Circuit quantum electrodynamics with a spin qubit

K. D. Petersson, L. W. McFaul, M. D. Schroer, M. Jung, J. M. Taylor, A. A. Houck & J. R. Petta

Electron spins trapped in quantum dots have been proposed as basic building blocks of a future quantum processor. Although fast, 180-picosecond, two-quantum-bit (two-qubit) operations can be realized using nearest-neighbour exchange coupling, a scalable, spin-based quantum computing architecture will almost certainly require long-range qubit interactions. Circuit quantum electrodynamics (cQED) allows spatially separated superconducting qubits to interact via a superconducting microwave cavity that acts as a 'quantum bus', making possible two-qubit entanglement and the implementation of simple quantum algorithms. Here we combine the cQED architecture with spin qubits by coupling an indium arsenide nanowire double quantum dot to a superconducting cavity. The architecture allows us to achieve a charge–cavity coupling rate of about 30 megahertz, consistent with coupling rates obtained in gallium arsenide quantum dots. Furthermore, the strong spin–orbit interaction of indium arsenide allows us to drive spin rotations electrically with a local gate electrode, and the charge–cavity interaction provides a measurement of the resulting spin dynamics. Our results demonstrate how the cQED architecture can be used as a sensitive probe of single-spin physics and that a spin–cavity coupling rate of about one megahertz is feasible, presenting the possibility of long-range spin coupling via superconducting microwave cavities.

The weak magnetic moment of the electron makes it difficult to couple spin qubits that are separated by a large distance. Approaches to transferring spin information include physically shuffling electrons with surface acoustic waves or using exchange-coupled spin chains, both of which are experimentally challenging to realize. An attractive alternative for realizing long-distance spin–qubit interactions is to interface spins with a superconducting microwave cavity in the cQED architecture. Unfortunately, direct coupling between a single spin magnetic dipole and the magnetic field of the cavity results in a spin–cavity coupling rate of $g_{\text{spin}}/2\pi \sim 10 \, \text{Hz}$, which is much too weak to be useful for quantum information processing. Recent experiments have explored coupling ensembles of spins to superconducting resonators, with the large number of spins, $N_{\text{spin}} \approx 10^{12}$, giving a $\sim N_{\text{spin}}^{1/2}$ enhancement in the spin–cavity coupling rate.

Another approach to spin–cavity coupling relies on the spin–orbit interaction. Spin–orbit coupling mixes spin and orbital degrees of freedom, resulting in spin states that have some orbital character, the spin–orbit doublets, $|\parallel\rangle$ and $|\perp\rangle$. Although electron spin states cannot be coupled directly to an electric field, the spin–orbit interaction enables electrical control by perturbing the orbital component of the electron wavefunction. Fast, coherent electrical control of spin states in quantum dots has been demonstrated in InAs nanowires where the spin–orbit interaction strength is large. The cQED architecture could be used to couple two distant InAs nanowire quantum dot spin qubits with the spin–orbit interaction enabling a significantly increased spin–cavity coupling rate, $g_{\text{cavity}}$. In this Letter, we present the first steps towards realizing this approach and couple the electric field of a high-quality-factor superconducting cavity to an InAs nanowire double quantum dot (DQD) device. We determine the charge–cavity coupling rate, $g_{\text{cavity}}$, for the molecular orbital states of a single excess charge in the DQD. Then, with each of the two quantum dots acting as a spin qubit, we perform fast electrical spin–state control followed by single spin read-out using the microwave cavity. Our results demonstrate that spin qubits, which require substantial magnetic fields for their operation, can be readily integrated into the superconducting cQED architecture and pave the way for long-range coupling of spin qubits via microwave cavities.

Our hybrid spin-qubit/superconducting device is shown in Fig. 1. We fabricate a half-wavelength superconducting Nb resonator (the cavity) with a resonance frequency of $f_0 = \omega_0/2\pi = 6.2 \, \text{GHz}$ and quality factor of $Q_0 \approx 2000$ (Supplementary Information, section 2). The amplitude and phase responses of the cavity are detected using a homodyne measurement with a microwave probe frequency $f_p$ (ref. 5). We couple a single InAs nanowire spin qubit to the electric field generated by the cavity. The qubit consists of a DQD defined in an InAs nanowire. A series of Ti/Au depletion gates create a simple double-well confinement potential containing $N_{\text{L}}$, $N_{\text{R}}$ electrons, where $N_{\text{L}}$ and $N_{\text{R}}$ are the numbers of electrons in

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**Figure 1** Hybrid DQD/superconducting resonator device. (a) Circuit schematic and micrograph of the hybrid device design. Transmission through the half-wavelength superconducting Nb resonator is measured using homodyne detection at a frequency $f_0$. Standard d.c. transport measurements are made possible by applying a source–drain bias, $V_{\text{SD}}$, to the DQD using a $\sim 4$–$8$–$\mu$H spiral inductor that is connected to the voltage node of the resonator. See Supplementary Information, section 1, for further details. (b, c) Scanning electron micrograph (b) and cross-sectional schematic view (c) of a typical nanowire DQD. The left and right barrier gates ($B_\text{L}$ and $B_\text{R}$), left and right plunger gates ($L$ and $R$), and middle gate ($M$) are biased to create a double-well potential within the nanowire. The drain contact of the nanowire, $D$, is grounded, and the source contact, $S$, is connected to an antinode of the resonator, oscillating at a voltage $V_{\text{Cavity}}(t)$. An a.c. voltage at a frequency $f_\text{ac}$ is applied to gate $M$ to generate an oscillating electric field, $E_0$.
the left- and right-hand dots, respectively. We tune the tunnel coupling, $t_{\text{DQD}}$, of the DQD by adjusting the voltage, $V_{\text{M}}$, on the middle barrier gate (M in Fig. 1b). A trapped electron in the DQD has an electric dipole moment of $d \approx 1,000e_{\text{Bohr}}$, where $e_{\text{Bohr}}$ is the Bohr radius and $e$ is the electronic charge.

For a spin in a single quantum dot, the calculation in ref. 18 predicts a spin–cavity coupling rate of $g_{\text{SC}} \approx g_{\text{D}}(E_{\text{R}}/\sqrt{\Delta E_0})(1/\lambda_{\text{SO}})$, where $E_{\text{R}}$ is the Zeeman splitting of the spin states, $\Delta E_0$ is the orbital level spacing, $l$ is the quantum dot size and $\lambda_{\text{SO}}$ is the spin–orbit length, which characterizes the strength of the spin–orbit interaction. Therefore, strong spin–cavity coupling requires two key components: a large charge–cavity coupling rate and a strong spin–orbit interaction. Charge–cavity coupling is achieved through the electric dipole interaction, as in experiments with superconducting qubits. An oscillating electric field, with amplitude $E_{0}$, periodically displaces the electron quantum dot potential by a distance $r_{\text{DQD}} = eE_{0}/\Delta E_0$ (Fig. 1c), which is dependent on the quantum dot confinement as determined by $\Delta E_0$ and $l$ (ref. 19).

To enhance the cavity electric field at the position of the DQD—and maximize the charge–cavity coupling rate—the source and drain contacts of the nanowire are connected directly to the voltage antinode and the ground of the resonator. In the presence of a strong spin–orbit interaction, the displacement of the electron can induce spin-state rotations at a rate $E_{0}/\hbar \times r_{\text{DQD}}$, where $\hbar$ denotes Planck’s constant divided by $2\pi$, with the linear dependence in $E_{0}$ due to the Van Vleck zero-field cancellation of the spin–orbit term. Strong spin–orbit coupling is achieved using InAs, which has a short spin-orbit length, $\lambda_{\text{SO}} \approx 100$ nm (ref. 20).

We first characterize the interaction between an electron trapped in a DQD and the electric field of the cavity, demonstrating a 30-MHz charge–cavity coupling rate with this device architecture. We focus on the cavity response near the $(M,N+1) \leftrightarrow (M, N+1)$ interdot charge transition in the many-electron regime ($M = 20$, $N = 20$; Supplementary Information, section 3). The DQD forms a two-level ‘artificial transition in the many-electron regime (refs 10, 22).

A qualitative understanding of the coupling between the quantum dot and the cavity can be obtained by considering the relevant energy scales in the system. The single-dot charging energy, $E_{\text{C}} \approx 12$ meV, is much larger than the relevant photon energy, $h\omega_{\text{L}} \approx 25$ μeV, and the cavity is largely unaffected by the DQD in Coulomb blockade. However, near interdot charge transitions (for example $(M,N+1) \leftrightarrow (M, N+1)$ or transitions with the source and drain electrodes (for example $(M,N) \leftrightarrow (M, N+1)$), the energy scales associated with the DQD are close to the cavity energy, and the cavity is damped, resulting in a phase shift in microwave transmission at the bare cavity frequency. In Fig. 2b, the DQD charge stability diagram is measured around the $(M,N+1) \leftrightarrow (M, N+1)$ transition by probing the phase response of the microwave cavity as a function of the gate voltages $V_{\text{L}}$ and $V_{\text{R}}$ (refs 10, 22).

Quantitative analysis of the cavity response requires a fully quantum mechanical model that accounts for photon exchange between the microwave field and the DQD. In cavity QED, the pertinent interactions are those between an atom with transition frequency $\omega_{\text{L}} = \Omega/\hbar$ and the photon field of the cavity, characterized by the resonance frequency $\omega_{\text{C}}$. The atom and cavity energy levels hybridize when the atom–cavity detuning, $\Delta = \omega_{\text{L}} - \omega_{\text{C}}$, is less than $\gamma_{\text{C}}$, leading to the Jaynes–Cummings ladder of quantum states. When the atom and cavity are detuned in the dispersive limit ($\Delta > \gamma_{\text{C}}$), the cavity field exhibits a phase shift in microwave transmission at the bare cavity frequency that is given by $\phi = -\arctan(2g_{\text{SC}}/\kappa\Delta)$, where $\kappa$ is the cavity decay rate. In Fig. 2d, we plot the phase response of the cavity for several values of the interdot tunnel coupling (see Supplementary Information, section 4 for the magnitude response). We observe a sign change in the phase as the atom–cavity detuning, $\Delta$, is varied from positive to negative values. We fit the phase and magnitude data to a master equation model (Supplementary Information, section 6.1) using a best-fit value of $g_{\text{SC}}/\Delta = 30$ MHz, an inhomogeneous broadening parameter, $\sigma_{\text{ph}}/\hbar = 5.1$ GHz, to account for low-frequency charge noise; and a $V_{\text{M}}$-dependent tunnel coupling that ranges from $1.8 \text{GHz}$ to $7.0 \text{GHz}$ (Fig. 2c). The charge–cavity coupling rate extracted here compares favourably to values obtained using Cooper pair box qubits ($g_{\text{SC}}/\Delta = 6$ MHz), transmon qubits ($g_{\text{SC}}/\Delta = 100$ MHz) and many-electron GaAs quantum dots ($g_{\text{SC}}/\Delta = 50$ MHz).

We characterize the strength of the spin–orbit interaction by operating the device as a spin qubit (Fig. 3). For simplicity, we label the charge states $(1,1)$ and $(0,2)$ (ref. 4). The ground state with two electrons in the right quantum dot is the singlet $S(0, 2)$. At negative detuning, the electrons are separated in a $(1,1)$ charge state, and the four relevant spin states are $|\uparrow\uparrow\rangle$, $|\downarrow\downarrow\rangle$, $|\uparrow\downarrow\rangle$, and $|\downarrow\uparrow\rangle$ (ref. 8). The level diagram is similar to a GaAs singlet–triplet spin qubit, with a key difference being that the $g$-factors for the two spins can vary significantly (Supplementary Information, section 6.2). Interdot tunnel coupling hybridizes the states with singlet character near $\varepsilon = 0$, and an external field results in Zeeman splitting, $E_{Z} = \tilde{g}\mu_{\text{B}}B$, of the spin states, where $\tilde{g}$ is the electronic $g$-factor, $\mu_{\text{B}}$ is the Bohr magneton and $B$ is the magnetic field.
Figure 3 | Spin-qubit spectroscopy. a. The measurement cycle for spin-qubit spectroscopy. The DQD is put in either the |↑↑⟩ or the |↓↓⟩ spin configuration as a result of Pauli spin blockade. EDSR transitions lift this blockade, resulting in current flow through the device. b. Spin-qubit spectroscopy as a function of a.c. driving frequency on the gate, \( f_G \), measured under Pauli blockade. Pauli blockade is lifted by EDSR driving when the microwave frequency is \( f_{\mu} = \frac{\mu B}{h} \), where \( \mu \) (\( i = 1, 2 \)) is the electronic g-factor of dot \( i \). Inset, current through the DQD as a function of gate voltages \( V_L \) and \( V_R \), with \( V_{SD} = 2.5 \) mV. Current is suppressed inside the dashed region owing to Pauli spin blockade. c. Energy levels of the spin–orbit qubit plotted as a function of \( \epsilon \). The data in b are acquired with \( \epsilon = \epsilon' = 1 \) meV.

Spin selection rules result in Pauli blockade at the two-electron transition, a key ingredient for spin preparation and measurement\(^{1,8,26}\) (Fig. 3b, inset). For example, state |↑↑⟩ cannot tunnel to S(0, 2) due to Pauli exclusion. Modulation of the confinement potential with a gate voltage results in spin–orbit-driven electric dipole spin resonance (EDSR) transitions that lift the Pauli blockade\(^{8,19}\). In Fig. 3b, we plot the current, \( I \), through the DQD with \( V_{SD} = 2.5 \) mV and the gates tuned in Pauli blockade (Fig. 3b, white dot in inset). Hyperfine fields rapidly mix spin states when \( E_2 = \frac{\mu B}{B_N} < B_\text{N} \), where \( B_\text{N} \approx 2 \) mT is the hyperfine field\(^8\). At finite fields, the leakage current is non-zero when the a.c. driving frequency on the gate, \( f_\text{G} \), satisfies the electron spin resonance condition \( E_2 = h f_\text{G} \). We observe two resonance conditions corresponding to single spin rotations in the left- and right-hand quantum dots, with g-factors of 8.2 and 10.6 (ref. 8).

In cQED with superconducting qubits, measurements of the spin response can be used for qubit read-out. For spin qubits, around \( \epsilon = 0 \), the DQD has a spin-state-dependent dipole moment due to Pauli blockade that allows spin-state read-out via the superconducting cavity\(^{27}\). We combine quantum control of the spins using EDSR and cavity detection of single-spin dynamics using the pulse sequence shown in Fig. 4a, b. Starting with the spin qubit in state |↑↑⟩, we pulse to negative detuning \( \epsilon = \epsilon' = -2 \) meV and apply a microwave burst of length \( \tau_B \) to drive EDSR transitions. For example, an EDSR \( \pi \)-pulse will drive a spin transition from |↑↑⟩ to |↓↓⟩. The resulting spin state is probed by pulsing back to \( \epsilon = 0 \) for a measurement time \( T_M \). The cavity is most sensitive to charge dynamics near \( \epsilon = 0 \) owing to the different a.c. susceptibilities of spin states |↑⟩⟩ and |↓⟩⟩ (Supplementary Information, section 6.3). In Fig. 4c, we plot the cavity phase shift as a function of \( f_\text{G} \) and \( B \). We again observe two features that follow the standard spin resonance condition, consistent with the d.c. transport data in Fig. 4b. By varying \( T_M \), we fit the measured phase response to theory and estimate a spin lifetime of \( T_\text{I} = 1 \) \( \mu \)s (Fig. 4d). We anticipate that the relaxation time is detuning dependent, with longer spin relaxation times away from \( \epsilon = 0 \) (ref. 28; Supplementary Information, section 5).

We demonstrate coherent control of the spin qubit and read-out via the cavity by varying the EDSR microwave burst length, \( \tau_B \). Figure 4e

Figure 4 | Coherent spin-state control and detection using the microwave cavity. a. Top: pulse sequence used for spin-state control and resonator readout, superimposed on the level diagram. Bottom: the a.c. susceptibility, \( \chi_\text{c} \), is dependent on the spin state of the DQD and allows for sensitive spin readout via the microwave cavity. b. Pulse sequence used for spin-state control and resonator read-out. Starting in state |↑↑⟩, an EDSR burst is applied at far detuning (\( \epsilon = \epsilon' = -2 \) meV). The resultant spin state is then measured at \( \epsilon = 0 \) by probing the cavity transmission using a weak continuous tone of frequency \( f_\text{R} \). c. Phase response of the cavity measured as a function of EDSR drive frequency, \( f_\text{G} \), and external field, \( B \), with \( \tau_B = 100 \) ns and \( T_M = 850 \) ns. EDSR transitions are observed in the phase response, in agreement with the d.c. transport data, with small differences in \( g_\text{G} \) and \( g_\text{R} \), attributable to the difference in sample tuning necessary to optimize the response. d. Measured phase shift as a function of \( T_M \), with \( \tau_B = 100 \) ns, \( B = 90 \) mT and \( f_\text{G} = 13.1 \) GHz. A fit to theory yields a spin relaxation time of \( T_\text{I} = 1 \) \( \mu \)s. e. Phase response of the cavity as a function of EDSR burst length, \( \tau_B \), and approximate driving power at the sample, \( P_{\text{G}} \), for \( B = 86 \) mT, \( f_\text{G} = 9.5 \) GHz and \( T_M = 1.75 \) \( \mu \)s. Data were taken at a different sample tuning from data in c–d. f. Rabi oscillations at different powers, indicated by the dashed lines in e. The data are shifted in phase by 0.45° for clarity. The solid curves are fits to a power-law decay. We obtain a minimum Rabi period of \( T_{\text{Rabi}} = 17 \) ns (Supplementary Information, section 5).
shows the measured phase as a function of $t_F$ and the gate drive power, $P_G$. We observe Rabi oscillations with a minimum period of 17 ns (Fig. 4f), consistent with an EDSR driving mechanism. These data show how the microwave field of the cavity is sensitive to the spin state of a single electron and that by using the qQED architecture quantum dot spin states may be coherently controlled and measured using microwave electric fields.

Long-distance coupling of spin qubits via a cavity will require a spin–cavity coupling rate that is larger than the cavity decay rate and the qubit decoherence rate. Although the method of spin-state read-out that we have demonstrated does not imply spin–cavity coupling, on the basis of our results we can estimate the effective spin–cavity coupling strength. From our measurements, we find that $g\mathcal{E}/\gamma = 0.2 \text{ MHz}$, which is four orders of magnitude larger than the coupling rate $g\mathcal{E}$ that would be obtained by coupling a single spin to the magnetic field of a microwave cavity. This spin–cavity coupling rate could be readily increased to $\sim 1 \text{ MHz}$ by increasing the cavity resonance frequency to $f_0 = 15 \text{ GHz}$ (which would proportionally increase both $g\mathcal{E}$ and $E_2$). Recent theoretical work also predicts an enhanced spin–cavity coupling for a single spin in a DQD biased at $\varepsilon = 0$ (ref. 29).

In addition to increasing the spin–cavity coupling rate, there is significant scope for improving the cavity decay rate and the qubit decoherence rate. Optimization of the resonator design will reduce the cavity decay rate to well below 1 MHz (Supplementary Information, section 2). There are several options for decreasing the qubit decoherence rate, which is at present limited by coupling to the nuclear spin bath. Dynamical decoupling has already been used to reduce the qubit decay rate to $\sim 1 \text{ MHz}$ in the InAs system. InAs could also be replaced by nuclear-spin-free Ge/Si core–shell nanowires where hole spin–orbit coupling is predicted to be large30. On the basis of our results, we anticipate that the strong-coupling regime for single spins can be reached, eventually allowing spin qubits to be interconnected in a quantum bus architecture.

METHODS SUMMARY

Samples were fabricated on high-resistivity, (100)-orientation silicon wafers with 250 nm of dry thermal oxide. Superconducting resonators were formed by first depositing a 100-mm-thick Nb layer followed by a 5-nm Au film without breaking vacuum. The Au film was removed, except in regions that were later contacted by either wire bonds or electron-beam lithography (EBL), by a chemical wet etch in a solution of hydrochloric acid and nitric acid. A half-wavelength resonator was defined using photolithography followed by etching in a solution of hydrochloric acid and nitric acid. A half-wavelength length of SiNx was deposited on top of the gate electrodes using plasma-enhanced chem-

Received 4 June; accepted 24 August 2012.

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Supplementary Information is available in the online version of the paper.

Acknowledgements Research at Princeton University was supported by the Alfred P. Sloan Foundation, the David and Lucile Packard Foundation, US Army Research Office grant W911NF-08-1-0189, DARPA QUEST award HR0011-09-1-0007 and the US National Science Foundation through the Princeton Center for Complex Materials (DMR-0981980) and CAREER award DMR-0846341. J.M.T. acknowledges support from ARO MURI award W911NF-09-1-0406.

Author Contributions K.D.P. fabricated the sample and performed the measurements. K.D.P., L.W.M. and A.A. developed the resonator fabrication and measurement processes. K.D.P., M.D.S. and M.J. developed the nanowire device fabrication processes. M.D.S. grew the nanowires. J.M.T. developed the theory for the experiment. K.D.P. and J.R.P. wrote the paper with input from the other authors. J.R.P. planned the experiment.

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