A tunable rf SQUID manipulated as flux and phase qubits

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Received 16 October 2009
Accepted for publication 23 October 2009
Published 14 December 2009
Online at stacks.iop.org/PhysScr/T137/014011

Abstract

We report on two different manipulation procedures of a tunable rf superconducting quantum interference device (SQUID). First, we operate this system as a flux qubit, where the coherent evolution between the two flux states is induced by a rapid change of the energy potential, turning it from a double well into a single well. The measured coherent Larmor-like oscillation of the retrapping probability in one of the wells has a frequency ranging from 6 to 20 GHz, with a theoretically expected upper limit of 40 GHz. Furthermore, here we also report a manipulation of the same device as a phase qubit. In the phase regime, the manipulation of the energy states is realized by applying a resonant microwave drive. In spite of the conceptual difference between these two manipulation procedures, the measured decay times of Larmor oscillation and microwave-driven Rabi oscillation are rather similar. Due to the higher frequency of the Larmor oscillations, the microwave-free qubit manipulation allows for much faster coherent operations.

PACS numbers: 74.50.+r, 03.67.Lx, 85.25.Dq

(Some figures in this article are in colour only in the electronic version.)

1. Introduction

Superconducting qubits are promising systems for the realization of quantum computation. Their coherent evolution, entanglement, storage and transfer of quantum information as well as quantum non-demolition readout are just a few examples of what has already been achieved [1–4]. One of the major challenges that superconducting qubits are now facing is to increase their coherence time in order to reach a useful number of quantum operations [5]. However, there is an alternative possibility of achieving the same goal: one can try to increase the qubit operation speed to make qubit gates shorter, thus leading to less restrictive requirements for the coherence time. Nowadays, qubit operation frequencies exceeding 500 MHz are difficult to reach with phase qubits manipulated by microwave signals, since the limited anharmonicity of the qubit energy potential gives rise to a leakage into higher excited states.

In this paper, we report experiments with a tunable rf superconducting quantum interference device (SQUID) and demonstrate two different manipulation procedures. The system under test, a double SQUID, can be operated in two different regimes: (i) as a flux qubit manipulated via fast (dc) pulses of magnetic flux, and (ii) as a phase qubit in which the quantum state evolution is controlled by microwave pulses of chosen amplitude and phase. The former manipulation approach without using microwaves allows one to reach very high oscillation frequencies, while operating the same circuit as a phase qubit offers the possibility of verifying the obtained results using a well-defined and studied manipulation technique. We will see in the following that the coherence times measured using the two different manipulation schemes are rather similar, suggesting that decoherence acts in a similar way in both cases. This conclusion draws attention to common sources of decoherence (presumably, dielectric loss due to two-level fluctuators in the material used for chip fabrication [6, 7]) and emphasizes the relevance of increasing the number of qubit rotations within the coherence time by increasing the oscillation frequency of the system.
The paper is organized as follows. A detailed overview of the system under test, composed of the double SQUID and the readout dc SQUID, is given in section 2. Section 3 describes the manipulation of the double SQUID by deforming its energy potential using fast pulses of magnetic flux. The measured coherent oscillations are presented here together with their theoretical interpretation. Section 4 reports the manipulation of the same device as a phase qubit. Rabi oscillations measured via microwave pulses of variable length are reported. The dependence of the Rabi oscillation frequency on microwave power is analyzed by taking into account the population of higher excited states.

2. System and experimental setup

The circuit that we studied consists of a superconducting Nb loop of inductance \( L = 85 \, \text{pH} \) interrupted by a small dc SQUID of inductance \( l = 6 \, \text{pH} \). The name for this kind of device varies between different authors from ‘double SQUID’ [8] to ‘combined rf–dc SQUID’ [9] or ‘modified rf SQUID’ [10]. The two Josephson junctions embedded into the dc SQUID are nominally identical except for unavoidable asymmetries originating in the fabrication process. Each junction has a critical current of \( 8 \, \mu A \) and a self-capacitance of \( 0.4 \, \text{pF} \). The system is manipulated via two magnetic bias fluxes \( \Phi_1 \) and \( \Phi_2 \) applied to the large and small loops, respectively. The detailed characterization of the device taking into account non-identical Josephson junctions and non-negligible inductance \( l \) has been reported elsewhere [11]. Both loops are designed fully gradiometrically with the intent of decreasing both the noise induced by external uniform magnetic fields as well as cross talk between the two bias fluxes. A photograph of the double SQUID together with its schematic representation is presented in figure 1(a). The double SQUID is defined by the pale-white area delimited by white solid lines at the center of the picture. The two large holes of \( 100 \times 100 \, \mu \text{m}^2 \) define the gradiometric main loop. In the pale-white area highlighted by the white dashed ellipse, one can identify two much smaller holes (\( 10 \times 10 \, \mu \text{m}^2 \)), defining the gradiometric inner dc SQUID. The two Josephson junctions, of dimensions \( 3 \times 3 \, \mu \text{m}^2 \), are visible at the center. The two coils on the left and right sides are used to control the bias flux \( \Phi_2 \); the mutual inductance between them and the double SQUID is \( 2.6 \, \text{pH} \). The coil inducing the bias flux \( \Phi_1 \), visible on the lower-central part of the picture, is wrapped around one of the two small loops, to which it has a mutual inductance of \( 6.3 \, \text{pH} \). The placement of coils and holes reduces the cross talk between \( \Phi_1 \) and \( \Phi_2 \) lines to less than \( 1 \% \), as verified experimentally. Each of the two coils inside the large loops forms a part of a superconducting transformer connecting the system to an unshunted readout dc SQUID (one for each side). The circuit was made by Hypres Inc. (Elmsford, NY, USA) using standard Nb/AlO\(_x\)/Nb technology, with a critical current density of \( 100 \, \text{A cm}^{-2} \) and SiO\(_2\) as dielectric material for junction isolation.

The inductance of the small loop \( l \) is chosen to be much smaller than that of the main loop \( L \), so that the two-dimensional energy potential defining the system can be approximated by a one-dimensional function of the parameter \( \delta \), corresponding to the phase difference across the inner dc SQUID. Moreover, the inductance of the main loop \( L \) and the critical current of each junction \( I_0 \) are chosen such that the parameter \( \beta_L = 2\pi L I_0 / \Phi_0 \) ranges between 1 and \( 5\pi / 2 \), resulting in a double-well potential. The two bias fluxes \( \Phi_1 \) and \( \Phi_2 \) are used to manipulate the energy potential profile. Changes in \( \Phi_1 \) modify the symmetry (figure 1(b)) of the potential, whereas changes in \( \Phi_2 \) tune the height of the barrier between the two local minima (figure 1(c)).

Measurements were performed at the base temperature of a dilution fridge stabilized at 15 mK. The experimental setup is reported in figure 2. The two currents generating the fluxes \( \Phi_1 \) and \( \Phi_2 \) were supplied via coaxial cables with 10 dB attenuators anchored at the 1 K-pot stage of the refrigerator. The microwave signal, applied to the double SQUID via the coil \( \Phi_2 \), was supplied through an extra coaxial cable with 10 dB attenuation at the 1 K-pot stage and 10 dB at base temperature. For biasing the readout dc SQUID, we used superconducting wires, metal powder filters [12] at base temperature and current dividers at the 1 K-pot stage, while the lines to measure the voltage response were equipped with low-pass filters instead of current dividers at the 1 K-pot stage.

3. Microwave-free manipulation

The qubit manipulating scheme reported in this section is based on changing the double-well potential to a single-well...
The final circuit state is read out by applying a current as the ‘portal’ region. After the initial qubit state is prepared, the potential is and.

Here we perform the quantum manipulation: By applying manipulation of the double SQUID. The ground state |0\rangle is defined to be the zero reference level.

The main point of the above procedure is the non-adiabatic transition between the lowest two energy levels, which occurs during the potential transformation. In figure 4 the energy values, expressed as a frequency, of the first six eigenvalues for the double-SQUID, are reported versus the control flux \( \Phi_c \). On that graph the energy of the ground state |0\rangle is defined to be the zero reference level. The symmetric double-well situation and single-well situation are depicted, respectively, on the left and right sides of the graph.

For a perfectly symmetric double well, the lowest two energy levels are degenerate, with an energy splitting \( \Delta \) between them depending exponentially on the inverse of the barrier height. For the more realistic case of non-perfect symmetry, we express the asymmetry \( \epsilon \) as the energy difference between the two minima of the potential. For a large potential barrier, the level splitting is dominated by \( \epsilon \) and remains essentially constant during the initial drop of the barrier.

When the barrier is almost completely removed (for \( \epsilon \sim \Delta \), which is indicated in figure 4 as the ‘portal’ region), the level splitting starts to increase and saturates at the single-well oscillation frequency. This process corresponds to ‘half’ a Landau–Zener process, starting from the degeneracy point and arriving at the large energy gap condition (see figure 4). In this process, the system is prepared and remains initially in a 50% superposition of the two lowest energy eigenstates. The initial phase of this superposition depends on the chosen initial state of the deep double-well potential. For example, an initial left (|L\rangle) (right (|R\rangle)) flux state is transformed in a 50% superposition with positive (negative) sign. Although unwanted non-adiabatic transitions to upper levels (from the third onward) are possible, they can be suppressed owing to the gap existing between the first two levels and upper levels in the portal region. This requires an appropriate choice of the pulse rise/fall times so that the transition is performed non-adiabatically for the computational state but adiabatically for upper levels.
After crossing the portal region, the system is maintained for the pulse duration in the deep single-well potential. Here, the lowest two levels are equally populated and the relative phase between them evolves with a rate given by the level splitting. This level splitting is tunable by the flux pulse amplitude, and can be described by the formula

$$\Delta E = \hbar \omega_0(\Phi_c) \approx \frac{\hbar}{\sqrt{2LC}} \sqrt{1 - \beta(\Phi_c)},$$  (1)

where $\beta(\Phi_c) = \beta_L \cos(\pi \Phi_c / \Phi_0)$ is the modification of the parameter $\beta_L$ via the flux $\Phi_c$. During this time the system is weakly responsive to fluctuations in the flux $\Phi_c$ and so it is naturally protected against noise.\(^4\) The final phase between the states is determined by the duration of the flux pulse, i.e. by the time elapsed in the single-well condition. At the end of the pulse the flux returns to the initial condition and the system goes back to the two-well state. The portal region is crossed again, and the inverse of the previous process occurs: the relative phase between the two energy states is transformed to the amplitude of the left/right states. After the pulse, the system’s flux state becomes frozen once again, and it is read out by an unshunted dc SQUID magnetometer that is weakly inductively coupled to the qubit loop.

The described procedure was repeated many times (from 100 to 10,000) in order to determine the probability of one of the projected states for a chosen combination of parameters, such as pulse height and duration. This is then repeated for different pulse durations $\Delta \tau$ in order to record the coherent oscillations of the qubit state (figure 5). From the curves collected for different pulse amplitudes, we extracted the oscillation frequency and the decay time. In figure 6 we show the measured oscillation frequency (full dots) versus the pulse amplitude. These data are in very good agreement with the theoretical expectation given by (1) for the single-well oscillator frequency (solid line). The measured decay time is on the order of a nanosecond, independent of the frequency.

The described manipulation scheme only allows for $R_z$ rotations around the qubit’s Bloch sphere. The full qubit control also requires phase control achieved by $R_z$ rotations.

\(^4\) As an example, a fluctuation of $1\sigma_\Phi$ in the flux $\Phi_c$ is responsible for a change in the oscillation frequency of only 125 kHz, corresponding to a percentage variation of $6.6 \times 10^{-4}$ at 19 GHz.

4. Microwave-induced manipulation

In addition to the above microwave-free manipulation, we have also operated the system as a phase qubit using microwave driving. The operating procedure is the same as that reported in [17]. The energy potential profile is strongly tilted via the flux bias $\Phi_x$, making one of the two local minima shallow enough to contain only a small number of energy levels. The two computational states of the qubit are defined to be ground $|0\rangle$ and first excited states $|1\rangle$ in the shallow well. The first excited state is populated by resonant absorption of photons from the microwave field, while the complete manipulation on the Bloch sphere can be performed.

Figure 5. Probability of measuring the state $|L\rangle$ as a function of pulse duration. The coherent oscillation shown here has a frequency of 14 GHz and a coherence time of approximately 1.2 ns.

Figure 6. Measured oscillation frequencies versus amplitude of the short flux pulse (full dots). The solid curve is a numerical simulation using the measured parameters of the circuit.
The flux of the double SQUID is measured with the dc An adiabatic, but fast, dc pulse is sent to the bias coil The dashed line represents a linear power–frequency were fitted with two The solid fitting curve takes into account both a non– Figure 7. Rabi oscillation of the double SQUID manipulated as a phase qubit by applying microwave pulses at 19 GHz. The oscillation frequency changes from 540 MHz to 1.2 GHz by increasing the power of the microwave signal by 10 dB. via microwave pulses of defined duration and phase and dc flux pulses. The main difference between our device and the conventional rf SQUID phase qubits [18] is the possibility to tune in situ the Josephson energy of the device via the flux coil $\Phi_x$. This additional tuning parameter allows one to modify the anharmonicity of the energy potential in a slightly different way to what is done by the flux $\Phi_x$. It is thus possible to minimize the leakage to higher excited states in a strongly deformed potential (three to four energy levels in the shallow well) without changing the number of levels inside the well nor the escape probability of each state. The readout procedure is performed in two steps:

1. An adiabatic, but fast, dc pulse is sent to the bias coil $\Phi_x$; the potential is thus deformed and consequently the barrier separating the two wells is reduced. The amplitude of the pulse is calibrated such that the transition from the shallow to the deep well is triggered for the ground state $|0\rangle$ with a probability of approximately 10%. That ensures a complete escape of the system from the lowest excited state $|1\rangle$, leading to a theoretical visibility close to 90%.

2. The flux of the double SQUID is measured with the dc SQUID coupled inductively to it, using the procedure already described in section 3.

The sequence of preparation, manipulation and readout is repeated between 100 and 10 000 times, depending on the desired statistical error. An example of three Rabi oscillations obtained for different powers of the microwave driving field at 19 GHz is reported in figure 7. The increase of the oscillation frequency with microwave power is clearly seen. The topmost oscillation has a frequency of $f_{\text{Rabi}} = 540$ MHz. Oscillations at lower frequencies are difficult to measure since the decay time is very short, between 1.5 and 2.0 ns. On the lower graph, the oscillation frequency is approximately equal to $f_{\text{Rabi}} = 1.2$ GHz. Further increasing the power leads, as we will see, to an unwanted population of higher excited states.

The correlation of the Rabi oscillation frequency versus amplitude of the microwave driving field is reported in figure 8. The deviation from the linearity visible at higher powers is a clear indication of populating higher excited states [19]. We also note that the qubit has been driven slightly off-resonance since the frequency of Rabi oscillations does not reach zero for zero power.

The experimental data in figure 8 were fitted with two curves:

- The dashed line represents a linear power–frequency relation with the addition of a term taking into account the microwave pumping out of resonance. The equation has the form

$$f_{\text{Rabi}} = \sqrt{(f_{01} - f_{\text{MW}})^2 + (kA)^2},$$

where $A$ is the normalized amplitude of the microwave driving signal in mW$^{1/2}$ and $k$ is a constant. The microwave frequency $f_{\text{MW}} = 19$ GHz is offset from the resonance frequency $f_{01}$ corresponding to the transition between the lowest two levels. The points used for the fit refer to a normalized microwave amplitude ranging between 0.1 and 0.4 mW$^{1/2}$. From the fit we obtained a value of $k = (2919 \pm 105)$ MHz mW$^{-1/2}$ and $f_{01} - f_{\text{MW}} = (349 \pm 45)$ MHz.

- The solid fitting curve takes into account both a non-zero population of higher excited states due to large driving field and the manipulation red detuned from the resonance. The dependence of the Rabi frequency on the power of the microwave signal is described in [20] (equation (22)). The function that we used for the solid curve fit has the form

$$f_{\text{Rabi}} = \sqrt{(f_{01} - f_{\text{MW}})^2 + (kA (1 - \beta A^2))^2}.$$

As before, $f_{01} - f_{\text{MW}}$ is the frequency out of resonance and $k$ is the linear power–frequency correlation, while the new coefficient $\beta$ depends on the details of the system. The result of the fit, this time made through all measured points, gives the values of $k = (2977 \pm 104)$ MHz mW$^{-1/2}$, $f_{01} - f_{\text{MW}} = (350 \pm 42)$ MHz and $\beta = 0.411 \pm 0.060$ mW$^{-1}$. Within the noted errors, both fits yield the same values for the common parameters.
We measured the energy relaxation time $T_1$ of the double SQUID when it was operated as a phase qubit. The occupation probability of the first excited state was measured after a variable time between a resonant microwave $\pi$-pulse and the readout pulse. The measured probability decay was fitted exponentially yielding $T_1 = 1.37 \text{ ns}$ (figure 9).

5. Conclusion

We presented measurements on a double SQUID manipulated both as a double-well flux qubit and as a phase qubit. The device manipulation as a flux qubit was performed by modification of its energy potential profile via fast dc pulses, whereas the manipulation as a phase qubit required the use of microwave signals to induce Rabi oscillations between the ground and first excited states of a shallow well in a strongly deformed potential. The measured coherence time of the Larmor oscillations obtained in the flux regime is about 1 ns, of the same order as the relaxation time $T_1$ in the phase regime. Since best available phase qubits already display relaxation times of several 100 ns, obtained by using appropriate materials in the fabrication processes [6], we suppose that also the coherence time of the Larmor oscillations obtained without microwaves could strongly benefit from the same treatment. Such a possible improvement of the coherence by two orders of magnitude, together with the much higher oscillation frequency in the microwave-free Larmor mode, should in principle allow one to reach the ultimate goal of $10^4$ single-qubit gate operations within the coherence time that is needed for the implementation of quantum algorithms.

Acknowledgments

This work was supported by the Deutsche Forschungsgemeinschaft (DFG) and the EU project EuroSQIP.

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