A two-qubit logic gate in silicon

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Quantum computation requires qubits that can be coupled in a scalable manner, together with universal and high-fidelity one- and two-qubit logic gates. Many physical realizations of qubits exist, including single photons, trapped ions, superconducting circuits, single defects or atoms in diamond and silicon, and semiconductor quantum dots, with single-qubit fidelities that exceed the stringent thresholds required for fault-tolerant quantum computing. Despite this, high-fidelity two-qubit gates in the solid state that can be manufactured using standard lithographic techniques have so far been limited to superconducting qubits, owing to the difficulties of coupling qubits and dephasing in semiconductor systems. Here we present a two-qubit logic gate, which uses single spins in isotopically enriched silicon and is realized by performing single- and two-qubit operations in a quantum dot system using the exchange interaction, as envisaged in the Loss–DiVincenzo proposal. We realize CNOT gates via controlled-phase operations combined with single-qubit operations. Direct gate-voltage control provides single-qubit addressability, together with a switchable exchange interaction that is used in the two-qubit controlled-phase gate. By independently reading out both qubits, we measure clear anticorrelations in the two-spin probabilities of the CNOT gate.

Figure 1c shows the stability diagram of the double-quantum-dot system with charge occupancy (N1, N2). The charge transitions of quantum dots D1 and D2, which are underneath gates G1 and G2, respectively, are distinguished by their gate voltage dependence and their capacitive coupling to the SET. We define qubit Q1 by loading a single electron into D1, so that N1 = 1; similarly for qubit Q2, we have N2 = 1.

To characterize the individual qubits, we bias the gate voltages such that the tunnel time of the respective qubit to the reservoir is approximately 100 µs and both qubits are measured in the (1, 1) charge state. Clear Rabi oscillations are observed as a function of microwave pulse time τ, for both qubits, as shown in Fig. 1d. Q1 has a dephasing time of T2 = 120 µs, with a coherence time T1 that can be extended up to 28 ms using CPMG (Carr–Purcell–Meiboom–Gill) pulses, and Q2 has a dephasing time of T2 = 61 µs (see Supplementary Information section 3).

We couple the qubits via the exchange interaction, as discussed in ref. 2, with an exchange coupling that is electrically controlled via the detuning energy ε (see Fig. 2a, b). We control the system in the (1, 1) region, and read out Q2 at the (1, 1)–(0, 1) transition and Q1 at the (0, 1)–(0, 0) transition.

The presence of a sharp interface and large perpendicular electric fields increases the energy of all excited states, and allows us to consider only the lowest five energy states (see also Supplementary Information section 2). We can consequently describe the system in the rotating-wave approximation in the basis [Q2, Q1], :ψ = [1, 1], [1, 1], [1, 1], [1, 1], [0, 2], with the effective Hamiltonian

\[ H = \begin{pmatrix} E_Z - \nu & \Omega & 0 & 0 \\ \Omega & \delta E_Z / 2 & 0 & \Omega \\ 0 & 0 & -E_Z + \nu & 0 \\ 0 & \Omega & -\delta E_Z / 2 & -t_0 \end{pmatrix} \]

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The anticrossing, as shown in Fig. 2c, results in different capacitive coupling, such that the individual dots can be easily distinguished. The tunnel coupling of the fourth transition ($N_1 = 3 \rightarrow 4$) of $D_3$ is relatively weak, which is due to valley and spin filling, because there is only one state in the lowest orbital that can be occupied. $Q_1$ and $Q_2$ are realized by depleting $D_1$ and $D_2$ to the last electron.

Figure 1 | Silicon two-qubit logic device, incorporating SET read-out and selective qubit control. a, b, Schematic (a) and scanning electron microscope coloured image (b) of the device. The quantum dot structure (labels $G_1$ and $G_{1,4}$) can be operated as a single or double quantum dot by appropriate biasing of gate electrodes $G_1$–$G_4$, where we choose here to confine the dots $D_{1,2}$ underneath gates $G_{1,2}$, respectively. The confinement gate $G_C$ runs underneath $G_1$–$G_2$ and confines the quantum dot on all sides except on the reservoir (R) side. Qubit operation is achieved via an ac current $I_{ac}$ through the ESR line, resulting in an ac magnetic field $B_{ac}$. c, Stability diagram of the double quantum dot obtained by monitoring the current $I_{SET}$ through the capacitively coupled SET. The numbers in parentheses are the change occupancies of $D_{2,1}$: ($N_2$, $N_1$).

The difference in distance to the SET results in different capacitive coupling, such that the individual dots can be easily distinguished. The tunnel coupling of the fourth transition ($N_1 = 3 \rightarrow 4$) of $D_3$ is relatively weak, which is due to valley and spin filling, because there is only one state in the lowest orbital that can be occupied. $Q_1$ and $Q_2$ are realized by depleting $D_1$ and $D_2$ to the last electron.

d, The quantum dot qubits can be individually controlled by electrically tuning the ESR resonance frequency using the Stark shift. Clear Rabi oscillations for both qubits are observed. All measurements were performed in a dilution refrigerator with base temperature $T = 50$ mK and a dc magnetic field of strength $B_0 = 1.4$ T.

where $E_Z$ is the mean Zeeman energy, $\delta E_Z$ is the difference in Zeeman energy between the dots, $\Omega$ is the Rabi frequency, $\nu$ is the microwave frequency and $\tau_0$ is the tunnel coupling; for simplicity we have scaled the system such that $h = 1$. In the experiments, we control $\epsilon$ by fast pulsing only on $G_1$. The read-out on $Q_2$ (R2) and control (C) bias points are depicted in Fig. 2a; we pulse close to the $(1, 1)$–$(1, 2)$ transition, which has the same energy level structure as the $(1, 1)$–$(0, 2)$ transition, shown in Fig. 2b. Single-qubit operations are realized with Rabi frequency $\Omega_r$, by matching $\nu$ to the resonance frequency of one of the qubits. The presence of exchange coupling between the qubits alters the Zeeman levels as shown in Fig. 2b, where the finite coupling between the qubits causes an anticrossing between the $(0, 2)$ and $(1, 1)$ states. We experimentally map out the energy levels in the vicinity of the anticrossing, as shown in Fig. 2c, d.

In Fig. 2c, we have initialized $Q_1$ and $Q_2$ to spin down and, by applying a $\pi$-pulse to $Q_1$ ($\pi_{X,Q_1}$), we map out the resonance frequency of $Q_2$ as a function of detuning. We measure exchange couplings of more than 10 MHz, above which $T_2^*$ of $Q_2$ becomes shorter than the $\pi$-pulse time $\tau_\pi = 1.5$ $\mu$s and so spin flips cannot occur.

Initialization of antiparallel spin states is possible by pulsing to the $(1, 2)$ and returning to the $(1, 1)$ charge states (labelled $I_{AP}$ in Fig. 2a).

In this sequence, an electron tunnels from $D_2$ to $D_1$ (I in Fig. 2d) followed by an electron tunnelling from the reservoir R to $D_2$ (II in Fig. 2d). After returning to $(1, 1)$, the electron from $D_2$ tunnels back to R (III in Fig. 2d), and one of the two electrons on $D_1$, which are in a singlet state, tunnels to $D_2$ (IV in Fig. 2d). With this initialization into an antiparallel spin state, when we apply a microwave pulse on $Q_2$, the spin-up fraction $f_1$ approaches $1/2$, except when the microwave frequency matches a resonance frequency of $Q_2$. Owing to the finite exchange interaction, there are two resonance frequencies. The lower frequency rotates the antiparallel state towards a combination of $(\uparrow, \downarrow)$ and $(\downarrow, \uparrow)$, where $Q_2$ always ends up as spin down. The higher frequency rotates the antiparallel state towards a combination of $(\downarrow, \downarrow)$ and $(\uparrow, \uparrow)$, where $Q_2$ always ends up as spin up. The results are depicted in Fig. 2d, which shows a decrease in $f_1$ at the lower branch and an increase of $f_1$ at the upper branch, demonstrating an exchange spin funnel, where both branches are visible (see Supplementary Information section 4 for further details). The two-qubit gate is conveniently realized using a quantum dot CZ gate (see Supplementary Information section 6 for theoretical details). This approach allows individual control over the qubits in the absence of interaction and its associated noise, while using the coupling to perform two-qubit operations with a frequency that can be much higher than the single-qubit Rabi rotation frequency.

As described by equation (1) and depicted in Fig. 2b, changing $\epsilon$ modifies the qubit resonance frequencies of $Q_1$ and $Q_2$ and introduces an effective detuning frequency $\nu_{1,4}(1, 1)$, such that one qubit acquires a time-integrated phase shift $\phi_{1,4}(1, 1)$ that depends on the $\overrightarrow{Z}$ component of the spin state of the other qubit, and vice versa. The exchange coupling and $\nu_{1,4}(1, 1)$ are maximized at the anticrossing $\epsilon = U_2$; see Fig. 2b. When
Figure 2 | Exchange spin funnel. a, Close up of the operation regime of the (1, 1)–(0, 2) charge states. We lowered the R–D2 coupling so that the tunnelling time is approximately 100 μs, matching the qubit experiments. In this range of weak R–D1 coupling, the emptying and filling of D1 is hysteretic with gate voltage, because the mutual charging energy becomes relevant25, as D1 can only tunnel when it aligns in energy with D2. R2 represents the read-out on Q2. IAP represents the antiparallel initialization. b, Schematic of the coupling between Q1 and Q2, using the exchange interaction at the |1, 1⟩–|0, 2⟩ transition. By electrically tuning the g factors of Q1 and Q2, we control the individual qubit resonance frequencies over 10 MHz. Here, we tune to a frequency difference of 40 MHz (the difference is exaggerated in the schematic for clarity) for individual qubit control. c, ESR spectrum of the |1⟩, |1⟩–|1⟩, |1⟩ transition as a function of increasing detuning. The data have been offset by a frequency ν0 = 39.14 GHz and the spin-up fractions are normalized for clarity. The dashed lines are fits using equation (1) and assuming h0 = 900 MHz, a Stark shift of 19 MHz V−1, and that the top gates have a lever arm of 0.2 eV V−1. d, As for c, but with an additional pulse of amplitude VAP (see schematic) so that we initialize antiparallel spin states (AP) and observe both the |1⟩, |1⟩–|1⟩ and |1⟩, |1⟩–|1⟩ transitions. The labels I–IV indicate electron tunnelling between R and D1, see text for details.

a CZ operation is performed such that ϕ12 + ϕ21 = π, the operation differs only by an overall phase from the basis CZ gate25. This overall phase can be removed using single-qubit pulses or via voltage pulses exploiting the Stark shift25. To realize a CNOT operation using the CZ gate, a CZ(π) rotation is performed in between two π/2-pulses on Q2 that have a phase difference ϕ12.

Figure 3a shows the spin-up fraction f1 of Q2 after applying a (π/2)X-pulse and a (π/2)Y-pulse on Q2 separated by an interaction time τcz with increasing exchange coupling, set via ε and tuned by the voltage Vcz. Plotting the frequency ν1 as a function of Vcz (Fig. 3b) gives a trend consistent with that observed via ESR mapping, as shown in Fig. 2c, d. The two-qubit dephasing time T2,cz is the free induction decay time of the arrows indicate the axis each data set corresponds to; the colouring of the data corresponds to that in c, indicating Vcz. The inset shows the number of possible CZ rotations NCZ. Although T2,cz decreases with coupling, the number of possible two-qubit rotations continues to increase.
the two-qubit system. At large values of detuning ($\epsilon \rightarrow \infty$) or when the interaction vanishes ($\epsilon \rightarrow 0$), $T_{2,\text{CZ}}^*\text{CZ}$ reduces to the single-qubit Ramsey $T_2$. We obtain $T_{2,\text{CZ}}^*$ by fitting an exponential to the decay of the oscillations in Fig. 3a. These values of $T_{2,\text{CZ}}^*$ are plotted along with the measured $V_{11}$ values in Fig. 3b. We find that the two-qubit dephasing rate $(T_{2,\text{CZ}}^*)^{-1}$ rises in step with the exchange coupling and $V_{11}$, which is to be expected because $\partial V_{11}/\partial V$ also increases with $V_{11}$, meaning that the qubit system becomes increasingly sensitive to electrical noise. Despite this, the total number of oscillations $N_{\text{CZ}} = \tau_{11} T_{2,\text{CZ}}^*$ also increases with $V_{11}$, as shown in Fig. 3b. In Supplementary Information section 7, we show an optimized sequence where $T_{2,\text{CZ}}^* = 8.3 \mu$s and $V_{11} = 3.14$ MHz, such that $N_{\text{CZ}} > 26$.

For all of the experiments described in Figs 2 and 3, we performed read-out only on $Q_2$, owing to its proximity to the reservoir used for spin selective read-out. However, to demonstrate a two-qubit CNOT gate it is desirable to read out the state of both qubits, allowing the observation of non-classical correlations. This requires a more complex pulsing protocol on $Q_1$ (Fig. 4a), to first read $Q_2$ at the charge transition (0, 0)--(0, 1), subsequently read and initialize $Q_1$ at the (0, 1)--(1, 1) transition, and initialize $Q_2$ at the (0, 0)--(0, 1) transition. The read-out ($R_{1,2}$) and control ($C$) bias points are shown on the charge stability map in Fig. 4b. Following two-qubit read-out of the previous state, the system is prepared as $|1\rangle$, after which single-qubit rotations are applied to each qubit to prepare any desired initial two-qubit state.

Figure 4c shows the measured states of $Q_1$ and $Q_2$ after applying the CZ gate as a function of $\tau_{2,\text{C}}$ with $Q_1$ initialized to $|1\rangle$ (top panel) and $|0\rangle$ (bottom panel). As expected, the control qubit $Q_1$ (red) is not perturbed when exchange is turned on because it is in a basis state; however, the target qubit $Q_2$ (blue) is initialized to $|1\rangle$ and then a $(\pi/2)_x$-pulse on $Q_2$ and so rotates about the equator of the Bloch sphere when exchange is turned on for time $\tau_{2,\text{C}}$. The strength of the exchange coupling is set so that $Q_2$ rotates about the Bloch sphere at double the frequency for $Q_1 = |0\rangle$ than for $Q_1 = |1\rangle$; this is reflected in the final state of $Q_2$ plotted in Fig. 4c (see Supplementary Information section 8 for further details). A $CZ$ gate is realized at $\tau_{2,\text{C}} = 480$ ns, when $\phi_{11} + \phi_{11} = \pi$, and this is converted to a CNOT gate (the target qubit is

Figure 4 | Two-spin correlations for a two-qubit logic gate. a, Pulsing protocol for two-qubit read-out and single- and two-qubit operations. After read-out of $Q_2$ ($R_2$) and $Q_1$ ($R_1$), we pulse back to ($R_2$) to ensure proper initialization. Individual qubit operations are performed with high $\epsilon$, whereas the CZ operation occurs in the presence of interaction. b, Stability diagram showing the operation regime. c, Spin-up fraction of both qubits after initializing $Q_2$ spin up (top) and spin down (bottom) using a microwave pulse and applying a controlled rotation using $Q_2$, as the target qubit. A CNOT gate is achieved in 480 ns, as indicated by the dotted purple line (see inset for the corresponding Bloch sphere animation). d, Two-spin probabilities as functions of the microwave pulse length on $Q_1$ after applying a CNOT gate (see inset for the corresponding Bloch sphere animation), showing clear anticorrelations between the two qubit spin states. The different plots correspond to different spin states of $Q_1$, as indicated. The black lines correspond to fits based on a CNOT gate, and include the experimental read-out errors (see Supplementary Information section 9). The green dotted lines correspond to the intended maximally entangled states.
flipped when the control qubit is $|\uparrow\rangle$ by applying $(\pi/2)_X$-pulses on $Q_2$ before and after the CZ.

We use the CNOT gate to create an entangled state of $Q_1$ and $Q_2$. To realize this, we initialize the qubits first to the $|\uparrow\rangle, |\downarrow\rangle$ state, then apply a varying microwave pulse time to rotate $Q_1$ into superposition states, with a Rabi time $\tau_{\text{Rabi}} = 2.4 \mu s$, and finally apply the CNOT gate. To demonstrate the CNOT gate, we convert the individual qubit spin-up fractions into two-spin probabilities: Fig. 4d shows the four possible two-spin probabilities. Clear oscillations are observed in the probabilities of the antiparallel states, $P(|\uparrow\rangle, |\downarrow\rangle)$ and $P(|\downarrow\rangle, |\uparrow\rangle)$, whereas these oscillations are almost absent in the probabilities of the parallel states, $P(|\uparrow\rangle, |\uparrow\rangle)$ and $P(|\downarrow\rangle, |\downarrow\rangle)$, thereby demonstrating the anticorrelations expected for the CNOT gate. The hints of oscillations in the symmetric spin states are probably due to read errors (which are included in the fitted line in Fig. 4d, see also Supplementary Information section 9); our current visibilities are not sufficient to demonstrate violation of the Bell inequality.

Future experiments will include improvements to the read-out fidelities, thus facilitating full two-qubit tomography. The qubit control fidelities could be further improved by lowering the sensitivity to electrical noise. Although these silicon qubits represent the smallest scalable two-qubit system reported so far, the complete fabrication process is compatible with standard CMOS (complementary metal–oxide–semiconductor) technology, and is also consistent with current transistor feature sizes, offering the prospect of realizing a large-scale quantum processor using the same silicon manufacturing technologies that have enabled the current information age.

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