Tuned transition from quantum to classical for macroscopic quantum states

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The boundary between the classical and quantum worlds has been intensely studied. It remains fascinating to explore how far the quantum concept can reach with use of specially fabricated elements. Here we employ a tunable flux qubit with basis states having persistent currents of 1 µA carried by a billion electrons. By tuning the tunnel barrier between these states we see a crossover from quantum to classical. Released from non-equilibrium, the system exhibits spontaneous coherent oscillations. For high barriers the lifetime of the states increases dramatically while the tunneling period approaches the phase coherence time and the classical regime is reached.

The quantum nature of quarks and atoms is as solidly established as the relevance of Newtonian mechanics for marbles and soccer balls. The boundary between the two worlds has been studied theoretically [1], but only a few experiments are available so far. It has been demonstrated that objects containing many atoms, such as large molecules [2], magnetic particles [3] or fabricated superconducting circuits [4, 5] can behave like single quantum particles. In this Letter we performed an experiment on a superconducting flux qubit, which is the ‘classical’ example of a macroscopic object that can be made to behave as a single quantum particle. It is characterized by two states with opposite macroscopic currents in a loop. We were able to control the tunnel barrier between these states over a very wide range. We tuned qubit energy levels below the barrier and at the same time effectively cool the sample to near zero temperature. This allowed us to study the qubit behavior when we go from the range of low barriers and strong quantum tunnelling to the regime where quantum tunnelling gradually disappears as the barrier is increased. In particular, we manage to observe the natural quantum oscillations manifested in the tunnelling of the long-living macroscopic magnetic moments. At very high barriers we see how these oscillations fade away as the barrier increased.

The flux qubit has a potential energy which consists of two degenerate wells [Fig. 1(a)] separated by a barrier $E_B$. Each well is connected with a macroscopic magnetic flux, with a sign (+/−) depending on being in the left or the right well, which can be detected on demand by a measurement apparatus. The zero-point energy $E_0$ of the qubit in each well can be made smaller than $E_B$ [6, 7]. Consequently, the barrier between the wells becomes classically impenetrable, and at low temperature the magnetic moment of the qubit can be flipped only via the quantum tunnelling process. This process is represented in Fig. 1(a) by the tunnelling coupling $\Delta$, which depends exponentially on the barrier height $E_B$. Thus, the flux qubit is particularly suited for testing quantum-to-classical crossover, as it integrates two seemingly contradictory features of the classical and quantum world: the macroscopic character of the quantum states with the fundamentally non-classical quantum tunnelling. Other attempts to study the quantum-to-classical transition on a single quantum system were based on increasing the effect of the noise on a system with fixed quantum interactions [8]. To probe the quan-

FIG. 1: Experiment to test macroscopic quantum coherence: (a) Potential energy of the flux qubit. The barrier height $E_B$ is large compared to the zero-point energy $E_0$ resulting in the strong localization of the well ground states $|L\rangle$ and $|R\rangle$. These states are connected to macroscopically distinguishable magnetic moments induced by the persistent currents in the flux qubit loop. The tunnel coupling $\Delta$ between the current states is exponentially dependent on $E_B$. The macroscopic quantum tunnelling $|L\rangle \leftrightarrow |R\rangle$ can be observed by a detector (DC-SQUID) sensitive to a change of the magnetic flux in the qubit loop. (b) Scanning electron micrograph of the flux qubit.
FIG. 2: Tunable flux qubit: (a) Schematics. The qubit (blue) is formed by three Josephson junctions one of which is a tunable double junction. The dark blue and red arrowed lines show the persistent currents connected to the states $|L\rangle$ and $|R\rangle$. The qubit state can be controlled by the bias lines $I_e$, $I_{e,dc}$, $I_o$ (black) and measured by the DC SQUID (green). (b) Potential energy (in units of the Josephson energy of the regular junctions) as a function of the two independent phase differences $\gamma_1$ and $\gamma_2$. (c) Sketch of the cross section of the potential energy along the white line connecting left and right wells through the saddle point (see (b)). The energy eigenstates $|0\rangle$ and $|1\rangle$ are superpositions of the persistent current states $|L\rangle$ and $|R\rangle$. (d) Energy diagram of the qubit vs. magnetic bias $\epsilon$.

Our flux qubit consists of three junctions symmetrically attached to a trap loop as shown in Fig. 2(a). The central junction is made tunable by replacing it by two junctions in parallel, thus providing control over $E_B$ and so $\Delta$. The trap loop is employed to capture a fluxoid (or $2\pi$-phase-winding number) [11], establishing a $\pi$ phase drop over the qubit junctions. If one fluxoid is trapped and the difference in flux in the two loop halves of the gradiometer $2f_e\Phi_0 = (f_1 - f_2)\Phi_0 \approx 0$ the system has a double well potential [Fig. 1(b)]. Here $\Phi_0$ is the magnetic flux quantum $h/(2e)$ and $f_e,1$, $f_e,2$ are the fluxes in units of $\Phi_0$. The ground states in each well of the potential are persistent current states $|L\rangle$ and $|R\rangle$ characterized by the currents $\pm I_p$ carried by the junctions, generating the before mentioned $+/-$ magnetic moments [Fig. 2(a)]. The energy eigenstates of the qubit are linear superpositions of $|L\rangle$ and $|R\rangle$ [Fig. 2(c)], following the Hamiltonian

$$H = -\hbar \frac{1}{2} \left( \epsilon (f_e, f_o) \sigma_z + \Delta (f_o) \sigma_x \right), \quad (1)$$

where $h\epsilon = 2I_p f_e\Phi_0$ is the magnetic energy bias and $\sigma_{x,z}$ are Pauli matrices. The critical currents of the four junctions are designed such that the two parallel junctions each have half the value of the critical current $I_0 \approx 700\,nA$ of the other two junctions. Applying insitu flux $f_o \Phi_0$ to the parallel junctions sets their total effective critical current to $I_0 \cos(\pi f_o)$, in this way allowing $f_o$ to control $E_B$ and $\Delta$ [12]. Note that, while $\Delta (f_o)$ depends strongly on $f_o$, the persistent current magnitude $I_p$ only shows a weak dependence. Qubit excitation is obtained by the magnetic field generated by current in the symmetrized-split $I_e$ line, acting on the qubit flux $f_e\Phi_0$. Similarly, the line $I_o$ together with the homogeneous field $B$ generated by an external coil, sets $f_o \Phi_0$ and changes $\Delta$. The geometrical symmetry leads to independent control of $\epsilon$ and $\Delta$. The qubit states are detected with a DC-SQUID which is coupled to the qubit by a shared wire with a mutual qubit-SQUID inductance $M \approx 6\,pH$.

Figure 3(a) shows the gap of the qubit for different $f_o$ and deduced from spectroscopy performed with the following protocol. First we set $\Delta$ with the field $B$ and apply a DC current $I_{e,dc}$ to tune the qubit frequency to $\nu_{qdb} \equiv (\Delta^2 + \epsilon^2)^{1/2} \approx 9\,GHz$. In the second step we apply a square current pulse $I_e$, shifting the qubit frequency, combined with a microwave excitation. Next, the qubit is returned to $\nu_{qdb} = 9\,GHz$ and a short bias current pulse $I_b$ is applied to the SQUID detector to measure the qubit state. The relative populations of the qubit ground and excited states determine the expectation value of the persistent current, generating an additional magnetic flux in the SQUID, resulting in a change of the SQUID switching probability (color scale) [13].
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The sequence starts at \( \nu_{\text{qb}} \sim 9 \text{ GHz} \), far above the effective noise temperature \( T_e \sim 50 - 100 \text{ mK} \) (\( \sim 1-2 \text{ GHz} \)) and the cryostat base temperature \( T_0 \equiv 20 \text{ mK} \). After waiting for a long enough time the qubit relaxes to its ground state. During later operations the qubit splitting can be reduced to values below \( T_e \) or even \( T_0 \). With the ability of producing fast energy shifts we achieved full coherent control of the qubit quantum state even for very small energy splittings for a duration limited by the relaxation time \( T_1 \). Coherent transitions below the thermal energy have been realized previously in superconducting qubits only with active microwave pulses [7] [12] similar to laser cooling used in atomic systems [16].

In Fig. 3(a) one can see that the gap covers nearly two decades, ranging from 150 MHz to 12 GHz, allowing us to study coherent transitions even in a regime of very small gaps. Over the same \( f_\alpha \) range \( I_p \) varies from 600 nA to 150 nA. Figs. 3(b,c) show spectra for two representative cases. For the regular flux qubit gap range (\( \Delta \sim 2 - 10 \text{ GHz} \)) our numerical simulations show that the qubit ground state levels lie above the barrier for the double well potential [Figs. 3(b,d)]. Only when the gap \( \Delta \) drops below 500 MHz the qubit levels fall below the barrier and the transitions between the wells become classically forbidden [Figs. 3(c,e)]. Note that \( \Delta \) closely follows an exponential dependence on \( f_\alpha \) over the full range of \( \Delta \), a feature exclusively associated with quantum tunneling.

In order to demonstrate the emergence of the classical opaqueness of the barrier and the macroscopic nature of the persistent current states we measured the relaxation time for a small gap \( \Delta = 200 \text{ MHz} \) as a function of the qubit frequency \( \nu_{\text{qb}} \) [Fig. 4]. From [1] it follows that \( |0\rangle |1\rangle \sim |1\rangle + (-\cos \theta) |L\rangle + \sin \theta |R\rangle \), where \( \tan \theta \equiv \Delta/\varepsilon \). Thus starting from \( \nu_{\text{qb}} = \Delta^2 \) the energy eigenstates are gradually transformed from (anti)symmetric superpositions of \( |L\rangle \) and \( |R\rangle \) states to being almost purely \( |L\rangle \) and \( |R\rangle \) at \( \nu_{\text{qb}} = 6 \text{ GHz} \gg \Delta \). The measurement shows nearly three orders of magnitude increase in lifetime of the excited state for the localized persistent current state \( |R\rangle \) compared to the delocalized superposition (\( |L\rangle - |R\rangle \))/\( \sqrt{2} \), reaching hundreds of \( \mu \text{s} \). These high values of \( T_1 \) demonstrate the extreme robustness of the persistent current states.

We used the experimental sequence shown in Fig. 5(a) for time-resolved detection of macroscopic quantum coherence. We start by tuning \( \Delta \) below 300 MHz with the magnetic field \( B \) to enter the deep tunnelling regime. Using \( I_{\text{dc}} \), we also tilt the double-well potential, preparing the qubit in its ground state \( |L\rangle \) with \( \nu_{\text{qb}} = 7 \text{ GHz} \). Subsequently, the double-well is made symmetric by means of a fast \( I_p \) pulse; in 0.3 ns the qubit is taken to its symmetry point. As the qubit energy changes fast relative to the tunnelling amplitude \( \Delta \), this transfer is non-adiabatic thus preserving the initial state occupation. These oper-
leads to a total destruction of the quantum phase between the persistent current states and the system is no-longer regarded as quantum. It is interesting to note that, for small gaps, sensitivity to $f_{\alpha}$ noise is strongly suppressed, and so the environment automatically chooses the basis of the macroscopic current states for dephasing.

Our measurements shows how the flux qubit can be gradually tuned from the quantum to the classical regime. With the increase of the tunnel barrier the ‘quannumness’ of the system manifested in coherent tunnelling is gradually lost. At the same time the lifetime of the persistent current states dramatically increases, which is naturally associated with macroscopic classical systems or classical bits. Also, over a large range of parameters the quantum and macroscopic properties are shown to coexist. Our experiment demonstrates the potential of fabricated quantum objects, where knobs are available to tune parameters in situ, for fundamental research as well as for applications.

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FIG. 5: Macroscopic quantum coherence of the persistent current states. (a) Measurement protocol: preparation of the qubit in $|L\rangle$ at $\nu_{qb} = 7$ GHz by strongly tilting the double-well potential; a fast shift to the symmetry point (symmetric double-well); free evolution for time $t$ and a fast shift back to $\nu_{qb}$, followed by the SQUID measurement pulse. (b) Time-resolved measurement of the tunnelling of the persistent current states for $\Delta = 262$ MHz (blue) and fit to $\cos(\Delta t)$ (red). (c-e) The colors indicate the switching probability of the SQUID. The horizontal scale represents the amplitude of the current pulse $I_c$ sweeping the qubit through the symmetry point where the macroscopic quantum coherence oscillations are clearly observed for $\Delta = 200$ MHz (c) and $\Delta = 150$ MHz (d). Residual oscillations are seen even for $\Delta = 90$ MHz (e) representing the boundary between quantum and classical state of the system.
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