A coherent spin–photon interface in silicon

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Electron spins in silicon quantum dots are attractive systems for quantum computing owing to their long coherence times and the promise of rapid scaling of the number of dots in a system using semiconductor fabrication techniques. Although nearest-neighbour exchange coupling of two spins has been demonstrated, the interaction of spins via microwave-frequency photons could enable long-distance spin–spin coupling and connections between arbitrary pairs of qubits (‘all-to-all’ connectivity) in a spin-based quantum processor. Realizing coherent spin–photon coupling is challenging because of the small magnetic–dipole moment of a single spin, which limits magnetic–dipole coupling rates to less than 1 kilohertz. Here we demonstrate strong coupling between a single spin in silicon and a single microwave–frequency photon, with spin–photon coupling rates of more than 10 megahertz. The mechanism that enables the coherent spin–photon interactions is based on spin–charge hybridization in the presence of a magnetic–field gradient. In addition to spin–photon coupling, we demonstrate coherent control and dispersive readout of a single spin. These results open up a direct path to entangling single spins using microwave–frequency photons.

Solid-state electron spins and nuclear spins are quantum mechanical systems that can be almost completely isolated from environmental noise. As a result, they have coherence times as long as hours and so are one of the most promising types of quantum bit (qubit) for constructing a quantum processor1–3. On the other hand, this degree of isolation poses difficulties for the spin–spin interactions that are needed to implement two-qubit gates. So far, most approaches have focused on achieving spin–spin coupling through the exchange interaction or the much weaker dipole–dipole interaction4–6. Among existing classes of spin qubits, electron spins in gate-defined silicon quantum dots have the advantages of scalability due to mature fabrication technologies and low dephasing rates due to isotopic purification7. Currently, silicon quantum dots are capable of supporting fault-tolerant control fidelities for single-qubit gates and high-fidelity two-qubit gates based on exchange6–12. Coupling of spins over long distances has been pursued through the physical displacement of electrons13–16 and through ‘super-exchange’ via an intermediate quantum dot17. The recent demonstration of strong coupling between the charge state of a quantum-dot electron and a single photon has raised the prospect of strong spin–photon coupling, which could enable photon-mediated long-distance spin entanglement18–20. Spin–photon coupling may be achieved by coherently hybridizing spin qubits with photons trapped inside microwave cavities, in a manner similar to cavity quantum electrodynamics with atomic systems and circuit quantum electrodynamics with solid-state qubits19–25. Such an approach, however, is extremely challenging: the small magnetic moment of a single spin leads to magnetic-dipole coupling rates of 10–150 Hz, which are far too slow compared with electron–spin dephasing rates to enable a coherent spin–photon interface25–30.

Here, we resolve this outstanding challenge by using spin–charge hybridization to couple the electric field of a single photon to a single spin in silicon25,31–34. We measure spin–photon coupling rates  of up to 11 MHz, nearly five orders of magnitude higher than typical magnetic-dipole coupling rates. These values of  exceed both the photon decay rate  and the spin decoherence rate , firmly anchoring our spin–photon system in the strong-coupling regime26,29,30.

Our coupling scheme consists of two stages of quantum-state hybridization. First, a single electron is trapped within a gate-defined silicon double quantum dot (DQD) that has a large electric-dipole moment. A single photon confined within a microwave cavity hybridizes with the electron charge state through the electric-dipole interaction35,36. Second, a micrometre-scale magnet (magnetron) placed over the DQD hybridizes electron charge and spin by producing an inhomogeneous magnetic field13–34. The combination of the electric-dipole interaction and spin–charge hybridization gives rise to a large effective spin–photon coupling rate. At the same time, the relatively low level of charge noise in the device ensures that the effective spin decoherence rate remains below the coherent coupling rate , a criterion that has hampered previous efforts to achieve strong spin–photon coupling37.

As well as demonstrating a coherent spin–photon interface, we also show that our device architecture is capable of single-spin control and readout. Single-spin rotations are electrically driven38 and the resulting spin state is detected through a dispersive phase shift in the cavity transmission, which reveals Rabi oscillations36.

Spin–photon interface

The device that enables strong spin–photon coupling is shown in Fig. 1a and contains two gate-defined DQDs fabricated using an overlapping aluminium gate stack (Fig. 1b). The gates are electrically biased to create a double-well potential that confines a single electron in the underlying natural-silicon quantum well (Fig. 1c). A plunger gate (P2) on each DQD is connected to the centre pin of a half-wavelength niobium superconducting cavity with a centre frequency of . We measure spin–photon coupling rates  of up to 11 MHz, nearly five orders of magnitude higher than typical magnetic-dipole coupling rates. These values of  exceed both the photon decay rate  and the spin decoherence rate , firmly anchoring our spin–photon system in the strong-coupling regime26,29,30.

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For simplicity, only one DQD is active at a time for all of the measurements presented here. The cavity is driven by a coherent microwave tone at frequency $f = f_c$ and power $P \approx -133$ dBm (corresponding to approximately 0.6 photons in the cavity, determined on the basis of AC Stark shift measurements of the spin–qubit frequency in the dispersive regime; see Extended Data Fig. 2).  

The normalized cavity transmission amplitude $A/A_0$ is displayed in Fig. 1d as a function of the voltages $V_{p1}$ and $V_{p2}$ on gates P1 and P2 of the first DQD (DQD1), which reveals the location of the (1, 0) $\leftrightarrow$ (0, 1) inter-dot charge transition (see Extended Data Fig. 3 for overall stability diagrams). Here $(n, m)$ denotes a charge state, with the number of electrons in the left (P1) and right (P2) dot being $n$ and $m$, respectively. The charge–photon coupling rate is estimated quantitatively by measuring $A/A_0$ as a function of the DQD level detuning $\epsilon$ (Fig. 1e). By fitting the data with the cavity input–output theory model using $\kappa/(2\pi) = 1.3$ MHz, we find $g_c/(2\pi) = 40$ MHz and $2\pi/\hbar = 4.9$ GHz, where $g_c$ is the inter-dot tunnel coupling and $\hbar$ is the Planck constant. A charge decoherence rate of $\gamma_c/(2\pi) = 35$ MHz is also estimated from the fit and confirmed independently using microwave spectroscopy with $2\pi/\hbar = 5.4$ GHz (refs 19, 20, 42). Fine control of the DQD tunnel coupling, which is critical for achieving spin–charge hybridization, is shown in Fig. 1f, in which $2\pi/\hbar$ is plotted as a function of the voltage $V_{p2}$ on the inter-dot barrier gate B2. A similar characterization of the second DQD (DQD2) yields $g_c/(2\pi) = 37$ MHz and $\gamma_c/(2\pi) = 45$ MHz at the (1, 0) $\leftrightarrow$ (0, 1) inter-dot charge transition. Owing to the higher impedance of the resonator, the values of $g_c$ measured here are much larger than in previous silicon DQD devices, which is helpful for achieving strong spin–photon coupling. In general, there are device–to–device variations in $\gamma_c$ (refs 19, 43). It is unlikely the slightly higher charge decoherence rate is a result of our device geometry, because the Purcell decay rate is estimated to be $\Gamma_P/(2\pi) \approx 0.02$ MHz $\ll \gamma_c/(2\pi)$. Excited valley states are not visible in the cavity response of either DQD, suggesting that they have negligible population. We therefore exclude valleys from the analysis below.

**Strong single spin–photon coupling**

We now demonstrate strong coupling between a single electron spin and a single photon, as evidenced by the observation of vacuum Rabi splitting. Vacuum Rabi splitting occurs when the transition frequency of a two-level atom $f_s$ is brought into resonance with a cavity photon of frequency $f_c$ (refs 21, 23). Light–matter hybridization leads to two vacuum–Rabi-split peaks in the cavity transmission. For our single-spin qubit, the transition frequency between two Zeeman-split spin states is $s = g_s\mu_B B_0$, where $g_s$ is the g-factor of the electron, $\mu_B$ is the Bohr magneton and $B_0$ is the magnetic field. To bring $f_s$ into resonance with $f_c$, we vary the external magnetic field $B_z$ along the z axis while measuring the cavity transmission spectrum $A/A_0$ as a function of the drive frequency $f$ (Fig. 2a). Vacuum Rabi splittings are clearly observed at $B_z = -91.2$ mT and $+92.2$ mT, indicating that $E_0 = h f_c$ at these field values and that the single spin is coherently hybridized with a single cavity photon. These measurements are performed on DQD1, with $2\pi/\hbar = 7.4$ GHz and $\epsilon = 0$. The dependence of $g_c$ on $\epsilon$ and $t_c$ is investigated below.

Assuming $g_c = 2$ for silicon, we estimate that an intrinsic field of about $92 \pm 2$ mT is needed for the single qubit, the transition frequency between two Zeeman-split spin states. Strong single spin–photon coupling is obtained by measuring $A/A_0$ as a function of the DQD level detuning $\epsilon$ (Fig. 1e).

To further verify the strong spin–photon coupling, we plot the cavity transmission spectrum at $B_z = 92.2$ mT (Fig. 2b). The two normal-mode peaks are separated by the vacuum Rabi frequency $2\pi g_c/(2\pi) = 11.0$ MHz, giving an effective spin–photon coupling rate of $g_c/(2\pi) = 5.5$ MHz. The photon decay rate at finite magnetic field is extracted from the line width of $A/A_0$ at $B_z = 90.3$ mT, at which $2\pi/\hbar = 6.7$ GHz is largely dominated by $f_c$, yielding $\kappa/(2\pi) = 1.8$ MHz. A spin decoherence rate of $\gamma_c/(2\pi) = 2.4$ MHz, with contributions from both charge
Strong single spin–photon coupling. a, $A/A_0$ as a function of the cavity drive frequency $f$ and an externally applied magnetic field $B_z$ for DQD1. Insets show data from DQD2 at the same values of $t_c$ and $\varepsilon$ and plotted over the same range of $f$. $B_z$ ranges from $-94$ mT to $-91.1$ mT (9.1 mT to 94 mT) for the left (right) inset. b, $A/A_0$ as a function of $f$ for DQD1 at $B_z = 90.3$ mT (red) and $B_z = 92.2$ mT (blue). c, $A/A_0$ as a function of $f$ for DQD2 at $B_z = 91.1$ mT (red) and $B_z = 92.6$ mT (blue). The spin–photon coupling rate $g_s/(2\pi)$ corresponds to half the frequency separation and so is 5.3 MHz for DQD1 and 5.3 MHz for DQD2.

Electrical control of spin–photon coupling

For quantum information applications it is desirable to turn qubit–cavity coupling rapidly on for quantum-state transfer and rapidly off for qubit-state preparation. Rapid control of the coupling rate is often accomplished by quickly modifying the qubit–cavity detuning $f_s - f_c$. Practically, such tuning can be achieved by varying the qubit transition frequency $f_s$ with voltage or flux pulses or by using a tunable cavity. These approaches are not directly applicable for control of the spin–photon coupling rate because $f_s$ depends primarily on magnetic fields that are difficult to vary on nanosecond timescales. In this section, we show that control of the spin–photon coupling rate can be achieved electrically by tuning $\varepsilon$ and $t_c$ (refs 32, 40).

We first investigate the dependence of $g_s$ on $t_c$. In Fig. 3a we show measurements of $A/A_0$ as a function of $B_z$ and $f$ for $\varepsilon = 0$, $\varepsilon = 20\mu$eV and $\varepsilon = 40\mu$eV. At $\varepsilon > 20\mu$eV (about 4.8 GHz), vacuum Rabi splitting is observed at $B_z = 92.1$ mT with a spin–photon coupling rate of $g_s/(2\pi) = 1.0$ MHz that is substantially lower than the value of $g_s/(2\pi) = 5.5$ MHz obtained at $\varepsilon = 0$. At $\varepsilon = 40\mu$eV (about 9.7 GHz), only a small dispersive shift is observed in the cavity transmission spectrum at $B_z = 91.8$ mT, suggesting a further decrease in $g_s$. These observations are qualitatively understood by considering that at $\varepsilon = 0$ the electron is delocalized across the DQD and forms molecular bonding (|−⟩) or anti-bonding (|+⟩) charge states (Fig. 3c). In this regime, the electron energy levels are well separated, and the electron spin experiences a large oscillating magnetic field, resulting in a substantial spin–photon coupling rate. By contrast, with $|\varepsilon| > t_c$, the electron is localized within one dot and it is natural to work with a basis of localized electronic wavefunctions $|L\rangle$ and $|R\rangle$, where $L$ and $R$ correspond to the electron being in the left and right dot, respectively (Fig. 3c). In this effectively single-dot regime, the displacement of the electron wavefunction by the cavity electric field is estimated to be about 3 pm for a single-dot orbital energy of $E_{\text{orb}} = 2.5$ meV (ref. 48), greatly suppressing the spin–photon coupling mechanisms. The large difference in the effective displacement lengths between the single-dot and double-dot regimes also implies an improvement in the spin–photon coupling rate at $\varepsilon = 0$ of approximately two orders of magnitude compared to $|\varepsilon| > t_c$. Alternatively, the reduction of $g_s$ may be...
interpreted as a result of suppressed hybridization between the $[-, \uparrow]$ and $[+, \downarrow]$ states due to their growing energy difference at larger $|e|$, as evident from Fig. 3c (see discussion below). Here $[-]$ ($\uparrow$) denotes an electron spin that is aligned (anti-aligned) with $B^\text{ext}$. These measurements highlight the important role of charge hybridization in the DQD.

Additional electric control of $g_e$ is enabled by voltage tuning $t_c$ (Fig. 1f). In Fig. 3b we show $g_e/(2\pi)$ and $\gamma_t/(2\pi)$ as functions of $2t_c/h$ at $e = 0$, as extracted from vacuum Rabi splitting measurements and microwave spectroscopy of the electron spin resonance (ESR) transition line width (Figs 2b, 4b, Extended Data Figs 4, 5). Both rates increase rapidly as $2t_c/h$ approaches the Larmor precession frequency $E_Z/h \approx 5.8$ GHz, and a spin–photon coupling rate as high as $g_e/(2\pi) = 11.0$ MHz is found at $2t_c/h = 5.2$ GHz. These trends are consistent with the DQD energy-level spectrum shown in Fig. 3c.33,34,41. With $2t_c/h \gg E_Z/h$ and $e = 0$, the two lowest energy levels are $[-, \downarrow]$ and $[+, \uparrow]$ and the electric-dipole coupling to the cavity field is small. As $2t_c$ is reduced and made comparable to $E_Z$, the ground state remains $[-, \downarrow]$, but the excited state becomes an admixture of $[-, \downarrow]$ and $[+, \downarrow]$ owing to the magnetic-field gradient $B^\text{M}_{\perp} - B^\text{M}_{\parallel} = 2B^\text{M}$ and the small energy difference between the states. The quantum transition that is close to resonance with $E_Z$ is now partially composed of a change in charge state from $- \rightarrow 0$ to $+ \rightarrow 0$, which responds strongly to the cavity electric field and gives rise to larger values of $g_e$. For $2t_c/h < E_Z/h$, a decrease in $t_c$ increases the energy difference between $[-, \downarrow]$, which reduces their hybridization and results in a smaller $g_e$. We note that hybridization with charge states increases the susceptibility of the spin to charge noise and relaxation, resulting in an effective spin decoherence rate $\gamma_t$ that is also strongly dependent on $t_c$ (Fig. 3b).33,34,41. Theoretical predictions of $g_e$ and $\gamma_t$ as functions of $2t_c/h$ based on measured values of $g_e$ and $\gamma_t$ (Fig. 1e) are in good agreement with the data (Fig. 3b).41. The discrepancy in the fit of $\gamma_t$ is discussed in Methods. The electric control of spin–photon coupling demonstrated here allows the spin qubit to switch quickly between regimes with strong coupling to the cavity and idle regimes in which the spin–photon coupling rate and susceptibility to charge decoherence are small.

**Dispersive readout of a single spin**

The preceding measurements demonstrate the ability to couple a single electron spin to a single photon coherently, potentially enabling long-range spin–spin couplings.46,47. For the device to serve as a building block for a quantum processor, it is also necessary to prepare, control and read out the spin state of the trapped electron deterministically. We first induce spin transitions by driving gate P1 with a continuous microwave tone of frequency $f_1$ and power $P_1 = -106$ dBm. When $f_1 \approx E_Z/h$, the excited-state population of the spin qubit $P_\uparrow$ increases and the ground-state population $P_\downarrow$ decreases. In the dispersive regime, in which the qubit–cavity detuning $\Delta_\text{cav} = E_Z/h - f_1$ satisfies $|\Delta_\text{cav}| \gg g_e/(2\pi)$, the cavity transmission experiences a phase response $\Delta \phi \approx -\tan^{-1}(|g_e/(\pi:\Delta_\text{cav})|)$ for a fully saturated ($P_\uparrow = 0.5$) qubit.43,44. It is therefore possible to measure the spin state of a single electron by probing the cavity transmission. As a demonstration, we spectroscopically probe the ESR transition by measuring $\Delta \phi$ as a function of $f_1$ and $B^\text{ext}$ (Fig. 4a). These data are acquired with $2t_c/h = 9.5$ GHz and $e = 0$. The ESR transition is clearly visible as a narrow feature with $\Delta \phi = 0$ that shifts to higher $f_1$ with increasing $B^\text{ext}$. $\Delta \phi$ also changes sign as $B^\text{ext}$ increases, consistent with the sign change of the qubit–cavity detuning $\Delta$ when the Larmor precession frequency $E_Z/h$ exceeds $f_1$. The nonlinear response in the small region around $B^\text{ext} = 92$ mT is due to the breakdown of the dispersive condition $|\Delta_\text{cav}| \gg g_e/(2\pi)$.

Finally, we demonstrate coherent single-spin control and dispersive spin–state readout. For these measurements, we choose $e = -1$ and $2t_c/h = 11.1$ GHz to minimize the spin decoherence rate $\gamma_t$ (Fig. 3b). Here the spin–photon coupling rate $g_e/(2\pi) = 1.4$ MHz (Fig. 3b). The external field is fixed at $B^\text{ext} = 92.18$ mT, which ensures that the system is in the dispersive regime with $\Delta_\text{cav} = 14$ MHz. A measurement of $\Delta \phi(f_1)$ in the low-power limit (Fig. 4b) yields a Lorentzian line shape with a full-width at half-maximum of 0.81 MHz, which corresponds to a low spin decoherence rate of $g_e/(2\pi) = 0.41$ MHz (refs 19, 42). Qubit control and measurement are achieved using the pulse sequence illustrated in Fig. 4c. Starting with a spin–down state $[-]$ at $e = 0$, the DQD is pulsed to a large detuning $\Delta_\text{cav} = 70$ MHz (about 17 GHz), which decouples the spin from the cavity. A microwave burst with frequency $f_b = 5.874$ GHz, power $P_b = -76$ dBm and duration $\tau_B$ is subsequently applied to P1 to drive a spin rotation.9,36,38. The DQD is then pulsed adiabatically back to $e = 0$ for a fixed measurement time $T_M$ for dispersive readout. To reinitialize the qubit, we choose $T_M = 20 \mu s \gg T_\text{rel}(e = 0)$, where $T_\text{rel}(e = 0) = 3.2 \mu s$ is the spin relaxation time measured at $e = 0$ (Extended Data Fig. 6). Figure 4d displays the time-averaged $\Delta \phi$ as a function of $T_M$, obtained with an integration time of 100 ms for each data point. We observe coherent single-spin Rabi oscillations with a Rabi frequency of $f_b = 6$ MHz. In contrast to readout approaches that rely on spin-dependent tunneling,9,38,49, our dispersive cavity-based readout corresponds in principle to quantum nondemolition readout.24. The readout scheme is also distinct from previous work that used a cavity-coupled InAs DQD, which detects the spin state through Pauli blockade rather than spin–photon coupling.26. In addition to enabling single spin–photon coupling, our device is capable of preparing, controlling and dispersively reading out single spins.

**Conclusion**

We have realized a coherent spin–photon interface at which a single spin in a silicon DQD is strongly coupled to a microwave-frequency
photon through the combined effects of the electric–dipole interaction and spin–charge hybridization (see Methods for a discussion of the prospects of applying the spin–photon interface to realize cavity-mediated spin–spin coupling). Spin–photon coupling rates of up to 11 MHz are measured in the device, exceeding magnetic-dipole coupling rates by nearly five orders of magnitude. The spin decoherence rate is strongly dependent on the inter-dot tunnel coupling $t$, and ranges from 0.4 MHz to 6 MHz, possibly limited by a combination of charge noise, charge relaxation and remnant nuclear field fluctuations. All-electric control of spin–photon coupling and coherent manipulation of the spin state are demonstrated, along with dispersive readout of the single spin, which lays the foundation for quantum non-demolition readout of spin qubits. These results could enable the construction of an ultra-coherent spin quantum computer with photonic interconnects and readout channels, with the capacity for surface codes, ‘all-to-all' connectivity and easy integration with other solid-state quantum systems such as superconducting qubits.24,46,47,50–52.

We note that two related preprints appeared after the submission of this Article: ref. 53 presents some of the results discussed here, and ref. 54 explores a different approach to spin–photon coupling and demonstrates coherent coupling of a single quantum dot to microwave–frequency photons.

Online Content Methods, along with any additional Extended Data display items and Source Data, are available in the online version of the paper; references unique to these sections appear only in the online paper.

Acknowledgements We thank A. J. Sigillito for technical assistance and M. J. Gullans for discussions. This work was supported by the US Department of Defense under contract H98230-15-CD0453, Army Research Office grant W911NF-15-1-0149, and the Gordon and Betty Moore Foundations EPiQS Initiative through grant GBMF4535. Devices were fabricated in the Princeton University Quantum Device Nanofabrication Laboratory.


Reviewer Information Nature thanks T. Meunier and the other anonymous reviewer(s) for their contribution to the peer review of this work.
METHODS

Device fabrication and measurement. The Si/Ge heterostructure consists of a 4-nm-thick Si cap, a 50-nm-thick Si1-xGex spacer layer, a 8-nm-thick natural-Si quantum well and a 225-nm-thick Si1-xGex layer on top of a linearly graded Si1-xGex, relaxed buffer substrate. Design and fabrication details for the superconducting cavity and DQDs are described elsewhere. The approximately 200-nm-thick Co micromagnet is defined using electron beam lithography and lift off. In contrast to earlier devices, the gate filter for P1 was changed to an L1–C–L2 filter, with L1 = 4 nH, C = 1–2 pF and L2 = 12 nH (ref. 43). This three-filter set allows microwave signals below 2.5 GHz to pass with less than 3 dB of attenuation. All data are acquired in a dilution refrigerator with a base temperature of 10 mK and electron temperature of \( T_e = 60 \) mK. The measurements of the transmission amplitude and phase response of the cavity (Figs 1, 4) are performed using a homodyne detection scheme. The measurements of the transmission spectra of the cavity (Figs 2, 3) are performed using a network analyser. The microwave drive applied to P1 (Fig. 4) is provided by a vector microwave source and the detuning pulses are generated by an arbitrary waveform generator, which also controls the timing of the microwave burst (Fig. 4d).

To maximize the magnetization of the Co micromagnet and minimize hysteresis, data at positive (negative) external applied magnetic fields (Fig. 2a) are collected after \( B_{\text{ext}}^{\text{off}} \) is first ramped to a large value of \(+300 \text{ mT} \) (\(-300 \text{ mT}\)). A small degree of hysteresis still remains for the micromagnet of DQD1, as can be seen by the different magnitudes of \( B_{\text{ext}}^{\text{off}} \) at which positive- and negative-field vacuum Rabi oscillations are observed (Fig. 2a). In Fig. 4a, the slope of the ESR transition is \( \Delta E_{\text{FSR}} / \Delta B = 44 \text{ MHz mT}^{-1} \), which is higher than the value (28 MHz mT\(^{-1}\)) expected for a fully saturated micromagnet. The slope of the transition suggests that the micromagnet is not fully magnetized and has a magnetic susceptibility of \( \chi_{\text{m}} = 3.2 \times 10^{-5} \text{ mT}^{-1} \). Theoretical models for spin–photon coupling and spin decoherence. Here we present a brief derivation of the analytical expressions for the spin–photon coupling rate \( g_s \) and the spin decoherence rate \( \gamma_s \). A more extensive discussion of spin–photon coupling and spin decoherence specific to our device architecture is presented in ref. 41. 

We start from the Hamiltonian that describes the DQD field in the \( x \) direction satisfies \( \langle B_{\text{ext}}^{\text{off}} + B_{\text{ext}}^{\text{on}} \rangle / 2 = 0 \), which is a good approximation given the geometry of the micromagnet and its alignment with the DQD. We add the electric-dipole coupling to the cavity with the Hamiltonian

\[
H_1 = g_s (a + a^\dagger) \tau_z \tag{2}
\]

where \( a \) and \( a^\dagger \) are the photon operators for the cavity. The electric-dipole operator can be expressed in the eigenbasis \( | n \rangle \) of \( H_1 \) as

\[
\tau_z = \sum_{n,n=0} d_{nm}|n\rangle\langle m| \tag{1}
\]

We then write the quantum Langevin equations for the operators \( a \) and \( a^\dagger \):
Combined with charge–photon coupling, the overall equations of motion (equation (2)) in a rotating frame with a drive frequency \( f \approx J \), assume the form

\[
\dot{\mathbf{a}} = -i \Delta a - \frac{\kappa}{2} \mathbf{a} + \sqrt{\kappa} \mathbf{a}_{\text{in}} - i \gamma g (\mathbf{a}_{\text{in}} + \mathbf{a}_{\text{out}})
\]

\[
\mathbf{a}_{\text{in}} = -ib \mathbf{a}_{\text{r}} - \gamma \sin \left( \frac{\Phi}{2} \right) \mathbf{a}_{\text{in}} + \gamma \sin (\Phi) \mathbf{a}_{\text{in}} - i \gamma g \mathbf{a}_{\text{in}}
\]

\[
\mathbf{a}_{\text{out}} = -ib \mathbf{a}_{\text{r}} - \gamma \cos \left( \frac{\Phi}{2} \right) \mathbf{a}_{\text{in}} + \gamma \sin (\Phi) \mathbf{a}_{\text{in}} - i \gamma g \mathbf{a}_{\text{in}}
\]

The \( \delta_1 \) and \( \delta_2 \) terms are defined as \( \delta_1 (\Phi) = (E_2 - E_0)/h - f \) and \( \delta_2 (\Phi) = (E_2 - E_0)/h - f \), where \( E_{0,1,2} \) correspond to the energy of the \( |0\rangle \), \( |1\rangle \) and \( |2\rangle \) state, respectively. Steady-state solutions to the above equations give the electric susceptibility of the spin qubit transition \( \chi_{0,1} = \sigma_0 / a = g \sqrt{(\delta_1 - \delta_2)} \), where we have identified a charge-induced spin decoherence rate \( \gamma_0 \) and a coherent spin–spin coupling to be implemented in our device architecture as well.

Prospects for long-range spin–spin coupling. The coherent spin–photon interface may be readily applied to enable spin–spin coupling across the cavity bus. Here we evaluate two possible schemes for implementing such a coupling, both of which have been demonstrated with superconducting qubits\(^{46-55}\). The first approach uses direct photon exchange to perform quantum-state transfer between two qubits\(^{56}\). The second approach to spin–spin coupling uses virtual photon exchange\(^{46}\). In this scheme, both spin qubits would operate in the dispersive regime, with an effective coupling rate of \( J \approx g^2 / \left( \Delta_1 + 1 / \Delta_2 \right) / 2 \), where \( \Delta_1 \) and \( \Delta_2 \) are the qubit–cavity detunings for qubits 1 and 2, respectively. Assuming that both qubits operate with an equal detuning \( \Delta_1 = \Delta_2 = 10 \text{GHz} \) to minimize Purcell decay, \( J = g^2 / 10 \). In Extended Data Fig. 8a, we plot the ratio \( g^2 / (\Delta_1 + 1 / \Delta_2) \) as a function of \( 2t_\text{ext}/h \approx 10 \text{GHz} \). The dominant spin–spin mechanism is probably hyperfine-induced dephasing by the \( 29\text{Si} \) nuclei in this regime (the coherence rate \( \gamma_0 / (2\Delta) + 0.4 \text{MHz} \) is close to the coherence rates commonly found with single-spin qubits in natural \( \text{Si} \); ref. 38), transitioning to isotopically purified \( 29\text{Si} \) host materials is likely to lead to an order-of-magnitude reduction in \( \gamma_0 / (2\Delta) \), as demonstrated recently\(^{39}\). Such an improvement will allow virtual-photon-mediated spin–spin coupling to be implemented in our device architecture as well.

Data availability. The data that support the findings of this study are available from the corresponding author on reasonable request. Source Data for Figs 1–4 and Extended Data Figs 2–8 are available with the online version of the paper.
**Extended Data Figure 1 | Micromagnet design.** To-scale drawing of the micromagnet design, superimposed on top of the SEM image of the DQD. The coordinate axes and the direction of the externally applied magnetic field $B^\text{ext}_z$ are indicated at the bottom. In this geometry, the DQD electron experiences a homogeneous z field $B^z_\text{ext} \approx B^M_\text{ext}$. The total x field $B^x$ that is experienced by the electron is spatially dependent, being approximately $B^x_\text{L} \approx B^M_\text{L}$ when the electron is in the L (R) dot ($|\epsilon| \gg t_c$) and $(B^M_\text{L} + B^M_\text{R})/2$ when the electron is delocalized between the DQDs ($\epsilon = 0$). The y field $B^y$ for the DQD electron is expected to be small compared to the other field components for this magnet design.
Extended Data Figure 2 | Photon number calibration. The ESR resonance frequency $f_{\text{ESR}}$, measured using the phase response of the cavity $\Delta \phi$ in the dispersive regime (Fig. 4b), is plotted as a function of the estimated power at the input port of the cavity $P$ (data). The device is configured with $g_s/(2\pi) = 2.4$ MHz and spin–photon detuning $\Delta/(2\pi) \approx -18$ MHz. The dashed line shows a fit to $f_{\text{ESR}} = f_{\text{ESR}}(P = 0) + (2n_{\text{ph}}g_s^2/\Delta)/(2\pi)$, where $n_{\text{ph}}$ is the average number of photons in the cavity, plotted as the top x axis. The experiments are conducted with $P \approx -133$ dBm (0.05 fW), which corresponds to $n_{\text{ph}} \approx 0.6$. The error bars indicate the uncertainties in the centre frequency of the ESR transition.
**Extended Data Figure 3 | DQD stability diagrams.** The cavity transmission amplitude $A/A_0$ (a, c) and phase response $\Delta \phi$ (b, d) are plotted as functions of $V_{P1}$ and $V_{P2}$ for DQD1 (a, b) and DQD2 (c, d), obtained with $f = f_c$. The $(1, 0) \rightarrow (0, 1)$ transitions are clearly identified on the basis of these measurements and subsequently tuned close to resonance with the cavity for the experiments described in the main text. The red circles indicate the locations of the $(1, 0) \rightarrow (0, 1)$ transitions of the two DQDs.
Extended Data Figure 4 | Spin decoherence rates at different DQD tunnel couplings. ESR line, as measured in the cavity phase response $\Delta \phi(f_s)$, is shown for different values of $2t_c/h$ in the low-power limit (data). $\xi = 0$ for every dataset. Dashed lines are fits with Lorentzian functions and $\gamma_s/(2\pi)$ is determined as the half-width at half-maximum of each Lorentzian. The spin–photon detuning $|\Delta| \approx 10 g_s$ for each dataset, to ensure that the system is in the dispersive regime.
Extended Data Figure 5 | Spin–photon coupling strengths at different DQD tunnel couplings. a, b, Vacuum Rabi splittings for $2t_{c}/h < f_{c}$ (a) and $2t_{c}/h > f_{c}$ (b), obtained by varying $B_{ext}^{p}$ until a pair of resonance peaks with approximately equal heights emerges in the cavity transmission spectrum $A/A_{0}$, $g_{s}/(2\pi)$ is then estimated as half the frequency difference between the two peaks. $\varepsilon = 0$ for every dataset. $g_{s}$ is difficult to measure for $5.2 \text{ GHz} < 2t_{c}/h < 6.7 \text{ GHz}$ owing to the small values of $A/A_{0}$ that arise from the large spin decoherence rates $\gamma_{s}$ in this regime.
Extended Data Figure 6 | Spin relaxation at $\varepsilon = 0$. The time-averaged phase response of the cavity $\Delta \phi$ is shown as a function of wait time $T_M$ (data), measured using the pulse sequence illustrated in Fig. 4c. The microwave burst time is fixed at $\tau_B = 80$ ns. The dashed line shows a fit using the function $\phi_0 + \phi_1 (T_1/T_M)[1 - \exp(-T_M/T_1)]$, which yields a spin relaxation time of $T_1 \approx 3.2 \mu$s. The experimental conditions are the same as for Fig. 4d.
Extended Data Figure 7 | Theoretical fits to vacuum Rabi splittings.

The calculated cavity transmission spectra (black solid lines) are superimposed on the experimentally measured vacuum Rabi splittings shown in Fig. 2b, c (data). The calculations are produced with $g_c/(2\pi) = 40 \text{ MHz}$, $\kappa/(2\pi) = 1.8 \text{ MHz}$, $\gamma_c/(2\pi) = 105 \text{ MHz}$, $B_z = B_{z,\text{ext}} = 209 \text{ mT}$, $B_{z,\text{ext}} = 92.2 \text{ mT}$ for DQD1 (DQD2). For DQD2, $B_{z,\text{ext}} = 92.6 \text{ mT}$. $B_{z,\text{ext}} = (B_{z,\text{ext}}^\text{M} - B_{z,\text{ext}}^\text{L})/2 = 15 \text{ mT}$ and $2\kappa/h = 7.4 \text{ GHz}$ for DQD1 (DQD2). For comparison, $A(f)/A_0$, simulated for a two-level charge qubit with a decoherence rate of $\gamma_c/(2\pi) = 2.4 \text{ MHz}$ coupled to a cavity with $\kappa/(2\pi) = 1.8 \text{ MHz}$ at a rate $g_c/(2\pi) = 5.5 \text{ MHz}$, is shown in a for thermal photon numbers of $n_{\text{th}} = 0.02$ (black dashed line) and $n_{\text{th}} = 0.5$ (red dashed line).
Extended Data Figure 8 | Prospect for long-range spin–spin coupling. a, The ratio $2g_s/(\kappa/2 + \gamma_s)$ as a function of $2t_c/h$, calculated using the data in Fig. 3b and $\kappa/(2\pi) = 1.8$ MHz. b, The ratio $g_s/\gamma_s$ as a function of $2t_c/h$, also calculated using the data in Fig. 3b.