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Information entropic superconducting microcooler

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We consider a design for a cyclic microrefrigerator using a superconducting flux qubit. Adiabatic modulation of the flux combined with thermalization can be used to transfer energy from a lower temperature normal metal thin film resistor to another one at higher temperature. The frequency selectivity of photonic heat conduction is achieved by including the hot resistor as part of a high frequency LC resonator and the cold one as part of a low-frequency oscillator while keeping both circuits in the underdamped regime. We discuss the performance of the device in an experimentally realistic setting. This device illustrates the complementarity of information and thermodynamic entropy as the erasure of the quantum bit directly relates to the cooling of the resistor.

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I. INTRODUCTION

For the purpose of quantum computing, the coherence properties of superconducting quantum bits (qubits) should be optimized by decoupling them from all noise sources as well as possible. However, many interesting experiments can also be envisioned when the decoupling is far from perfect. One such experiment closely related to coherence optimization is using a qubit as a spectrometer1,3,2,4 for the environment of the device in an experimentally realistic setting. This device illustrates the complementarity of information and thermodynamic entropy as the erasure of the quantum bit directly relates to the cooling of the resistor.

II. FLUX-QUBIT COOLER AND THE THERMODYNAMIC CYCLE

We study a flux qubit coupled inductively to two different loops, as shown in Fig. 1(a). In loop \( j \) (\( j=1,2 \)), we have a resistor \( R_j \) in series with an inductor \( L_j \) and a capacitor \( C_j \). These form two damped harmonic oscillators. The resistors are in general at different temperatures \( T_1 \) and \( T_2 \). The coupling of the qubit to the admittances of the two loops, \( Y_1 \) and \( Y_2 \), is assumed to be sufficiently large to dominate the relaxation of the qubit. This assumption can be easily validated by, e.g., increasing the mutual inductance. The flux qubit is an otherwise superconducting loop except for three or four Josephson junctions with suitably chosen parameters. In particular, one of the junctions is made smaller than the others to form a two-level system. When biased close to half of the flux quantum \( \Phi_0 = h/2e \), the qubit can be described (in persistent current basis) by the Hamiltonian

\[
\frac{\hbar}{2} = -\frac{1}{2}(\Delta \sigma_x + \varepsilon \sigma_z),
\]

(1)

where \( \sigma_x \) and \( \sigma_z \) are Pauli matrices, \( \hbar \varepsilon = 2I_p(\Phi - \Phi_0/2) \) is the flux-tunable energy bias, and \( \Phi \) is the controllable flux threading the qubit loop. Away from \( \Phi = \Phi_0/2 \), the energy eigenstates have the persistent currents \( \pm I_p \) circulating in the loop. The tunneling energy \( \hbar \Delta \) results in an anticrossing at \( \Phi = \Phi_0/2 \) and there the expectation value of current of the energy eigenstates is zero. The transition angular frequency of the qubit is \( \omega = \sqrt{\varepsilon^2 + 4\Delta^2} \).

Consider the ideal cycle shown in Figs. 1(b) and 1(c) where the bias of the flux qubit is swept slowly (at a fre-
quency $f$ slower than $\Delta/2\pi$) between two extreme values $\epsilon_1$ and $\epsilon_2$ corresponding to two different energy level separations $h\omega_1$ and $h\omega_2$. Let us further assume that $\omega_1 \approx \omega_{iC}$ and $Q_1 \gg 1$, where $\omega_{iC} = (\omega_1^2 + \omega_2^2)/\omega_1^2$ and $Q_1 = (\omega L_{Cj}/R_j$. This choice guarantees that the qubit mainly couples to resistor $R_1$ ($R_2$) at bias point 1 (2). We emphasize that although both resistors are considered baths from the point of view of qubit relaxation, only resistor 2 is strictly speaking assumed to be a heat reservoir, whereas resistor 1 has to be “small” in order to cool down. The cooling cycle consists of steps O, P, Q, and R. First, in step O, the qubit has the angular frequency $\omega_2$ and is allowed to thermalize. Because of the bandwidth limitations imposed by the reactive elements, the qubit tends to thermalize with resistor $R_2$ to temperature $T_2$. However, ideally, the qubit splitting is large enough such that the thermal population is small. In the next step P, the flux bias is adiabatically changed to point 1 such that the level populations do not change but the energy eigenstates do. The sweep is assumed to be, however, faster than relaxation. In point 1, the angular frequency is reduced to $\omega_1$. Because the level populations and therefore the Boltzmann factors do not change, i.e., $h\omega_1/k_BT_2 = h\omega_2/k_BT_2$, the qubit must now be at a lower temperature $T_2$ given by $T_2 = T_2\omega_1/\omega_2$ in order to compensate for the change of the qubit splitting. Note that the quantum mechanical adiabaticity implies also thermodynamic adiabaticity: while the energy eigenstates change, the quantum mechanical adiabaticity implies also thermodynamic adiabaticity: because the energy eigenstates change, the level populations (and also the entropy) do not. In step Q, the qubit and resistor 1 are allowed to thermalize to temperature close to $T_1$, which results in heating of the qubit and in cooling of resistor 1 if $T_2 < T_1$. At this point, the ideally pure quantum state of the qubit gets erased and the stored information is lost. The entropy of the qubits increases, but locally the entropy of resistor 1 decreases such that one can say that some information is “stored” in the resistor as it cools but naturally with some loss. Finally, in step R, the qubit is adiabatically shifted back to frequency $\omega_2$, which results in heating of the qubit to the effective temperature $T_1 = T_2\omega_1/\omega_2$, which is assumed to be higher than $T_2$. The excess energy is dumped to admittance 2 when the cycle starts again from the beginning. Note that due to the condition $T_2 < T_1$, resistor 1 can never be cooled below $T_2\omega_1/\omega_2$. Since there is no isothermal stage in the above cycle, the present device is not even in principle a Carnot cooler but rather an Otto-type device.

The density matrix of the qubit with the transition angular frequency $\omega$ at temperature $T [\beta = (k_B T)^{-1}]$ is given by

$$\rho_{eq}(\beta, \epsilon) = \frac{1}{2} \left[ I + \left( \frac{\Delta}{\omega} \sigma_x + \frac{\epsilon}{\omega} \sigma_z \right) \tanh \left( \frac{\beta h \omega}{2} \right) \right],$$

where $I$ is the identity matrix. Using this, the cooling power and the efficiency of the ideal cycle in Fig. 1(c) can be easily calculated. It is given by the area of the shaded region in the entropy-temperature plane below points P and Q. In principle, one could solve for the effective temperature of the qubit along the line between points P and Q as a function of entropy given by $S = -k_B \ln \rho$. Alternatively, we can simply note that the expectation value of the energy stored in the qubit in point P is $E_P = \text{Tr} \{ \rho_{eq} \beta_2 \omega_2 / \omega_1, \epsilon \} H_1 \}$. While after relaxation, we have $E_Q = \text{Tr} \{ \rho_{eq} \beta_1, \epsilon \} H_1 \}$. where $H_1 \approx H(\epsilon_1)$ is the Hamiltonian at point 1. We thus get for the ideal cooling power $P = (E_Q - E_P) f$, i.e.,

$$P/f = -\frac{h \omega_1}{2} \tan \left( \frac{\beta_2 h \omega_1}{2} \right) + \frac{h \omega_1}{2} \tan \left( \frac{\beta_1 h \omega_1}{2} \right),$$

where $f$ is the pump frequency. The cooling power achieves the maximum value of $h \omega_1 f/2$ when the thermal population in step O (and P) is small and when the population in step Q is large, i.e., when $\beta_2 h \omega_2 \gg 1$ and $\beta_1 h \omega_1 \ll 1$. Naturally, a practical device has to be designed to always fulfill the first condition, in which case the smallest achievable temperature is on the order of $h \omega_1 / k_B$ below which the cooling power decreases rapidly. Another figure of merit is the ratio $\eta$ of the heat removed from resistor 1 divided by the heat added to resistor 2. It can be obtained as the ratio of the shaded area divided by the sum of the hatched area and the shaded area, i.e., $\eta = (E_Q - E_P) / (E_Q - E_0)$, where $E_0 = \text{Tr} \{ \rho \beta_2, \epsilon \} H_2 \}$ and $E_R = \text{Tr} \{ \rho \beta_1, \epsilon \} H_1 \}$. This simplifies neatly to $\eta = \omega_1 / \omega_2 < 1$, which is in harmony with the second law of thermodynamics.

### III. BLOCH EQUATION

For a more quantitative analysis, we have to consider the details of the relaxation rates due to the baths. The golden rule transition rates between the instantaneous $\epsilon$-dependent eigenstates due to resistor $j$ are given by

$$\Gamma_{i,i}^{j} = \frac{2\pi}{\hbar} |\langle 0; \epsilon | dH/d\Phi | 1; \epsilon \rangle|^2 M_j^2 S_j(\pm \omega_j)$$

$$= \frac{2\pi}{\hbar} \frac{\Delta^2}{\omega} M_j^2 S_j(\pm \omega_j),$$

where the positive sign corresponds to relaxation. The total relaxation rate is $\Gamma_{i}^{\text{tot}} = \Gamma_{i}^{1} + \Gamma_{i}^{2}$. Here, the unsymmetrized noise spectrum is given by

$$S_j(\omega) = \frac{1}{2\pi} \int_{-\infty}^{\infty} e^{-i\omega t} \langle \delta I_j(0) \delta I_j(t) \rangle dt$$

$$= \frac{1}{2\pi} \left[ \frac{2\hbar \omega \Re Y_j(\omega)}{1 - \exp(-\beta \hbar \omega)} \right],$$

where $\Re Y_j(\omega) = R_j^{-1} \left[ 1 + Q_j^2 \left( \frac{\omega}{\omega_{LCj}} - \frac{\omega}{\omega_j} \right)^2 \right]$ is the real part of admittance of circuit $j$. The total relaxation rate is thus

$$\Gamma_{i}^{\text{tot}} = \frac{2\hbar \Delta M_j^2 \coth \left( \frac{\beta \hbar \omega}{2} \right)}{R_j \hbar \omega \left[ 1 + Q_j^2 \left( \frac{\omega}{\omega_{LCj}} - \frac{\omega}{\omega_j} \right)^2 \right]}.$$

To model the behavior of the device, we utilize the Bloch master equation given in our case by
\[ \dot{M} = -B \times \dot{M} - \Gamma^1_{th} (\dot{M} \cdot \dot{M}) \dot{M} - \Gamma^2_{th} (\dot{M} \cdot \dot{M}) \dot{M} - \Gamma_3 \dot{M} \cdot \dot{M}, \]

where \( \dot{M} = \text{Tr}(\dot{\sigma} \rho) \) is the “magnetization” of the qubit and \( \ddot{B} = \Delta \ddot{x} + \varepsilon \ddot{z} \) is the fictitious magnetic field. Note, however, that the \( z \) component of \( \ddot{B} \) and \( \dot{M} \) do correspond to real magnetic field and magnetization, respectively. In Eq. (7), \( \dot{M} \) and \( \dot{M} \cdot \dot{M} \) are the components of the magnetization parallel and perpendicular to \( \dot{B} \), respectively. These are explicitly

\[ \dot{M} = \frac{1}{\omega^2} (\Delta M_\perp + \varepsilon M_\parallel) (\Delta \ddot{x} + \varepsilon \ddot{z}), \]

\[ \dot{M} = \frac{\varepsilon^2 M_\perp - \Delta \varepsilon M_\parallel + \Delta M_\parallel - \Delta \varepsilon M_\perp \ddot{z}}{\omega^2} + \frac{\Delta^2 M_\parallel - \Delta \varepsilon M_\perp \ddot{z}}{\omega^2}. \]

Here, \( \dot{M}_T \) stands for the \( e \)-dependent equilibrium magnetization of a qubit at temperature \( T \) given explicitly by

\[ \dot{M}_T = \left( \frac{\Delta \ddot{x} + \varepsilon \ddot{z}}{\omega} \right) \tanh \left( \frac{\beta B \omega}{2} \right), \]

and \( \Gamma_3 = (\Gamma^1_{th} + \Gamma^2_{th})/2 \) is the dephasing rate. We neglect pure dephasing due to the intentionally large dominating thermalization rate. Equation (7) describes the dynamics of a quantum two-level system coupled to two dissipative baths. The baths tend to relax the qubit toward instantaneous equilibrium with two competing rates. Equations of this type are often used in the stationary case, but they also yield very good predictions for strong driving as, for instance, in the case of the Landau-Zener interference.\(^{16}\) As is obvious from Eq. (7), the qubit actually tends to relax toward an effective \( e \)-dependent equilibrium magnetization \( \Gamma^1_{th} \dot{M}_T + \Gamma^2_{th} \dot{M}_T \) at the rate \( \Gamma^1_{th} + \Gamma^2_{th} \).

IV. SIMULATION RESULTS

To illustrate the practical potential of the device, we show in Fig. 2 the simulated cooling power with sinuosoidal driving of \( e(t) \) compared to the ideal case along with the actual loop in the entropy temperature plane. The heat flow \( P_{ej} \) from resistor \( j \) to the qubit is simply obtained by integrating the product of the thermalization rate and the energy deficit, i.e., \( P_{ej} = \int dV_{\text{th}} \text{Tr} \{ \rho \sigma, [\dot{H}(t), H] \} \text{Tr} \{ \rho \sigma \} \}. \) The density matrix of \( e(t) \) is solved numerically using the Bloch equation (system is followed over a few periods until it has converged to the limit cycle). We see that the actual simulated behavior does not significantly deviate at low \( f \) from the ideal behavior and that cooling powers on the order of femtowatts can be achieved with reasonable sample parameters. The oscillator behavior at high \( f \) is interpreted as the Landau-Zener interference.\(^{16,17}\) The ideal operation frequency for cooling would be the highest frequency where the cycle is still fairly adiabatic, i.e., about 1 GHz is the optimal frequency in the present example.

\[ P_1 = \int_0^\infty \frac{d\omega}{2\pi} [4\hbar \omega^2 M^2 \text{Re} Y_1(\omega) \text{Re} Y_2(\omega)(n_1(\omega) - n_1(\omega))], \]

where \( n_1(\omega) = [\exp(\beta \hbar \omega) - 1]^{-1} \) are the boson occupation factors and \( M \) is the mutual inductance between the loops. For detuned well underdamped resonators, the photonic heat conduction turns out to be quite negligible. For instance, for the values of Fig. 3 with \( M=5 \) \( \text{pH} \) and \( R_1=R_2=1 \) \( \Omega \), we get only \( P_1 < 2 \times 10^{-18} \) \( \text{W} \) even if \( T_1=0 \) \( \text{K} \) and \( T_2=300 \) \( \text{mK} \). Figure 2 illustrates the calculated equilibrium temperature versus operation frequency obtained numerically by finding the
While the capacitance values are similar to those in typical flux qubits.\textsuperscript{2} To satisfy the conditions of the above numerical example, we need quite large mutual inductances which, however, can be easily achieved using, e.g., kinetic inductance.\textsuperscript{21} The strong driving also requires rather large inductance between the microwave line and the qubit, which should not result in uncontrolled relaxation. For instance, $M_{\text{mw}} = 5 \, \text{pH}$ coupling (which is realistic) to the control line is acceptable as it would result in at most $3 \times 10^7 \, \text{s}^{-1}$ relaxation rate assuming a $50 \, \Omega$ environment at 0.3 K. This choice will not degrade the performance of the device significantly since driving is much faster. Yet, sufficiently strong driving corresponding to the example in Fig. 2 can be achieved with a modest $3 \, \mu \text{A}$ ac current. Fabrication process most likely will require three lithography steps.

\section*{V. CONCLUSIONS}

In conclusion, we have described a method of using a superconducting flux qubit driven strongly, yet adiabatically, at microwave frequency to cool an external metal resistor. Here we considered LC resonators to achieve the required frequency selectivity but a coplanar waveguide resonator or a mechanical oscillator could be used in principle, too. We demonstrated by a numerical example that it is possible to observe the associated temperature decrease experimentally. This effect is directly related to the loss of information and thus to the increase of entropy of the quantum bit.

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