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Quantum Superposition of Macroscopic Persistent-Current States

Caspar H. van der Wal, A. C. J. ter Haar, F. K. Wilhelm, R. N. Schouten, C. J. P. M. Harmans, T. P. Orlando, Seth Lloyd, J. E. Mooij

Microwave spectroscopy experiments have been performed on two quantum levels of a macroscopic superconducting loop with three Josephson junctions. Level repulsion of the ground state and first excited state is found where two classical persistent-current states with opposite polarity are degenerate, indicating symmetric and antisymmetric quantum superpositions of macroscopic states. The two classical states have persistent currents of 0.5 microampere and correspond to the center-of-mass motion of millions of Cooper pairs.

When a small magnetic field is applied to a superconducting loop, a persistent current is induced. Such a persistent supercurrent also occurs when the loop contains Josephson tunnel junctions. The current is clockwise or counterclockwise, thereby either reducing or enhancing the applied flux to approach an integer number of superconducting flux quanta $\Phi_0$. In particular when the enclosed magnetic flux is close to half-integer values of $\Phi_0$, the loop may have multiple stable persistent-current states, with at least two of opposite polarity. The weak coupling of the Josephson junctions then allows for transitions between the states. Previous theoretical work (2–4) proposed that a persistent current in a loop with Josephson junctions corresponds to the center-of-mass motion of all the Cooper pairs in the system and that quantum mechanical behavior of such persistent-current states would be a manifestation of quantum mechanical behavior of a macroscopic object. In a micrometer-sized loop, millions of Cooper pairs are involved. At very low temperatures, excitations of individual charge carriers around the center of mass of the Cooper-pair condensate are prohibited by the superconducting gap. As a result, the coupling between the dynamics of persistent supercurrents and many-body quasi-particle states is very weak. Josephson junction loops therefore rank among the best objects for experimental tests of the validity of quantum mechanics for systems containing a macroscopic number of particles (3, 5, 6) [loss of quantum coherence results from coupling to an environment with many degrees of freedom (7)] and for research on the border between classical and quantum physics. The potential for quantum coherent dynamics has stimulated research aimed at applying Josephson junction loops as basic building blocks for quantum computation (qubits) (8–11).

We present microwave spectroscopy experiments that demonstrate quantum superpositions of two macroscopic persistent-current states in a small loop with three Josephson junctions (Fig. 1A). At an applied magnetic flux of $\frac{1}{2}\Phi_0$, the system behaves as a particle in a double-well potential, where the classical states in each well correspond to persistent currents of opposite sign. The two classical states are coupled via quantum tunneling across the barrier between the wells, and the loop is a macroscopic quantum two-level system (Fig. 1B) (12). The energy levels vary with the applied flux as shown (Fig. 1C). Classically, the levels cross at $\frac{1}{2}\Phi_0$. Tunneling between the wells leads to quantum mechanical eigenstates that at $\frac{1}{2}\Phi_0$ are symmetric and antisymmetric superpositions of the two classical persistent-current states. The symmetric superposition state is the quantum mechanical ground state with an energy lower than the classical states; the antisymmetric superposition state is the loop’s first excited state with an energy higher than the classical states. Thus, the superposition states manifest themselves as an anticrossing of the loop’s energy levels near $\frac{1}{2}\Phi_0$. We performed spectroscopy on the loop’s two quantum levels (Fig. 2) and our results show the expected anticrossing at $\frac{1}{2}\Phi_0$ (Fig. 3) (13). We also studied the resonance-line shapes and found behavior similar to microscopic quantum two-level systems (14, 15) (Fig. 4).

Detecting quantum superposition. In our experiments, the magnetic flux generated by the loop’s persistent current was measured with an inductively coupled direct-current superconducting quantum interference device (DC-SQUID) (Figs. 1 and 2), while low-amplitude microwaves were applied to induce transitions between the levels. We observed narrow resonance lines at magnetic field values where the level separation $\Delta E$ was resonant with the microwave frequency. The DC-SQUID performs a measurement on a single quantum system. Thus, we should expect that the measurement process limits the coherence of our system. While the system is pumped by the microwaves, the SQUID actively measures the flux produced by the persistent currents of the two states. Detecting the quantum levels of the loop is still possible because the meter is only weakly coupled to the loop. The flux signal needs to be built up by averaging over many repeated measurements on the same system (Fig. 2B), such that an ensemble average is effectively determined. We measure the level separation, i.e., energy rather than flux, as we perform spectroscopy, and we observe a change in averaged flux when the microwaves are resonant with the level separation (the peaks and dips in Figs. 2B and 3A). We also chose to work with an extremely underdamped DC-SQUID with unsaturated junctions to minimize damping of the quantum system via the inductive coupling to the SQUID.

Similar observations were recently made by Friedman et al. (16) who performed spectroscopy on excited states in a loop with a...
Each well has persistent currents of opposite sign, with a magnitude \( I_p \) very close to \( I_c \), of the weakest junction and with energies \( E = \pm I_p (\Phi_{\text{ext}} - 1/2 \Phi_0) \) (dashed lines in Fig. 1C). The system can be pictured as a particle with a mass proportional to \( C \) in the Josephson potential; the electrostatic energy is the particle’s kinetic energy. The charging effects are conjugate to the Josephson effect. For low-capacitance junctions (small mass) quantum tunneling of the particle through the barrier gives a tunnel coupling \( t \) between the persistent-current states. In the presence of quantum tunneling and for \( E_c/E_C \) values between 10 and 100, the system should have two low-energy quantum levels \( E_0 \) and \( E_1 \), which can be described using a simple quantum two-level picture (10, 11).

The system was realized by microfabricating an aluminum micrometer-sized loop with unshunted Josephson junctions (30). Around the loop, we fabricated the DC-SQUID magnetometer (Fig. 1A), which contains smaller Josephson junctions that were as underdamped as the junctions of the inner loop. Loop parameters estimated from test junctions fabricated on the same chip and electron microscopy inspection of the measured device give \( I_c = 570 \pm 60 \) nA and \( C = 2.6 \pm 0.4 \) fF for the largest junctions in the loop and \( I_p = 0.82 \pm 0.1 \), giving \( E_c/E_c = 38 \pm 8 \) and \( I_p = 450 \pm 50 \) nA. Due to the exponential dependence of the tunnel coupling \( t \) on the mass (i.e., the capacitance \( C \)) and the size of the tunnel barrier, these parameters allow for a value for \( t/I_c \) between 0.2 and 5 GHz. The parameters of the DC-SQUID junctions were \( I_c = 109 \pm 5 \) nA and \( C = 0.6 \pm 0.1 \) fF. The self inductance of the inner loop and the DC-SQUID loop were estimated to be 11 \( \pm 1 \) picoHenry (pH) and 16 \( \pm 1 \) pH, respectively, and the mutual inductance \( M \) between the loop and the SQUID was 7 \( \pm 1 \) pH (31). The flux in the DC-SQUID is measured by ramping a bias current through the DC-SQUID and recording the current level \( I_{\text{sw}} \) where the SQUID switches from the supercurrent branch to a finite voltage (Fig. 2A). Traces of the loop’s flux signal were recorded by continuously repeating switching-current measurements while slowly sweeping the flux \( \Phi_{\text{ext}} \) (Fig. 2B). The measured flux signal from the inner loop will be presented as \( I_{\text{sw}} \), which is an averaged value directly deduced from the raw switching-current data (32).

**Results.** Figure 3A shows the flux signal of the inner loop, measured in the presence of low-amplitude continuous-wave microwaves at different frequencies \( f \). The rounded step in each trace at \( 1/2 \Phi_0 \) is due to the change in direction of the persistent current of the loop’s ground state (see also Fig. 1C). Symmetrically around \( \Phi_{\text{ext}} = 1/2 \Phi_0 \), each trace shows a peak and a dip, which were absent when no microwaves were applied. The positions of the peaks and dips in \( \Phi_{\text{ext}} \) depend on microwave frequency but not on amplitude. The peaks and dips result from microwave-induced transitions to the state with a persistent current of opposite sign. These occur when the level separation is resonant with the microwave frequency, \( \Delta E = hf \). In Fig. 3B, half the distance in \( \Phi_{\text{ext}} \) between the resonant peak and dip \( \Delta \Phi_{\text{res}} \) is plotted for all the frequencies \( f \). The relation between \( \Delta E \) and \( \Phi_{\text{ext}} \) is linear for the high-frequency data. This gives \( I_p = 484 \pm 2 \) nA, in good agreement with the predicted value. At lower frequencies, \( \Delta \Phi_{\text{res}} \) significantly deviates from this linear relation, demonstrating the presence of a finite tunnel splitting at \( \Phi_{\text{ext}} = 1/2 \Phi_0 \). To fit to Eq. 1 yields \( t/I_c = 0.3 \pm 0.03 \) GHz, in agreement with the estimate from fabrication parameters. The level separation very close to \( 1/2 \Phi_0 \) could not be measured directly because at this point the expectation value for the persistent current is zero for both the ground state and the excited state (Fig. 1C). Nevertheless, the narrow res-
The full step height of the rounded step at amplitude is ε in microwave amplitudes. The saturated dip units (a.u.), followed by a saturation for larger amplitudes, is very wide. The internal variable is coupled to a dissipative environment, and the associated effective mass (i.e., the capacitance across the SQUID) is very large. The internal degree of freedom has negligible intrinsic dephasing and the associated mass (i.e., the capacitance of the junctions of the SQUID) is very small. Consequently, this variable exhibits quantum behavior. The classical external degree of freedom of the SQUID performs a measurement on the SQUID’s inner quantum variable, which in turn is weakly coupled to our quantum loop. We therefore expect that the SQUID contributes dominantly to the loop’s dephasing and damping with the present setup. The choice for an underdamped DC-SQUID resulted in very wide switching-current histograms. The width of the histogram corresponds to a standard deviation in the flux readout of $11 \times 10^{-3} \Phi_0$. The uncertainty in flux readout is much larger than the flux signal from the inner loop $2M I_\text{SW} \approx 3 \times 10^{-3} \Phi_0$. Therefore, we can only detect the loop’s signal by averaging over many switching events (Fig. 2).

The loss of dip amplitude in Fig. 4 is probably due to a small contribution to the effective $\Phi_{\text{ext}}$ from the circulating current in the DC-SQUID. The SQUID is operated at 0.76 $\Phi_0$ in its loop, where its circulating current depends on the bias current due to its nonlinear behavior (I). This means that data recorded by switching events on the low $I_{\text{bias}}$ side of the FWHM of the $I_{\text{SW}}$ histogram in Fig. 2A differs in flux bias on the inner loop from that of the high side with $20 \times 10^{-6} \Phi_0$.

Fig. 2. (A) Current-voltage characteristic (inset) and switching-current histogram of the underdamped DC-SQUID. The plot with bias current $I_{\text{bias}}$ versus voltage $V$ is strongly hysteretic. The $I_{\text{bias}}$ level where the SQUID switches from the supercurrent branch to a finite voltage state—the switching current $I_{\text{SW}}$—is a measure for the flux in the loop of the DC-SQUID. Switching to the voltage state is a stochastic process. The histogram in the main plot shows that the variance in $I_{\text{SW}}$ is much larger than the flux signal of the inner loop’s persistent current, which gives a shift in the averaged $I_{\text{SW}}$ of about 1 nA (see Fig. 2B). (B) Switching-current levels of the DC-SQUID versus applied flux. The inset shows the modulation of $I_{\text{SW}}$ versus the flux $\Phi_{\text{SQUID}}$ applied to the DC-SQUID loop (data not averaged, one point per switching event). The main figure shows the averaged level of $I_{\text{SW}}$ (solid line) near $\Phi_{\text{SQUID}} = 0.76 \Phi_0$. At this point, the flux in the inner loop $\Phi_{\text{ext}} \approx \sqrt{3V/\Phi_0}$. The rounded step at $\Phi_{\text{ext}} = \sqrt{3V/\Phi_0}$ indicates the change of sign in the persistent current of the loop’s ground state. Symmetrically around $\sqrt{3V/\Phi_0}$ the signal shows a peak and a dip, which are only observed without measurement in the presence of continuous-wave microwaves (here 5.895 GHz). The peak and dip are due to resonant transitions between the loop’s two quantum levels (Fig. 3). The background signal of the DC-SQUID that results from flux directly applied to its loop (dashed line) is subtracted from the data presented in Figs. 3A and 4A.
Resonance lines at low $V_{AC}$ (i.e., with a FWHM $< 20 \cdot 10^{-6} \Phi_0$) cannot be observed as the peaks and dips smear out when averaging over many switching events. The loss of dip amplitude and the apparent saturation of the FWHM at low $V_{AC}$ is probably dominated by this mechanism for inhomogeneous line broadening.

The width of the rounded steps in the measured flux in Fig. 3A is much broader than expected from quantum rounding on the scale of the value of $t$ that was found with spectroscopy (see also Fig. 1C). We checked the temperature dependence of the step width, measured in the absence of microwaves. We found that at temperatures above 100 mK, the step width is in agreement with the thermally averaged expectation value for the persistent current ($I_{th} = I_p \tanh (\Delta E / 2k_B T)$) where we use the level separation $\Delta E$ and $I_p$ found with spectroscopy. However, at low temperatures the observed step width saturates at an effective temperature of about 100 mK. We checked that the effective temperature for the SQUID’s switching events did not saturate at the lowest temperatures. The higher effective temperature of the loop is a result of the loop being in a nonequilibrium state. The population of the excited state could be caused by the measuring SQUID or other weakly coupled external processes.

Concluding remarks and future prospects. The data presented here provide clear evidence that a small Josephson junction loop can behave as a macroscopic quantum two-level system. The application of an underdamped DC-SQUID for measuring the loop’s magnetization is a useful tool for future work on quantum coherent experiments with Josephson junction loops. The present results also demonstrate the potential of three-junction persistent-current loops for research on macroscopic quantum coherence and for use as qubits in a quantum computer. This requires quantum-state control with pulsed microwaves and development of measurement schemes that are less invasive. Circuits that contain multiple qubits with controlled inductive coupling are within reach using present-day technology.

References and Notes

12. The double-well potential shown in Fig. 1B is the relevant potential for the tunnel transitions between the system’s two persistent-current states. The potential is the sum of the Josephson energies of all the junctions and the zero-point energy of a Josephson phase degree of freedom that is perpendicular to the tunnel direction between the two wells (70, 71). This zero-point energy is nearly constant along the tunnel trajectory between the wells, and it has a plasma frequency much higher than the level separation between the lowest two quantum levels. It should therefore be included in the effective potential for transitions between the two wells. For the sample parameters mentioned in the text the lowest two quantum levels are well below the top of the effective tunnel barrier.
13. The results presented here do not exclude alternative theories for quantum tunneling in macro-realistic theories [37]. This would require a type of experiment as proposed by Leggett et al. (5).
Triple Vortex Ring Structure in Superfluid Helium II

Demosthenes Kivotides, Carlo F. Barenghi,* David C. Samuels

Superfluids such as helium II consist of two interpenetrating fluids: the normal fluid and the superfluid. The helium II vortex ring has generally been considered merely as a superfluid object, neglecting any associated motion of the normal fluid. We report a three-dimensional calculation of the coupled motion of the normal-fluid and superfluid components, which shows that the helium II vortex ring consists of a superfluid vortex ring accompanied by two coaxial normal-fluid vortex rings of opposite polarity. The three vortex rings form a coherent, dissipative structure.

Vortex rings \( (I) \) have long been studied as ideal examples of organized flow structures. A large body of literature has been concerned with vortex rings in a zero-viscosity (inviscid) fluid in which the vortex core thickness is much smaller than the ring’s radius. This mathematical idealization is realized in a quantum fluid \( (2, 3) \), helium II, which is a superposition of two fluid components: the normal fluid (which is a fluid with nonzero viscosity) and the superfluid (an inviscid fluid). The concept of the superfluid vortex ring \( (+) \) or loop has contributed to many advances in superfluidity, ranging from vortex creation \( (5, 6) \) to turbulence \( (7–10) \). An example of this is the fundamental issue of quantum mechanical phase coherence and the onset of dissipation. Ions injected into superfluid helium II move without friction, provided that the speed does not exceed a critical value \( (11) \) above which superfluid vortex rings are created \( (5) \). Vortex creation \( (12, 13) \) and motion \( (14, 15) \) have been studied theoretically using various models and are also being investigated by atomic physicists in the context of Bose-Einstein condensation in clouds of alkaline atoms \( (15, 16) \). The concept of the vortex ring has been applied to interpretations of the nature of the roton \( (17–19) \) and the superfluid transition itself \( (20) \). Finally, vortex rings are important in the study of superfluid turbulence, which manifests itself as a disordered tangle of superfluid vortex loops (distorted vortex rings). Superfluid vortex lines may also end at walls, or at free surfaces, without forming closed loops. For simplicity, we will consider here a circular superfluid vortex ring, but our results should also apply to all superfluid vortex lines.

Recent experiments, such as the observation of decay rates of superfluid vorticity \( (21, 22) \) consistent with the decay rates of Navier-Stokes turbulence, motivate our study of the dynamical coupling between the superfluid vorticity and the normal-fluid component. Superfluid vorticity scatterers \( (23) \) the thermal excitations that make up the normal fluid, producing a mutual friction acting on the velocity fields \( V_t \) and \( V_n \) of the two fluid components of helium II. Although the superfluid vorticity can be detected directly by the second sound technique \( (21) \), very little is actually known about the normal-fluid flow because we have no practical flow visualization techniques near absolute zero. We present results of a three-dimensional calculation in which \( V_t \) and \( V_n \) are determined self-consistently. The calculation reveals the surprising triple structure of the helium II vortex ring. We also discuss the implications of this finding for the interpretation of current turbulence experiments.

Our method is based on an improvement over the vortex dynamics approach of Schwarz \( (24, 25) \), who modeled a superfluid vortex line as a curve \( S(\xi, t) \) that obeys the...