Charge Relaxation in a Single-Electron Si/SiGe Double Quantum Dot

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We measure the interdot charge relaxation time $T_1$ of a single electron trapped in an accumulation mode Si/SiGe double quantum dot. The energy level structure of the charge qubit is determined using photon-assisted tunneling, which reveals the presence of a low-lying excited state. We systematically measure $T_1$ as a function of detuning and interdot tunnel coupling and show that it is tunable over four orders of magnitude, with a maximum of 45 μs for our device configuration.

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Semiconductor quantum dots have been widely used as probes of fundamental quantum physics and to implement charge and spin qubits [1,2]. Coherent manipulation and two-qubit entanglement have been demonstrated in GaAs double quantum dots (DQDs) and error correction techniques such as dynamic decoupling have been employed to suppress decoherence [3–6]. As an alternative host material, Si holds promise for ultrafast spin qubits due to weak spin-orbit coupling, a centrosymmetric lattice (no piezoelectric phonon coupling), and an established route to isotopic purification [7–11]. Spin lifetimes of 6 s have been measured in Si and isotopically purified $^{28}$Si crystals can support spin coherence times as long as 4 s [7,12].

While Si closely approximates a “semiconductor vacuum” for electron spins, its electronic band structure leads to potential complications that are absent in the conventional GaAs/AlGaAs two-dimensional electron gas (2DEG) system [13]. First, the ~3 times larger effective mass of electrons in Si requires depletion gate patterns to be scaled down significantly in order to achieve orbital level spacings comparable to those obtained in GaAs. Second, the band structure of bulk Si consists of six degenerate valleys, which introduces an additional decoherence pathway [14]. Valley degeneracy is partially lifted by uniaxial strain in a Si/SiGe heterostructure [15]. However, the energy splitting between the lowest two valleys is highly sensitive to device specifics, such as step edges in the quantum well [11,16,17]. Detailed measurements of the low-lying energy level structure, and the time scales that govern energy relaxation between these levels, are therefore needed in Si quantum dots [18].

In this Letter, we systematically measure the interdot relaxation time $T_1$ of a single electron trapped in a Si DQD as a function of detuning $\varepsilon$ and interdot tunnel coupling $t_c$. We demonstrate a four order of magnitude variation in $T_1$ using a single depletion gate and obtain $T_1 = 45 \mu s$ for weak interdot tunnel couplings [19]. We also use photon-assisted tunneling (PAT) to probe the energy level structure of the single-electron system, demonstrating spectroscopy with an energy resolution of ~1 μeV. In contrast with single-electron GaAs dots, we observe low-lying excited states ~55 μeV above the ground state, an energy scale that is consistent with previously measured valley splittings [11,16].

Measurements are performed on an accumulation mode Si/SiGe DQD. We apply a top gate voltage $V_T = 2$ V to accumulate carriers in a Si quantum well located ~40 nm below the surface of the wafer [see Fig. 1(a)]. The resulting 2DEG has an electron density of $4 \times 10^{11}$ cm$^{-2}$ and a mobility of 70 000 cm$^2$/V s. A 100 nm thick layer of Al$_2$O$_3$ separates the top gate from the depletion gates, which are arranged to define a DQD and a single-dot charge sensor.

We first demonstrate single-electron occupancy using radio frequency (rf) reflectometry [20]. A single quantum dot is coupled to a resonant circuit with resonance frequency $f_c = 431.8$ MHz [see Fig. 1(a)] and used as a high sensitivity charge detector [21]. The reflected amplitude $A$ is a sensitive function of the conductance of the single-dot sensor $g_Q$, which is modulated by charge transitions in the DQD. We map out the DQD charge stability diagram in Fig. 1(b) by plotting the numerical derivative $dA/dV_L$ as a function of $V_L$ and $V_R$. No charging transitions are observed in the lower left corner of the charge stability diagram, indicating that the DQD has been completely emptied of free electrons. We identify this charge configuration as (0,0), where $(N_L, N_R)$ indicates the number of electrons in the left and right dots.

The device is operated as a single-electron charge qubit near the (1,0) ↔ (0,1) interdot charge transition. Charge dynamics are governed by the Hamiltonian $H = (\varepsilon/2)\sigma_z + t_c\sigma_x$, where $\sigma_i$ are the Pauli matrices. We demonstrate tunable interdot tunnel coupling in the single-electron regime by measuring the left dot occupation $P_{(1,0)}$ as a function of detuning [see the inset of Fig. 1(b)] [1,3,22]. Qubit occupation is described by

$$P_{(1,0)} = \frac{1}{2} \left( 1 + \cos \left( \frac{2 \pi}{T_1} \right) \right).$$

$$T_1 = \frac{\hbar}{2} |t_c| \frac{1}{\Delta E}.$$
we observe a small resonance away from the interdot charge transition, suggesting the existence of a low-lying excited state in one of the dots. In particular, we observe a small resonance bias triangles is not a smooth function of gate voltage. 

\[ P_{(1,0)} = \frac{1}{2} \left[ 1 - \frac{e}{\Omega} \tanh \left( \frac{\Omega}{2k_B T_e} \right) \right] \]  

where \( k_B \) is Boltzmann’s constant, \( T_e \sim 100 \text{ mK} \) is the electron temperature, and \( \Omega = \sqrt{E^2 + 4T_e^2} \) is the qubit energy splitting [19,23,24]. With \( V_N = 225 \text{ mV} \), the interdot charge transition is thermally broadened as \( 2T_e < k_B T_e \). Increasing \( T_e \) by adjusting \( V_N \) leads to further broadening of the interdot transition. For \( V_N = 290, 300, \) and \( 310 \text{ mV} \) we extract \( 2T_e = 3.8, 5.9, \) and \( 9.0 \text{ GHz} \) by fitting the data to Eq. (1). These results show that the interdot tunnel coupling can be sensitively tuned in the single-electron regime in Si.

We investigate the DQD energy level structure in Fig. 2(a), where we plot the current \( I \) as a function of \( V_L \) and \( V_R \) with a fixed source-drain bias \( V_{SD} = 700 \text{ µeV} \) [25]. In contrast with GaAs devices, the current in the finite bias triangles is not a smooth function of gate voltage. In particular, we observe a small resonance \( \sim 60 \text{ µeV} \) away from the interdot charge transition, suggesting the existence of a low-lying excited state in one of the dots. In a few electron GaAs DQD, orbital excited states are typically several meV higher in energy than the ground state [26].

Higher energy resolution is obtained using PAT spectroscopy, in which microwaves drive charge transitions when the photon energy matches the qubit splitting \( hf = \Omega \), where \( f \) is the photon frequency and \( h \) is Planck’s constant. PAT transitions are directly observed as deviations from the ground state occupation in measurements of \( P_{(1,0)} \) as a function of detuning [compare Fig. 2(b) and the inset of Fig. 1(b)]. For \( f \lesssim 15 \text{ GHz} \), the PAT peaks are symmetric around \( \epsilon = 0 \) and shift to larger detuning with increasing photon energy, consistent with a simple two level interpretation [19,27]. However, for \( f \gtrsim 15 \text{ GHz} \), an additional PAT peak emerges at negative detuning and is not accompanied by a corresponding PAT peak at positive detuning. Figure 2(c) shows the extracted transition frequencies as a function of detuning.

The data are fit using a three level Hamiltonian that includes the left dot ground state (1,0)\(_L\), the right dot ground state (0,1)\(_R\), and a right dot excited state (0,1)\(_e\), as sketched in the inset of Fig. 2(b) (see the Supplemental Material [28]). We obtain best fit values of \( \tau \approx 1.9 \text{ GHz} \) and \( \Delta = 55 \text{ µeV} \), consistent with the data shown in the inset of Figs. 1(b) and 2(a). Within the \( \sim 1 \text{ µeV} \) resolution of our measurement, we do not observe anticrossings associated with (0,1)\(_e\). The energy eigenstates obtained from the model are plotted as a function of detuning in Fig. 2(a), where we plot the current \( I \) as a function of \( V_L \) and \( V_R \) near the (1,0) \( \leftrightarrow (0,1) \) charge transition. A cut through the finite bias triangle (right) indicates the presence of a low-lying excited state. (b) \( P_{(1,0)} \) plotted as a function of detuning \( \epsilon \) for different excitation frequencies. For \( f \gtrsim 15 \text{ GHz} \), a new PAT peak emerges (gray arrow) corresponding to the (1,0)\(_g \leftrightarrow (0,1) \) transition. The appearance of this PAT peak is accompanied by the suppression of the (1,0)\(_g \leftrightarrow (0,1) \) PAT peak (black arrow) at positive detuning. (c) Transition frequencies as a function of detuning and (d) energy level diagram extracted from data in (c). The data in (c) are best fit with an interdot tunnel coupling \( \tau = 1.9 \text{ GHz} \) and an excited state energy \( \Delta = 55 \text{ µeV} \).
Fig. 2(d). For comparison, an excited state is observed in the left quantum dot in a second device (Device 2), with 
\( \Delta = 64 \mu eV \) (see the Supplemental Material [28]). For both devices, the excited state energy is highly sensitive to \( V_N \) and \( V_C \), suggesting that it is not purely orbital in origin [17].

Several additional features observed in the data are explained by the three level model. The \((0, 1)_g \rightarrow (0, 1)_e\) intradot charge transition [see the dotted line in Fig. 2(c)] is not visible since the charge detector is only sensitive to interdot charge transitions. We also note that the \((0, 1)_e \rightarrow (1, 0)_g\) PAT peak is not visible at positive detuning. At low temperatures, the qubit population resides in the ground state \((0, 1)_g\), preventing microwave transitions from \((0, 1)_e\) to \((1, 0)_g\). Finally, the \((0, 1)_g \rightarrow (1, 0)_g\) PAT peak is suppressed when \( \varepsilon > \Delta \) due to population trapping in \((0, 1)_e\).

We measure the interdot charge relaxation time \( T_1 \) by applying microwaves to \( V_L \) with a 50\% duty cycle and varying the pulse period \( \tau \) [see Fig. 3(a)]. We focus on the \((1, 0)_g \rightarrow (0, 1)_g\) transition at negative detuning, where the high energy state \((0, 1)_g\) is not populated. Simulations of \( P_{(1,0)} \) as a function of time \( t \) for \( \tau = 1 \mu s \) are shown in Fig. 3(a) for three realistic values of \( T_1 \). During the first half of the pulse cycle, microwaves drive the \((1, 0)_g \rightarrow (0, 1)_g\) charge transition, resulting in an average \( P_{(1,0)} = 0.5\). The microwave excitation is then turned off, leading to charge relaxation during the second half of the pulse cycle, with \( P_{(1,0)} \) approaching 1 on a time scale set by \( T_1 \).

In the inset of Fig. 3(b), we plot \( P_{(1,0)} \) as a function of detuning for \( \tau = 10 \) ns and \( \tau = 100 \) \( \mu s \). As expected, the PAT peak is smaller for longer periods due to charge relaxation. Specifically, in the limit \( \tau \ll T_1 \), there is not sufficient time for relaxation to occur during the second half of the pulse cycle, leading to a time averaged value of \( P_{(1,0)} = 0.5\). In contrast, in the limit \( \tau \gg T_1 \), relaxation happens quickly, leaving \( P_{(1,0)} = 1\) for the majority of the second half of the pulse cycle. Due to experimental limitations, such as frequency dependent attenuation in the coax lines and finite pulse rise times at small \( \tau \), we are unable to drive the transitions to saturation for some device configurations. To extract \( T_1 \) we therefore fit the raw \( P_{(1,0)} \) data as a function of \( \tau \) to the form

\[
P_{(1,0)} = P_{\text{max}} + (P_{\text{min}} - P_{\text{max}}) \frac{2T_1(1 - e^{-\tau/(2T_1)})}{\tau},
\]

where \( P_{\text{max}} \) and \( P_{\text{min}} \) account for the limited visibility of the PAT peaks (see Ref. [19] and the Supplemental Material [28]). Extracted \( T_1 \) values are insensitive to the rescaling of the data via \( P_{\text{max}} \) and \( P_{\text{min}} \).

The interdot charge relaxation rate is strongly dependent on the interdot tunnel coupling. This variation is directly visible in the data shown in Fig. 3(b) for \( V_N = 225, 250, \) and 265 mV. To facilitate a direct comparison of the data, we plot the normalized electron occupation \( P_{(1,0)}' = 0.5 + 0.25(P_{(1,0)} - P_{\text{min}})/(P_{\text{max}} - P_{\text{min}}) \), using the values of \( P_{\text{min}} \) and \( P_{\text{max}} \) extracted from fits to Eq. (2) (see the Supplemental Material [28]). In Fig. 3(c), we plot \( T_1 \) over a wide range of \( V_N \) for two different excitation frequencies. We see a longer characteristic relaxation time for larger interdot barrier heights, with a maximum observed value of 45 \( \mu s \). The same overall trend is observed in data from Device 2 [see Fig. 3(d)] where the interdot tunnel coupling was tuned using \( V_C \). Interdot tunnel coupling is only measurable in charge sensing when \( 2t_c > k_B T_c \) [23].

For Device 1 [see Fig. 3(c)] we obtain \( 2t_c = 2.4, 3.8, \) and 5.9 GHz for \( V_N = 280, 290, \) and 300 mV and for Device 2 [see Fig. 3(d)] we obtain \( 2t_c = 3.2 \) GHz for \( V_C = -325 \) mV.

The detuning dependence of \( T_1 \) is investigated in Fig. 4(a), where we plot \( T_1 \) as a function of \( f \propto \Omega \) for the \((0, 1)_g \rightarrow (1, 0)_g\) transition [3, 22]. Data are taken at \( f = 12.3, 19.5, 25.9, \) and 30.0 GHz, as indicated by the arrows in the energy level diagram in the upper panel of Fig. 4(a). Our data indicate that \( T_1 \) increases weakly as a function of detuning for the range of frequencies accessible in our cryostat.

To further investigate the excited state, we measure \( T_1 \) for the \((0, 1)_g \rightarrow (1, 0)_g\) and \((0, 1)_e \rightarrow (1, 0)_g\) relaxation processes at the same values of \( f \) [see the bottom panel of Fig. 4(b)]. In contrast with the \((0, 1)_g \rightarrow (1, 0)_g\) relaxation process, \((0, 1)_e\) can relax via two distinct pathways [see the top panel of Fig. 4(b)]. The first relaxation process is a direct transition from \((0, 1)_e \rightarrow (1, 0)_g\) with a rate \( \Gamma_e\), while
the second pathway proceeds via intradot charge relaxation to $(0,1)_g$ with a rate $\Gamma_g$, followed by an interdot transition to $(1,0)_g$ with a rate $\Gamma^*_g$. We find that the $(0,1)_e \rightarrow (1,0)_g$ relaxation is faster than the $(0,1)_g \rightarrow (1,0)_g$ relaxation for the same energy splitting. The shorter excited state lifetime is consistent with either a fast direct relaxation rate $\Gamma_e^*$ or a fast intradot relaxation followed by an interdot transition. Neglecting valley physics, one would expect $\Gamma_e > \Gamma_g$ due to the more extended excited state orbital wave function. Considering valley physics, and assuming that direct interdot charge relaxation is limited by the intervalley relaxation rate, then $\Gamma_e < \Gamma_g$. The role of valley states is somewhat debated and the nature of the right dot excited state is unclear [18]. Therefore, to obtain a rough estimate of $\Gamma_i$, we assume $\Gamma_i = \Gamma_e$ (since the level detuning is the same). Taking the measured excited state $T_i = 55 \text{ ns}$ at $f = 210 \text{ GHz}$, we obtain a lower bound estimate of $\Gamma_i = 1.5 \times 10^7 \text{ s}^{-1}$ [29].

We modify the results of Raith et al. to allow the calculation of phonon mediated charge relaxation rates considering only intravalley relaxation in the far detuned limit ($|\epsilon| > \tau_e$), assuming Gaussian wave functions for the nonhybridized charge states, with dot radius $a$ and dot separation $2d$ (see the Supplemental Material [28] and Ref. [30]). The electron-phonon coupling Hamiltonian in a Si quantum well takes the form

$$H_{e \rightarrow ph} = \sum_{Q,\lambda} \frac{\hbar |Q|}{2 \rho V c^2_{\lambda}} D_Q^\lambda (a_Q^\lambda e^{iQ \cdot r} - a_Q^\dagger e^{-iQ \cdot r}),$$

where

$$D_Q^\lambda = (\Xi_d \hat{e}_Q^\dagger \cdot \hat{Q} + \Xi_u \hat{e}_Q^\dagger \hat{Q}_d).$$

Here $a_{Q,\lambda}$ ($a_{Q,\lambda}^\dagger$) is the annihilation (creation) operator for phonons belonging to branch $\lambda$ ($\lambda = TA1, TA2$ for transverse phonons and $\lambda = LA$ for longitudinal phonons) with wave vector $Q$, and speed of sound in Si $c_{\lambda}$. $V$ is the volume of the Si quantum well layer and $\rho$ is the density of Si. $\Xi_u$ and $\Xi_d$ are the shear and dilatation deformation potential constants, and $\hat{e}_Q$ and $\hat{Q}$ are the phonon unit polarization vector and the phonon unit wave vector [30]. Using realistic parameters, $T_i$ values calculated in this model are in order of magnitude agreement with our data (see the Supplemental Material [28]). However, the predictions are exponentially sensitive to $a$ and $d$, quantities that are difficult to accurately determine. Moreover, the model predicts a relaxation rate that increases with energy splitting for the range of detunings accessed in our experiment, following the power law $\Gamma_i = 1/T_i \approx \Omega^3$, whereas we observe a rate that decreases weakly with increasing detuning (see the Supplemental Material [28]). This discrepancy may be due to a detuning dependent $\tau_e$ or contributions from other relaxation channels, such as charge noise [31].

In summary, we have measured charge relaxation times in a single-electron Si/SiGe DQD, demonstrating a four order of magnitude variation of $T_i$ with gate voltage. Energy level spectroscopy indicates the presence of a low-lying excited state. From the estimated dot radius $a \approx 38 \text{ nm}$, we expect orbital level spacings on the order of $1 \text{ meV}$, a factor of 18 larger than the value obtained from PAT spectroscopy ($\Delta = 55 \mu \text{eV}$) [32]. This suggests that the low-lying excited state is a valley-orbit mixed state [33].

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