 + \Delta t}
\]

and, similarly, for \(I_{c2}\) the index 1 is replaced by 2, \(I_{c0}\) by \(\alpha I_{c0}\) and \(\delta_0\) by \(\delta_0\). Here \(\Delta t\) is the short period of time we have chosen, \(r_0\) and \(r_1\) are random numbers distributed between \(-1\) and 1 and the \(\delta_i\)’s denote the strength of the fluctuations. If \(I_{c1}\) or \(I_{c2}\) becomes negative, it is reset to 0. \(\delta_0\) describes synchronous fluctuations of both junctions, whereas \(\delta_1, \delta_2\) describe independent fluctuations. If the coherence length is larger than

![FIG. 2: Current through a Josephson junction as a function of the phase difference for the case considered in Ref. 14. Here the unit of current is \(we\xi_A H^*_c/\Phi_0\), where \(w\) is the cross section of the wire, \(\xi_A\) the coherence length in the strongly superconducting wire and \(H_c\) the thermodynamic transition field. The length of the weak link is \(2d\xi_A = 0.002\xi_A\) and the curves are for several values of \(\Gamma\), a parameter that defines the ratio between the strengths of superconductivity in both materials. The curves are marked with the number \(\log_{10}(\Gamma/d)\). Only the curve for \(\Gamma = 10^{-3}d\) is nearly sinusoidal. In order to make the curve visible, the current values for the case \(\Gamma = 10^{-3}d\) were multiplied by 400.](image-url)
The values of $\gamma_i$ behave as the coordinates of a particle in two dimensions (with anisotropic mass) which feels the potential in this graph, in addition to an anisotropic viscosity and a Langevin force. In this graph the darker areas represent lower potentials. This potential is periodic in $\gamma_1+\gamma_2$, with period $4\pi$. The parameters used to evaluate this graph are $L = I_{c1} = I_{c2} = 1, \beta_1 = -\beta_2 = 0.7, \varphi_2 = \pi/2$.

In order to reduce the number of parameters in the model we took $\tau = \alpha = \beta_1 = -\beta_2 = 0.7, \varphi_2 = \pi/2$. From here, it can be obtained that the average of the square of the change of $I_{c1}$ during the period $\Delta t$ will be $k_BT/\kappa \tau_1$. It then follows that $\tau_1 \delta_i^2 = 2k_BT/\kappa$. However, since $\kappa$ depends on the details of the junction, we will regard $\tau$ and the $\delta_i$’s as independent parameters.

The precise form of the fluctuations of $I_{c1}$ is not important. I have changed the distributions of the random numbers $r_i$ from uniform to Gaussian, to bimodal and to asymmetric (retaining the values of the average and the variance) and the difference in the results obtained from different distributions is not noticeably larger than the typical difference obtained for different runs with the same distribution. All that matters is how often and for how long the potential barriers in Fig. 3 are low compared to the noise level. The scattering in the results, though, seems to depend on the distribution used.

One might gain some intuition by comparing the orders of magnitude of the energies involved in the problem. Fluctuations in the electromagnetic field and in the order parameter are both of the order of $k_BT$; the energy required to break superconductivity ($\sim \kappa I_{c0}^2$) should be significantly, but not exceedingly, larger than $k_BT$. In this way, only when many random steps accumulate to diminish the order parameter, superconductivity is broken at a junction. For the parameters used here, $I_{c1}$ typically vanished in one out of $10^4$ steps.

III. RESULTS

The equations above determine the evolution of the phase differences $\gamma_{1,2}$ for given sets of model parameters. This evolution was followed, using Euler iterations, during $10^8$ steps. The average voltages are then obtained by dividing the change of $\gamma_{1,2}$ by the elapsed time. In order to cancel out possible biases in the random numbers, we evaluated four sets of phase differences, which were obtained by reversing the sign of either $I_{c1}$ or $I_{c2}$ in Eq. 8, and then took the average of the voltages obtained for the four sets. The values obtained for $V_1$ and $V_2$ were nearly the same, and the reported voltage values are their averages.

Figure 4 shows our main result, obtained for a given set of parameters. In spite of the scattering in the results, it can be safely concluded that the dc voltage does not vanish, that it is a function of the flux, that it vanishes for integer and half-integer values of $\Phi/\Phi_0$, and has maxima in between.

I have studied the dependence of the voltage on each of the parameters of the model. This study has been limited to variations of a single parameter each time, with the others fixed at their values as in Fig. 4. For $\delta_0$, $\tau$, and $\eta$, the voltage is insignificant below some threshold, then rises to a maximum and decreases slowly. This is expected, because for low values of these parameters the phase differences have no opportunities to jump between consecutive minima of the “potential”, whereas for high values there are too many opportunities and thermal equilibrium is reached. Figure 7 shows the voltage as a function of $\delta_0$. Qualitatively, the same shape is ob-
voltage is roughly proportional to \( -\beta + 0.25\beta^3 \) and, for \( \delta_1 \leq 0.03, \) to \( e^{-80\delta_1}. \)

In principle, the results should not depend on \( \Delta t, \) but in practice the steps cannot be too large if the derivatives of \( \gamma_i \) and \( V_i \) are approximated by ratios of discrete differences; \( \Delta t \) also cannot be too small, since then too many steps are required until \( I_{c1,2} \) become small a statistically significant number of times. In our calculations we took \( \Delta t = 0.01. \) In order to estimate the accuracy of our results, we repeated some calculations for \( \Delta t = 0.1. \) The voltages obtained were lower by about 20%.

An important difference between our model and other flashing ratchets is that here the variations of \( I_c \) in time are themselves due to spontaneous fluctuations, and it is not obvious who is the agent that invests the work required for rectification. Some models have been studied in which rectification is performed by fluctuations, but those are manifestly nonequilibrium fluctuations, whereas our heuristic considerations suggest that in the present case rectification is not subject to any detailed form of \( I_N \) or of the fluctuations of \( I_c, \) so that the question remains open.

**IV. APPLICATION TO EXISTENT EXPERIMENT**

Recent measurements on mesoscopic rings composed by segments of unequal widths, near the critical temperature, found a dc voltage with apparently the same flux-dependence as in our Fig. 4. Thus far we have not explained their results, since the voltages we have found are smaller than \( I_{c0}R_1 \) by four orders of magnitude, whereas the experimental results are smaller than the corresponding product by only one order of magnitude (provided that we identify \( R_1 \) with the resistance of a segment in its normal state and \( I_{c1} \) with the maximum supercurrent that can flow in it).

Nevertheless, bearing in mind the differences between the experiment and our model, the very existence of a non zero flux-dependent dc voltage in our case is already remarkable. The experiment differs from our model in two essential respects. The first difference is that their rings contain only two segments, without the strongly superconducting branches (see Fig. 1). Conceivably, by assuming that \( I_{j1} \) depends only on \( \gamma_j \) and by assuming that fluctuations affect the size of \( R_1 \) only (with the shape of the functions remaining invariant), essential features of the ring behavior are lost.

The second difference between the experiment and our model is that the experiment used an external ac current \( I_{c}(t) \) to destroy superconductivity. This feature is readily incorporated into our model. Assuming that both segments give equal contributions to the self inductance of the ring, all we have to do is add \( I_c(t)/2 \) at the right hand side of Eq. (2). By adding this term, our system becomes a rocked ratchet.

Figure 5 shows the dc voltages obtained for our model.
with an external current of the form \( I(t) = I_0 \sin(\omega t) \). Some of the parameters used to evaluate these results \((\alpha, C, L)\) were estimated from the available experimental data; the remaining parameters are arbitrary. The results bear close resemblance to those obtained for other periodically rocked ratchets\[12,22\] for rocking periods \(2\pi/\omega\) that are larger by two or more orders of magnitude than the characteristic time \(\hbar/2eI_0R_1\) of the circuit, we obtain the upper curve, independently of \(\omega\). However, for periods that are comparable to \(\hbar/2eI_0R_1\), we find that there is a nearly periodic pattern of rocking strengths for which the rectifying effect is present.

In the experiment, \(\hbar/2eI_0R_1 \sim 10^{-11}s\), whereas the shortest period of the applied ac current was \(10^{-6}s\). The relevant curve is therefore the upper one. Indeed, the experiment found that the frequency of the ac current has no influence. The upper curve in Fig. 6 is remarkably similar to Fig. 6 in Ref. 21, and this time the orders of magnitude coincide (for the voltages, and also for the currents).

The heuristic argument used in Fig. 3 to predict the existence of a dc voltage is not applicable now: in the presence of a large external current, the sign of the voltage is opposite to what that argument would predict.

Even in the presence of an external current, the noise and the supercurrent fluctuations play a central role. For \(I_0 = 4\) and the other parameters as in the upper curve of Fig. 6, setting \(\eta = \delta_0 = 0\) leads to a voltage decrease by one order of magnitude; for \(I_0 \lesssim 2.1\), the voltage vanishes when we set \(\eta = \delta_0 = 0\).

The experimental temperature is related to the parameters in our model through the values of \(I_0\). The typical deviation of \(I_c\) from its equilibrium value is of the order of \(\delta_0 \tau^{1/2} \sim (k_B T / \kappa)^{1/2}\). A significant response is obtained when \(I_0\) is larger than—but comparable to—this typical deviation, i.e., for temperatures slightly below the onset of superconductivity, as in this experiment.

Let us finally discuss two additional experiments in which nonuniform superconducting loops lead to a flux dependent dc voltage. Long ago de Waele et al.\[23\] measured dc voltage on a double point contact SQUID. The asymmetry was not in the weak links, as in the present article, but rather in the thick parts of the loop: one part was a niobium foil and the other a tin needle. No controlled driving ac current was supplied to the circuit; apparently, the role of the ac current was substituted by existing electromagnetic radiation in the lab. When the circuit was appropriately shielded, the dc voltage was no longer measurable. As in the cases considered here, the highest dc voltage was found when the temperature was slightly lower than the critical temperature of the weaker superconducting material (Sn). The highest dc voltage was comparable to the product of the maximal current through a point contact times its normal resistance.

Weiss et al.\[24\] measured rectification by an YBCO SQUID for a wide range of applied frequencies (ac and rf). The asymmetry was mainly due to different maximal currents of the junctions; this asymmetry leads to a rocking ratchet. Their ac results, shown in their Fig. 5(a) are remarkably similar to the upper curve in our Fig. 6 including orders of magnitude. In later experiments\[25\] the same group studied the influence of temperature and repeated the experiment for a niobium circuit. However, all the temperatures were considerably below the critical temperature, so that only variations in our parameter \(\eta\) were significant, while the fluctuations of \(I_c\) had no noticeable influence. In addition, they used tunnel junctions, so that our \(\beta\) is also expected to vanish.

In view of the differences among the reviewed experimental loops and our “zero-dimensional” model, their similar behavior is remarkable and suggests that the features encountered in our results will be present in any sort of nonuniform superconducting loop near its transition temperature.

**V. SUMMARY**

Superconductors near their transition temperature may be regarded as a special class of thermal ratchets, in the sense that the fluctuations of its order parameter can act as a flashing asymmetric “potential” and thus rectify the Johnson noise. We have found a simple model that qualitatively reproduces the experimental results of Ref. 21. For applied frequencies that are comparable to the “natural” frequency of the circuit, a complex pattern is predicted. In the absence of external current, the voltage decreases by about three orders of magnitude and may change sign, but our simulations and our heuristic considerations indicate that it does not vanish.
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