

# Coherent manipulations of charge-number states in a Cooper-pair box

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We have investigated coherent properties of an artificial two-level system realized in a small-Josephson-junction circuit. Two charge-number states of a small superconducting electrode connected to a reservoir via a Josephson junction were used as the two relevant states. Coherent manipulations of the charge-number states by applying a gate-voltage pulse were demonstrated, and decoherence of the system was studied through a spin-echo-type experiment. The result suggested that low-frequency energy-level fluctuations due to  $1/f$  background charge noise is a dominant source of the dephasing.

## I. INTRODUCTION

A solid, or a solid-state circuit, contains an extremely large number ( $\sim N_A$ ; Avogadro constant) of degrees of freedom. They are usually strongly interacting, and the lifetime of a particular elementary excitation is rather short. Some of them, for example, nuclear spin degrees of freedom, are well isolated from others, but have many identical ones, and are not easily accessed individually. Consequently, it is not straightforward to select a single degree of freedom out of the huge set and coherently manipulate it in order to observe its quantum mechanical behavior.

Nevertheless, in a simple electric circuit containing a small Josephson junction, we can realize a single two-level system by using interplay between the Josephson effect and the charging effect [1–4]. Indeed, two collective variables in the Josephson junction, the charge, i.e., a number of Cooper pairs which have passed the junction, and the phase difference across the junction are conjugate variables like a position and a momentum of a particle. If one of them is confined, then the other fluctuates quantum mechanically. Moreover, the coupling of those collective variables to other degrees of freedom can be designed to be quite small. Coherent manipulations of quantum states in such a two-level system have been recently demonstrated experimentally [5,6].

In addition, recently there has been much interest in such a two-level system as a possible candidate of a quantum bit (qubit) for quantum computing [7]. Several theoretical proposals have been made on implementations of quantum computing using Josephson-junction qubits [8–15]. As a qubit realized in a solid-state device, those qubits are believed to have an advantage in the scala-

bility to an integrated quantum circuit. On the other hand, whether other requirements for quantum computing, such as precise control, accurate quantum-state readout, and long decoherence time, are fulfilled is not guaranteed a priori. Therefore, progresses in experiments to really demonstrate the potential of those qubits are anticipated.

In this paper, our experimental studies on a Josephson-junction charge two-level system are reviewed. After introducing the two-level system, a Cooper-pair box, a quantum-state readout process used in our experiment is described. We have demonstrated coherent manipulations of superposition of two charge-number states by using gate-voltage pulses [5,16,17]. However, there are still many unknown factors about the decoherence in the system. A well-known technique in NMR physics, i.e., spin-echo technique, has been adapted as “charge echo” in our qubit to study the effect of low-frequency energy-level fluctuations on the dephasing of the qubit [18].

## II. A COOPER-PAIR BOX

For constructing a qubit in a solid-state device, it seems to be a natural choice to use a superconductor. In a superconductor, conduction electrons form Cooper pairs and condense into a single macroscopic ground state, that is, a superconducting ground state. Consequently, a superconducting energy gap  $\Delta$  opens in the density of states of the electronic excitations. The gap prevents low-energy excitation caused by perturbations and helps to maintain the coherence of the system.

However, a single ground state is not enough for a qubit; we need at least two energy levels within the excitation gap. How about using two superconductors, then?

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Suppose there are two superconductors coupled via a thin insulating barrier (Josephson junction; Fig. 1(a)). Let us define charge-number states with  $|n\rangle$  where  $n$  is a number of excess Cooper pairs in the right electrode. Because of the conservation of the total charge, the left electrode has  $(-n)$  excess Cooper pairs. The charge-number states are the ground state of the two superconductors with such a particular charge distribution. Each pair of nearest-neighbor charge-number states is coupled by coherent transfer of a Cooper pair through the junction. The transfer energy is given by Josephson energy  $E_J = \hbar\Delta/(8e^2R_J)$  [19], where  $R_J$  is the normal-state resistance of the junction. In the case of a large junction in which the charging effect is negligible, all the charge(-number) states are degenerate and, thus, form a one-dimensional tight-binding array in the charge space (Fig. 1(b)). As a result, the phase difference  $\theta$  between the two superconductors, a conjugate variable to the charge number  $n$ , becomes a good quantum number, and we obtain a single Bloch band in the phase space (Fig. 1(c)). This means that there is still a single available state within the superconducting gap.

On the other hand, if we reduce the size of the junction, the charging energy due to the excess Cooper-pair charge becomes dominant, and the degeneracy between charge states is lifted. We consider a small Josephson junction biased with gate voltage  $V_g$  through gate capacitance  $C_g$ . The simple circuit is called a Cooper-pair box (Fig. 1(d)) [20]. The single-electron charging energy of the “box” electrode is given by  $E_C \equiv e^2/(2C_\Sigma)$ , where  $C_\Sigma \equiv C_J + C_g$  is the total capacitance of the box electrode, a sum of the junction capacitance and the gate capacitance. The electrostatic energy of the charge state  $|n\rangle$  is  $4E_C(n - Q_g/2e)^2$  which is quadratic in  $n$  as shown in Fig. 1(e) for  $Q_g/e = 1$ . Here,  $Q_g \equiv C_g V_g$  represents the charge induced by the gate voltage.

When  $E_C > E_J \gg k_B T$ , we can treat the system as an effective two-level system by taking into account only the two lowest-energy charge states. Without losing generality, we focus on the states  $|0\rangle$  and  $|1\rangle$  at around  $Q_g/e = 1$ . In Fig. 1(f), electrostatic energies of the two states are plotted with dashed lines as a function of  $Q_g$ . Because of the Josephson coupling between the two charge states, the eigenenergies drawn with solid lines show anticrossing at  $Q_g/e = 1$ , where the two charge states are degenerate. The next-lowest-energy state ( $|-1\rangle$  or  $|2\rangle$ ) is well separated from the two eigenstates by about  $8E_C - E_J/2$  to  $4E_C$  for  $|Q_g/e - 1|$  from 0 to 0.5.

If we map the system on a spin  $\frac{1}{2}$  in a magnetic field, the effective Hamiltonian is written as

$$H = \frac{1}{2}\delta E(Q_g)\sigma_z - \frac{1}{2}E_J\sigma_x, \quad (1)$$

where  $\delta E(Q_g) \equiv 4E_C(Q_g/e - 1)$  is the electrostatic-energy difference between the two charge states, and  $\sigma_z$  and  $\sigma_x$  are Pauli matrices. The fictitious magnetic field

is given by  $(E_J, 0, -\delta E(Q_g))$ .

Cooper-pair-box samples are fabricated by electron-beam lithography and shadow evaporation of Al films [21]. Because of the collective nature of the superconducting condensate and the short electric shielding length due to the high electron density in metallic superconductors, the two important parameters  $E_J$  and  $E_C$  do not strongly depend on microscopic disorders in the device and are well controlled by the geometry of the fabricated sample. A scanning-electron micrograph of a sample and a schematics are shown in Fig. 2. The bright parts are the Al electrodes evaporated on a substrate covered with an insulating  $\text{SiN}_x$  layer. A thin Al strip in the middle is the box electrode. A Cooper pair tunnels through an  $\text{AlO}_x$  tunnel barrier sandwiched between the box and the reservoir electrodes. In our samples, two parallel junctions are used to modulate the effective  $E_J$  by applying magnetic field through the SQUID loop. The total junction area is about  $0.01 \mu\text{m}^2$  and the resistance is about  $10 \text{ k}\Omega$ . On the right end of the box electrode, there is another small junction for readout of the qubit state. The detail is discussed in the next section. There are two gate electrodes for controlling energy levels of the qubit. One of them is connected to a room-temperature high-speed pulse generator via an on-chip coplanar line and a low-loss coaxial cable. The other gate for a dc gate voltage as well as the reservoir and the probe electrode are connected to dc-signal lines, which are filtered by thin lossy coaxial cables [22].

The samples were embedded in a shielded copper block and mounted on a cold finger of a dilution refrigerator whose base temperature is about  $30 \text{ mK}$ . Typically, the superconducting gap  $\Delta$  in the Al electrodes is  $50\text{--}60 \text{ GHz}$  in a frequency scale ( $= \Delta/\hbar$ ).  $E_C$  and  $E_J$  of the samples are about  $25 \text{ GHz}$  and  $10\text{--}20 \text{ GHz}$ , respectively.

So far, we have considered only “even-number” state, in which the box contains even number of electrons and they all form Cooper pairs. However, if  $\Delta < E_C$ , an “odd-number” state with a single unpaired electron in the box becomes a ground state at around  $Q_g/e = 1$  [23,24]. Hence, for a qubit experiment,  $\Delta > E_C$  is basically required.

### III. READOUT OF THE CHARGE STATES

Before describing the coherent dynamics of the qubit, we first discuss the scheme we used for the quantum-state readout. For the readout, an additional electrode is attached on the box electrode via a tunnel junction (Fig. 2). The “probe” electrode is voltage-biased with an appropriate bias voltage  $V_b$ , so that two quasiparticles tunnel out sequentially only when there is one excess Cooper pair in the box [25–27]. This is an analogy of a selective-field ionization in atomic physics. A Cooper-pair box in  $|1\rangle$  state is selectively “ionized” by

the electric field across the probe junction. The upper limit for  $V_b$  is to avoid the third quasiparticle tunneling and the lower limit is to supply enough bias energy for the first and the second quasiparticle tunneling each accompanied with a breaking of a Cooper pair in the box [28]. The quasiparticle tunneling rates for the first and the second processes are estimated from Fermi's golden rule as  $\Gamma_{qp1(2)} = \frac{\Delta E_{qp1(2)}}{e^2 R_b} W\left(\frac{\Delta E_{qp1(2)}}{2\Delta}\right)$ , where  $\Delta E_{qp1(2)} \equiv eV_b + (-)E_C + 2E_C(1 - Q_g/e)$  is the energy dissipated in the quasiparticle tunneling, and  $W(x)$  is a factor related to the BCS quasiparticle density of states, which is  $\sim 1$  for  $x \geq 1$  and  $\sim 0$  for  $x < 1$  and  $k_B T \ll \Delta$  [29]. Because of the charging effect, in the  $|0\rangle$  state, quasiparticle tunneling is Coulomb-blockaded. Therefore, by using this quasiparticle tunneling process, we can distinguish the two charge states. Also, this process initializes the qubit state to  $|0\rangle$  state after the readout. The two quasiparticles left in the box immediately recombine, and the box is initialized to be a ground state with no excess Cooper pair.

The readout is basically on the charge basis (along  $z$ -axis of the pseudo spin). Therefore, if the qubit basis is tilted from the charge basis, the readout cannot be a perfect one. Suppose in our qubit we have  $E_C = 5E_J$  and bias the qubit at  $Q_g/e = 0.5$ . Then, the fidelity of the readout is  $(\delta E/\Delta E)^2 \sim 99\%$ , where  $\Delta E \equiv \sqrt{E_J^2 + \delta E^2}$ .

Although the setup is so simple and the scheme really worked in the experiments described below, there are two obvious drawbacks in the readout scheme. One is the permanent coupling of the “detector” to the qubit. Because of that, the coherence of the qubit is disturbed by the stochastic escape of the first quasiparticle at a rate of  $\Gamma_{qp1}$ , that sets the upper limit of the decoherence time of the qubit in our setup. Therefore, to see the effect of coherent dynamics in the qubit, a condition  $E_J \gg \hbar\Gamma_{qp1}$  should be fulfilled. That is possible with a large probe-junction resistance  $R_b$ . Typically, we use  $R_b$  of 30–50 M $\Omega$ , which is more than a thousand times larger than  $R_J$ . Then,  $\hbar\Gamma_{qp1}$  is about 20 MHz. On the other hand, too small  $\Gamma_{qp1}$  means that it takes a long time before the readout takes place stochastically. If there were other faster relaxation mechanisms of the qubit states, the readout result could be incorrect.

In principle, though we have not successfully demonstrated, we can effectively switch on and off the detector by changing  $V_b$  and using the strong nonlinearity in  $W(x)$ . Then, it is also possible to decrease  $R_b$  and increase  $\Gamma_{qp1}$  so that “stronger” measurement can be done when the detector is on.

The other drawback is the small output signal of the readout process. The output is only two electrons tunneling through the probe junction. In principle, we can store the electrons in a small electrode and detect the charge by using, for example, a single-electron transistor. Then, we could do a single-shot readout. However,

that would result in additional complications because of the charging effect in the additional storage electrode. Instead, in our experiments, we used ensemble measurement. Nominally identical quantum-state controls and readouts were repeated with a repetition time  $T_r$ , and the output electrons were accumulated to give rise to a measurable amount of dc current through the probe junction. For a completion of the readout and a preparation of the initial state for the next control,  $T_r$  should be sufficiently longer than  $(1/\Gamma_{qp1} + 1/\Gamma_{qp2})^{-1}$ . On the other hand, the maximum output current  $2e/T_r$  should not be too small compared to the current measurement noise floor.

This probe-junction readout scheme has one unique feature. Generally, for a quantum-state readout, a dissipative process is required in order to transfer quantum information into classical information. However, often the dissipative element in the detector can cause some amount of back action even in the off-state of the detector, that may decohere the qubit. In the present case, the dissipative process used in the detector is the energy relaxation of a quasiparticle in the continuous density of states above the gap energy. The quasiparticle injected in a high-energy state is strongly damped by electron-electron interaction and electron-phonon interaction. Since this intra-electrode process only works on the tunneling quasiparticles and not on the charge degree of freedom of the Cooper-pair box, there is no back action before the quasiparticle tunneling takes place. Moreover, the state after the first quasiparticle tunneling is not in the Hilbert space spanned by the two-charge states. The system travels outside of the qubit Hilbert space and then comes back to it after the second quasiparticle tunneling and recombination of two quasiparticles left in the box.

The single readout process with a small output without amplification may seem to be useless. For a final output of a quantum computer, we would definitely need an amplifier that produces a measurable output signal from a quantum-scale signal. However, we would like to point out a possible application of this type of readout in quantum computing. In the quantum error-correction scheme [30–32], although we need many ancilla qubits prepared in  $|0\rangle$  state and many readouts of those qubits, the signals of the readout are not necessarily amplified to the level that we can recognize. We might be able to devise a feedback quantum-state control for the error correction by using the small unamplified signal. The readout process can also be used as an efficient initialization process. Our measurement is not a projective measurement, but always ends with  $|0\rangle$  state and an output signal. So, in principle, we can also feedback the signal to the measured qubit to mimic a projective measurement.

#### IV. COHERENT MANIPULATIONS OF THE CHARGE STATES

For coherent manipulations of the charge states, any modulation of the parameters in the Hamiltonian (Eq. (1)) can be used, as long as the modulation is fast enough to be nonadiabatic [33]; otherwise the initially prepared ground state just follows the ground state adiabatically and comes back to the same state except for the Berry phase [34]. Here we do not discuss the so-called geometrical control using the Berry phase [35–37] and focus on dynamical control of the quantum state.

In our experiments,  $\delta E$  was modulated by a gate-voltage pulse. In principle, we can also modulate  $E_J$  by applying a time-dependent magnetic field on a SQUID-type Josephson junction (Fig. 2). For an observation of coherent oscillations between two degenerate charge states, a single gate-voltage pulse was applied. Starting from the initial bias point  $Q_g/e \sim 0.5$  ( $|\delta E| \gg E_J$ ) where we prepare the initial state  $\sim |0\rangle$ , the system is abruptly brought by the pulse to the degeneracy point  $Q_g/e = 1$ , i.e.,  $\delta E(Q_g) = 0$  (cf. Fig. 1(f)). For a finite pulse duration  $\Delta t$ , the quantum state of the qubit oscillates between the two charge states. In the pseudo-spin model, the spin starts from an up-spin state along  $z$ -axis and rotates around  $x$ -axis with a frequency of  $E_J/\hbar$ . After the pulse, the system comes back to the initial gate voltage with a final state, a superposition of charge states whose coefficient depends on  $\Delta t$ . Because of the large  $\delta E(Q_g)$ , the pseudo spin now precesses almost horizontally. Thus, even by using the quasiparticle readout which has much slower time constant  $1/\Gamma_{qp1}$  compared to the precession period, we can readout the  $z$ -component of the spin. The result of the quantum-state control was measured by the probe-junction scheme. As a function of the pulse width  $\Delta t$ , current through the probe junction oscillated reflecting the oscillations of quantum states between two charge states [5,16].

We can also use an ac gate-voltage pulse for the quantum-state control utilizing Rabi oscillations, as people usually do in atomic physics and NMR physics. When the ac driving frequency matches the energy separation between the two eigenstates, coherent quantum-state evolution between the two states takes place involving photon energy of the driving field. In our experiment, because of the lack of technology to create a short-enough ac pulse — short enough for applying a pulse within the decoherence time — we instead used an added signal of continuous-wave microwave and a dc pulse. The dc pulse shifted the bias point from off-resonant to resonant with the microwave frequency. The observed Rabi-oscillation frequency depended on the power of the ac driving field, and Rabi oscillations involving multiple photons also took place [17].

#### V. DECOHERENCE IN A COOPER-PAIR BOX

##### A. decoherence at degeneracy point

To know the decoherence time and the mechanism, we first tried to see the decay of the coherent oscillations as a function of the pulse width. This corresponds to the decoherence at the degeneracy point  $Q_g/e = 1$ . We have observed those oscillations up to 2-5 ns depending on samples [16,38]. However, the signals were very noisy and not good enough for deducing a reliable decay envelope. The nonideal shape of the decay envelope could be partly due to the irregular shape of the applied pulse. Furthermore, the lifetime of the oscillations was not so reproducible from run to run even for the same sample, perhaps because of instability of the offset charge (see below). Therefore, we have had only very limited information about the decoherence at the degeneracy point.

##### B. Free-induction decay

We also looked at an off-degeneracy point ( $|\delta E| \gg E_J$ ) to know the mechanism of the decoherence [18]. For that purpose, we applied two short ( $\Delta t = 80$  ps) pulses with a delay time  $t_d$  in between.  $\Delta t$  and  $t_d$  were still much shorter than  $1/\Gamma_{qp1}$ . The height of the pulses was adjusted so that each pulse works approximately as a  $(\pi/2)_x$ -pulse, which rotates the pseudo spin by 90 degrees around  $x$ -axis. Starting from the ground state, the first  $(\pi/2)_x$ -pulse creates an equal-weight superposition of the two eigenstates. During the delay time, the qubit phase evolves because of the energy difference  $\Delta E$  between the two eigenstates. And the second  $(\pi/2)_x$ -pulse projects the phase information on  $z$ -axis, i.e., the readout axis. As a function of  $t_d$ ,  $z$ -component of the final state oscillates as  $\cos(\Delta E t_d/\hbar)$ . This is a direct analogy of a Ramsey interference experiment in atomic physics or a free-induction-decay experiment in NMR [40]. In the experiment, the oscillating signal decayed rapidly within a few-hundred picoseconds (Fig. 3(a)). This time scale was much shorter than that of the single-pulse experiment at the degeneracy point. This fact could be explained if the dominant decoherence source is the dephasing due to charge fluctuations. Because of the energy-level anti-crossing,  $\Delta E$  is less sensitive to the charge fluctuations at the degeneracy point than at the off-degeneracy point. Thus, the dephasing due to the charge fluctuations is much weaker at the degeneracy point.

##### C. Charge echo

The situation in the free-induction-decay experiment is very similar in the case of NMR. Although we measured

a single two-level system, because of the time-ensemble averaging process in the measurement scheme, even slow fluctuations of qubit energy levels could contribute to the apparent decoherence as an inhomogeneity in the ensemble. In NMR, a large ensemble of same spins is measured at the same time as magnetization of the sample. If there is an inhomogeneity in the applied magnetic field, the spins precess at different velocities, and the oscillation in the averaged signal decays rapidly. However, in NMR, a clever technique called spin echo is known for suppressing the effect of inhomogeneity [39,40]. We have adapted the idea to our charge qubit and called it charge echo [18]. In the echo experiment, three pulses were applied on the qubit. The second pulse inserted between two  $(\pi/2)_x$ -pulses works as a  $(\pi)_x$ -pulse, which flips a sign of the phase. As a result, the phase evolutions in the first ( $t_{d1}$ ) and the second part ( $t_{d2}$ ) of the delay time are cancelled each other, and the effect of the fluctuations is erased as long as the fluctuations of the energy level is slow compared to the delay time. In fact, the echo signal was observed only when the two parts of the delay time are almost same (Fig. 3(b)), and the echo signal survived up to about 5 ns when we changed the delay time (Fig. 5). The result showed that the echo technique recovers a part of the coherence buried in the inhomogeneity and, at the same time, that the low-frequency part of the energy-level fluctuations was dominant in the apparent decoherence.

#### D. Decoherence sources

As a source of the low-frequency energy-level fluctuations, the so-called background charge fluctuations are considered.  $1/f$  charge noise is very common in single-electron devices and is usually attributed to random motions of charge impurities around the device either in the substrate or the junction barrier, although the exact mechanism is not known [41,42]. Indeed, the sample studied here also showed  $1/f$  noise when the dc current through the probe junction was measured. Because of the narrow bandwidth of the measurement setup, the  $1/f$  noise spectrum could be measured only below  $\sim 100$  Hz. However, here we assume that the  $1/f$  noise spectrum extended to infinite frequency and estimate the dephasing of the qubit due to the noise [43]. Under an assumption of gaussian fluctuations [44], the dephasing factor as a function of the delay time ( $t_d = \tau$ ) in the free-induction decay is given as

$$\langle \exp i\varphi(\tau) \rangle = \exp \left[ -\frac{1}{2\hbar^2} \int_{\omega_m}^{\infty} d\omega S_{\Delta E}(\omega) \left( \frac{\sin(\omega\tau/2)}{\omega/2} \right)^2 \right]. \quad (2)$$

Here,  $S_{\Delta E}(\omega) = (4E_C/e)^2 S_Q(\omega)$  is a power-spectrum density of the energy-level fluctuations related to that

of  $1/f$  charge fluctuations  $S_Q(\omega) \equiv \alpha/\omega$ , and  $\omega_m$  is a low-frequency cutoff corresponding to a finite averaging time (20 ms) in the experiment. Similarly, for the echo experiment with  $t_{d1} = t_{d2} = \tau/2$ , we obtain

$$\exp \left[ -\frac{1}{2\hbar^2} \int_{\omega_m}^{\infty} d\omega S_{\Delta E}(\omega) \left( \frac{\sin^2(\omega\tau/4)}{\omega/4} \right)^2 \right]. \quad (3)$$

The effect of the echo technique is clearly seen when we compare the integrands of the two formulas for  $S_{\Delta E}(\omega) \propto 1/\omega$  as shown in Fig. 4. The low-frequency part contribution in the free-induction decay experiment is partly cancelled in the echo experiment, but not completely.

Coupling of the qubit to the external circuit also causes decoherence [8]. The qubit can emit or absorb a photon (to be precise, a bosonic excitation in electromagnetic modes in the circuit) to and from the bosonic degrees of freedom in the electromagnetic environment. In the present case, it takes place as a photon-assisted Cooper-pair tunneling between the box and the reservoir electrodes. Voltage fluctuations in the environment can also dephase the qubit in an elastic manner. However, the estimation of those effects gives a time scale of about 100 ns, which is much longer than that of our experiment. Those mechanisms seem not dominant in the present case.

As we have discussed in Sec. III, the readout process using quasiparticle tunneling through the probe junction stochastically disturbs the coherence of the qubit at a rate  $\Gamma_{qp1}$ . Thus, we expect decrease of the average coherence as  $\exp(-\Gamma_{qp1}\tau)$ . In the present experiment,  $\Gamma_{qp1}$  was estimated to be about  $(8 \text{ ns})^{-1}$ .

In Fig. 5, all the estimations of decoherence as well as the experimental result are plotted. The decoherence due to the  $1/f$  background charge fluctuations gives a gaussian decay [49]. If we fit the echo-experiment result, we obtain  $\alpha = (0.64 \times 10^{-3}e)^2$ . The number is quite close to what we observed in the low-frequency noise measurement ( $\alpha = (1.3 \times 10^{-3}e)^2$ ) and a typical numbers given in the literature ( $\alpha = (0.3 \times 10^{-3}e)^2$ ). Considering the uncertainty in the noise spectrum in the frequency range not accessible by the direct noise measurement, it is difficult to compare the estimation directly with the data. However, the result suggests that the background charge fluctuations could be a dominant source of the decoherence of our qubit [45].

#### VI. CONCLUDING REMARKS

In our experiments, coherent manipulations of charge-number states in a Cooper-pair box were demonstrated. Quantum state of a single qubit was controlled by voltage pulses applied on the gate electrode. A few successive rotations of the state vector around different axis

were used to conduct a free-induction-decay experiment and a charge-echo experiment. At the same time, however, problems for future experiments, for example, to study the quantum coherent properties and entanglement among multiple qubits, were also clarified. The decoherence time at the off-degeneracy point seems to be limited by background charge fluctuations in the device. It was shown that a part of the apparent decoherence due to the low-frequency fluctuations can be cancelled by charge-echo technique. However, as long as we use the present design of the qubit and the same operation scheme, a long decoherence time for doing many-step quantum-state operations would only be possible if we could suppress the background charge noise drastically. Moreover, the readout scheme we used was far from ideal. The detector was permanently coupled to the qubit, and an ensemble average over many readouts had to be taken because of the small output signal. A detector, which can be disconnected from qubits during the coherent control and has a single-shot measurement resolution, is desired.

Recently experimental progresses on those issues have been reported. In Yale and Chalmers, RF-SET (radio-frequency single-electron transistor) has been developed and proposed as a detector for a Cooper-pair box [46–48]. Saclay group proposed and demonstrated a use of switching current measurement across a Cooper-pair box to read out the qubit state [6,49,50]. In both schemes, the detector can be well decoupled from the qubit when it is off, and is expected to have a single-shot measurement resolution when it is switched on. In the latter readout scheme, the readout axis is  $x$ -axis [51]. Thus, the readout works at the degeneracy point where the qubit basis is along  $x$ -axis and the energy levels are immune to charge fluctuations in the first order. In the experiment in Saclay, the qubit was operated at the degeneracy point and showed a long decoherence time of about 500 ns [6]. The ratio  $E_J/E_C$  was much larger than that of our experiment. This also helped to increase the immunity of the qubit to charge fluctuations and made a long decoherence time available even in the presence of background charge noise.

In view of those progresses and the better understandings of the physics behind, it seems that in the near future we will see more complicated quantum-state manipulations among multiple qubits in the Josephson-junction charge-qubit system [52]. Demonstrations of a single-shot measurement, a two-qubit gate operation, and measurement of quantum correlation, or entanglement, would be the next important steps. At the same time, more engineering issues such as a precision of the control and a cross-talk between different control gates would become important for the many-step manipulations.

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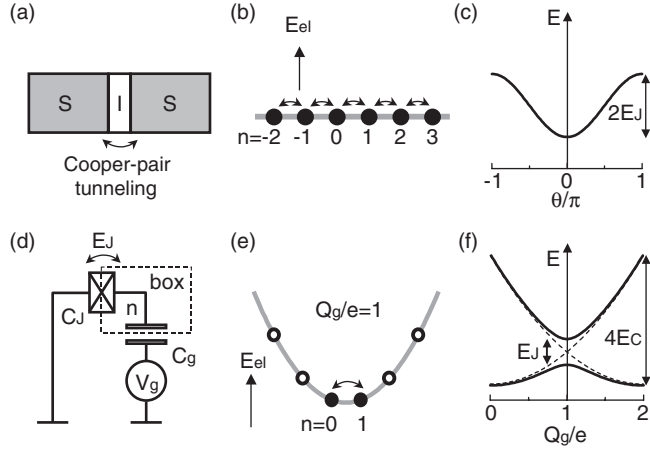


FIG. 1. Josephson effect and a Cooper-pair box. (a) Josephson junction. (b) Degenerate charge-number states. (c) Energy band in the phase space. (d) Schematic circuit diagram of a Cooper-pair box. The rectangle with a cross represents a Josephson junction. (e) Electrostatic energy of the charge states. (f) Energy bands in the charge space.

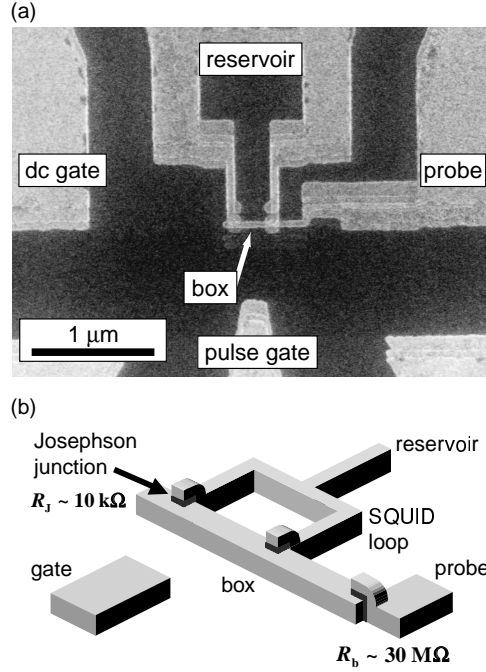


FIG. 2. (a) Scanning-electron micrograph and (b) schematic of a Cooper-pair box with an additional probe electrode.



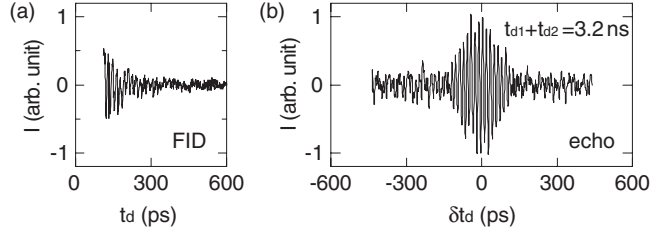


FIG. 3. Charge-echo experiment: (a) Free-induction decay (FID) signal vs delay time  $t_d$ . (b) Echo signal vs. difference  $\delta t_d$  between the first ( $t_{d1}$ ) and the second delay time ( $t_{d2}$ ).

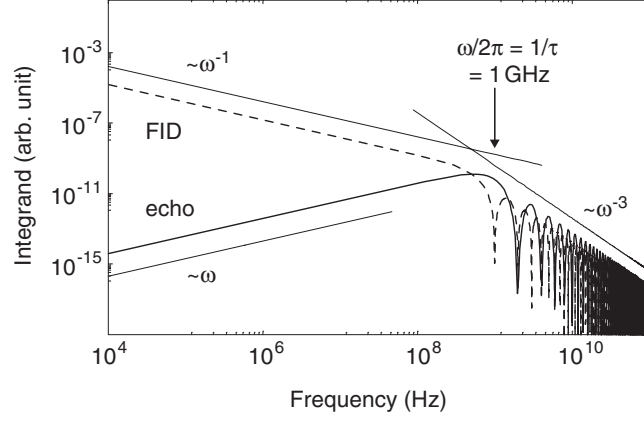


FIG. 4. Integrands in the dephasing formula for the free-induction-decay (FID) experiment (Eq. (2); dashed curve) and for the echo experiment (Eq. (3); solid curve).  $\tau$  is set to be 1 ns.

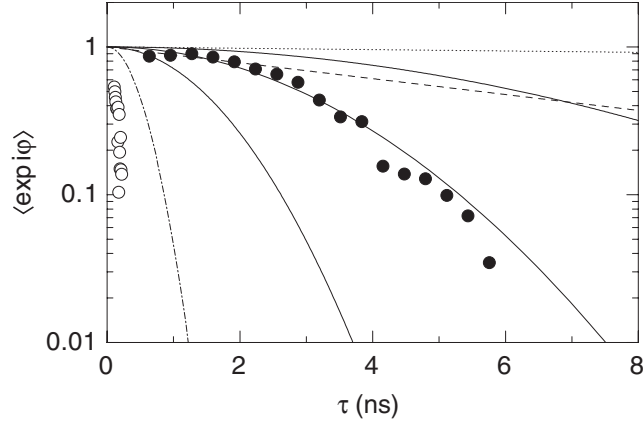


FIG. 5. Decay of the normalized amplitude of the echo (filled circles) and the free-induction-decay signals (open circles). Estimated decoherence due to electromagnetic environment (dotted line) and the readout process (dashed) are shown. The solid lines are estimations for the echo experiment in the presence of  $1/f$  background charge noise with  $\alpha = ((0.3, 0.64, 1.3) \times 10^{-3} e)^2$  from top to bottom, respectively. The dash-dotted curve is the calculated dephasing in the free-induction decay for  $\alpha = (0.64 \times 10^{-3} e)^2$ .