Gate-defined quantum dots were recognized early on as a promising platform for quantum information and many materials have been investigated as hosts for the quantum dots. Initial research mainly focused on the low-disorder semiconductor gallium arsenide. Steady progress in the control and understanding of this system culminated in the initial demonstration and optimization of spin qubit operations and the realization of rudimentary analogue quantum simulations. However, the omnipresent hyperfine interactions in group III–V materials seriously deteriorate the spin coherence. Considerable improvements to the coherence times could be achieved by switching to the group IV semiconductor silicon, in particular when defining spin qubits in an isotopically purified host crystal with vanishing concentrations of non-zero nuclear spins. This enabled single-qubit rotations with fidelities beyond 99.9% and the execution of two-qubit logic gates with fidelities up to 98%, underlining the potential of spin qubits for quantum computation. Nevertheless, quantum dots in silicon are often formed at unintended locations, and control over the tunnel coupling determining the strength of two-qubit interactions is limited. Moreover, the absence of a sizable spin–orbit coupling for electrons in silicon requires the inclusion of microscopic components such as on-chip striplines or nanomagnets close to each qubit, which complicates qubit definition for electrons in silicon, this is absent for holes, and excited states can be well separated in energy. In silicon, unfavourable band alignment prevents strain engineering of low-disorder quantum wells for holes, restricting experiments to metal–oxide–semiconductor structures. Research on germanium has mostly focused on self-assembled nanowires and has demonstrated single-shot spin readout and coherent spin control. However, strained germanium can reach hole mobilities of $\mu > 10^6 \text{cm}^2 \text{V}^{-1} \text{s}^{-1}$, and undoped germanium quantum wells were recently shown to support the formation of gate-controlled hole quantum dots. Now, the crucial challenge is the demonstration of coherent control in this platform and the implementation of qubit–qubit gates for scalable quantum information with holes.

Here we make this step and demonstrate single- and two-qubit logic with holes in planar germanium. We fabricate devices on silicon substrates, using standard manufacturing materials. We grow undoped strained germanium quantum wells, measured to have high hole mobilities $\mu > 5 \times 10^6 \text{cm}^2 \text{V}^{-1} \text{s}^{-1}$ and a low effective hole mass $m_h = 0.09 m_e$, extrapolated to reach $m_h = 0.05 m_e$ at zero density, with $m_e$ the electron rest mass. This allows us to define quantum dots of comparatively large size, and we find excellent control over the exchange interaction between the two dots. We operate in a multi-hole mode, reducing challenges in tuning and characterization, which is advantageous for scaling. We make use of the spin–orbit interaction for qubit driving and perform single-qubit rotations at frequencies exceeding 100 MHz. This advantage of fast driving becomes further apparent in coherently accessing the Hilbert space of a two-qubit system.
For example, in silicon the execution of a controlled NOT (CNOT) gate implemented with an on-chip stripline has shown using microsecond long pulses, and this timescale can be reduced to 0.2–0.5 μs by incorporating nanomagnets. Here we demonstrate that the spin–orbit coupling of holes in germanium together with the sizable exchange interaction enables a CNOT within 75 ns.

A scanning electron microscope image of the germanium two-qubit device is shown in Fig. 1a. To accumulate holes and define quantum dots, we used a high-mobility Ge quantum well and controlled by the electric gates. The direction of the external field, indicated by the black arrow. Schematic cross-section of the system, where quantum dots are formed below plunger gates P1 and P2, while the different tunnelling rates can be controlled by barrier gates BS, BD, and BC. c. Transport current through the double dot as a function of plunger gate voltages for weak (top) and strong (bottom) interdot coupling, mediated by a virtual tunnel gate. d. Charge stability diagram of the qubit operation point, where the dashed lines correspond to the charge transitions. The detuning axis is indicated by the dotted line, with label R corresponding to the qubit readout point. To allow coherent control of the isolated spin states, a two-level voltage pulse on gates P1 and P2 is used to detune the dot potentials and prevent tunnelling to and from the dots during the manipulation phase (label M). c. Transport current through the double dot as a function of plunger gate voltage for positive (left) and negative (right) bias. Pauli spin blockade becomes apparent from the suppression of the transport current for the positive bias direction, up to the single–triplet energy splitting of $E_g = 0.6$ meV. f. Illustration of the energy landscape in our double-quantum dot system. g. Resonance frequency, $f_{\text{res}}$, of the two qubits as a function of the external magnetic field, showing the individual qubit resonances.

To determine the control fidelity, which describes the accuracy of our quantum gates, we implement randomized benchmarking of the single-qubit Clifford group (Fig. 2c). The measured decay curve of the qubit state as a function of sequence length $m$ is shown in Fig. 2d, from which we extract a single-qubit control fidelity of $F = 99.3\%$, using gate times $t_g = 20$ ns and $t_{\pi/2} = 10$ ns. If $F > 99\%$, with the infidelity for π/2 gates being approximately twice as low as for the π gates, on account of the difference in pulse length.

We extensively characterize the coherence in our system at an exchange coupling of $J/h = 20$ MHz and find $T_2^{\text{coh}} = 833$ ns and $T_1 = 419$ ns, which can be extended by performing a Hahn echo to $T_2^{\text{coh}} = 1.9$ μs and $T_1 = 0.8$ μs (data in Extended Data Fig. 7), as indicated in Fig. 2f. These coherence times compare favourably to $T_2^{\text{coh}} = 130$ ns for germanium hut wires and $T_2^{\text{coh}} = 270$ ns for holes in silicon. Electrons in GaAs have even shorter dephasing time, with $T_2^{\text{coh}} = 10$ ns. The limited $T_2^{\text{coh}}$ in GaAs is due to hyperfine interactions,
which can be mitigated to a large extent by using nuclear notch filtering\textsuperscript{29}, leading to $T_\text{2} = 800$ µs. This source of dephasing can be avoided altogether by using group IV materials with nuclear spin-free isotopes\textsuperscript{30}. This has led to $T_\text{2} = 28$ ms for electrons in isotopically purified silicon\textsuperscript{13}, and isotopic purification may also increase the quantum coherence in germanium. Furthermore, we observe spin lifetimes of $T_{1,\text{1}} = 9$ µs and $T_{1,\text{2}} = 3$ µs. We have found that these lifetimes increase exponentially when lowering the tunnel coupling between each qubit and its respective reservoir (Extended Data Fig. 8), and relaxation times of $T_\text{1} > 100$ µs have been reported for germanium nanowires\textsuperscript{19,20}, both giving good prospects for increasing the relaxation time by closing the reservoir barrier during operation.

When the manipulation of both qubits is combined, the coupling of the two qubits (exchange interaction $J$) becomes apparent. As is illustrated in Fig. 3a, the resonance frequency of each of the qubits is shifted when the other qubit is prepared in its $|1\uparrow\rangle$ state. The strength of this interaction depends on the inter-dot tunnel coupling $t_{\text{cd}}$ as well as the detuning $\varepsilon$ of the dot potentials. By changing the amplitude of voltage

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**Fig. 2** Qubit control, gate fidelity and quantum coherence of planar germanium qubits. a, Measurement sequence used for the Rabi driving measurements. Measurement cycles with EDSR pulses are alternated with reference cycles without a microwave tone, allowing an efficient background current subtraction. Each cycle is repeated $N$ times, such that measurement and reference cycles alternate at a typical lock-in frequency of $f_{\text{res}} = 89.75$ Hz. b, Colour map of the differential bias current $\Delta I_{\text{SD}}$ as a function of microwave pulse time $t_{\text{p}}$ and power $P$, where clear Rabi rotations on Q1 can be observed. a.u., arbitrary units. c, Schematic illustration of the (interleaved) randomized benchmarking sequence applied to Q1. C corresponds to a single Clifford gate, with $m$ being the total number of applied random Clifford gates. d, Differential bias current as a function of $m$ for the randomized benchmarking sequence on Q1. The extracted control fidelity is $F_c = (99.3 \pm 0.03)\%$. e, Gate fidelities for the $\pi$ and $\pi$/2 gates. Error bars correspond to 1σ. f, Spin coherence and lifetime times for Q1 and Q2. Error bars correspond to 1σ.

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**Fig. 3** Tunable exchange coupling and operation at the charge symmetry point. a, Illustration of the relevant energy levels in our hole double quantum dot with zero (green) and finite (black) exchange coupling between the dots. Six energy levels are considered: the four different (1,1)-charge states as well as the (2,0), and (0,2), singlet charge states in which both holes occupy the same quantum dot. Four individual transitions can be driven, corresponding to the conditional rotations of the two-qubit system. The size of the exchange interaction is equal to $h = f_1 f_2 - f_1 f_3 - f_2 f_3$. b, Measurement pulse cycles used to map out the exchange splitting of Q1 (top) and Q2 (bottom). As a result of the demodulation of the alternating cycles, transition $f_{\text{los}}$ gives a negative signal and transition $f_{\text{ds}}$ results in a positive signal. c, d, EDSR spectra of Q1 (c) and Q2 (d) as a function of the detuning $\varepsilon$. The exchange splitting can be tuned to a minimum at $\varepsilon = 0$ and increases closer to the $(m, n) - (m+1, n-1)$ and $(m, n) - (m-1, n+1)$ charge transitions. e, Exchange interaction as a function of $\varepsilon$ as extracted from c, d. Fitting the exchange coupling yields an interdot tunnel coupling $t_{\text{cd}} = 1.8$ GHz and charging energy $U = 1.46$ meV. f, The interdot tunnel coupling can also be controlled by gate BC. Changing the potential on this gate, while keeping $\varepsilon = 0$, allows good control over the exchange interaction between the two qubits. g, Coherence time $T_2^*$ of both qubits as a function of detuning voltage $V$. When the slope of the resonance line is equal to zero, the qubit is expected to be, to first order, insensitive to charge noise. Solid lines indicate fits of the data to $\frac{1}{T_2^*} = \frac{1}{T_2^*} + \frac{1}{T_2^*}\text{res} + \alpha a$, with $a$ as a scaling factor. It can be observed that $T_2^*$ is indeed longest when the slope of the resonance line is closest to zero. Error bars correspond to 1σ. h, Resonance frequency of transition $f_1$ and $f_2$, as a function of detuning voltage.
rotations can naturally be performed by selectively driving each of the four transitions. A CX gate is achieved at $t_{CX} = t_{\pi}$ on $f_3(f_4)$. A small off-resonant driving effect can be observed, which we mitigate by tuning $t_{CX} = t_{\text{resonant}} = t_{4\pi,\text{off-resonant}}$.

control qubit with applied phase corrections. We observe a larger signal amplitude on Q1 after 0 and 4π rotations on Q2 as compared with a 2π rotation on Q2. This 4π periodicity is in agreement with fermionic statistics and suggests an echoing pulse correcting residual environmental coupling. The full π phase shift on Q2 for a conditional 2π rotation on Q1, as a result of the $\theta/2$ phase that is accumulated by the control qubit, demonstrates the application of a coherent CX gate.

The demonstration of a universal gate set with all-electrical control and without the need of any microscopic structures offers good prospects to scale up spin qubits using holes in strained germanium. The hole states do not suffer from nearby valley states, and the quantum dots are contacted by superconductors that may be shaped into microwave resonators for spin–photon coupling. This provides opportunities for a platform that can combine semiconducting, superconducting and topological systems for hybrid technology with fast and coherent control over individual hole spins. Moreover, the demonstrated quantum coherence and level of control make planar germanium a natural candidate to engineer artificial Hamiltonians for quantum simulation, going beyond classically tractable experiments.

**Online content**

Any methods, additional references, Nature Research reporting summaries, source data, extended data, supplementary information, acknowledgements, peer review information; details of author contributions and competing interests; and statements of data and code availability are available at https://doi.org/10.1038/s41586-019-1919-3.

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Methods

Fabrication process
Our Ge/SiGe heterostructures are grown on a 100-mm n-type Si(001) substrate, using an Epsilon 2000 (ASM) RP-CVD reactor, as described in ref. 23. The device’s Ohmic contacts and the electrostatic gates are defined by electron beam lithography, electron beam evaporation and lift-off of Al and Ti/Pd. Ohmic contacts consist of a 20-nm-thick Al layer, followed by a 17-nm-thick Al2O3 gate dielectric grown by atomic layer deposition at 300 °C. Next, the first layer of Ti/Pd (40 nm) gates is deposited, followed by 17 nm of Al2O3, and the second layer of overlapping Ti/Pd (40 nm) gates. Finally, vias connecting the lower gate layer are etched through the top Al2O3 layer, followed by the deposition of 1-μm-thick Al99Si1 bond pads to protect the device during bonding.

Experimental set-up
All measurements are performed in a Bluefors dry dilution refrigerator with a base temperature of $T_{\text{base}} = 10$ mK. Constant d.c. voltages are applied with battery-powered voltage sources, and the voltages on gates P1 and P2 are combined with an a.c. voltage by a bias-tee with a cut-off frequency of 3 Hz. The a.c. voltage for gate P1 is generated by an arbitrary waveform generator (AWG) Tektronix AWG5014C, combined with a microwave signal generated by a Keysight PSG8267D vector source. The a.c. voltage for gate P2 is solely the waveform generated by the AWG. EDSR pulses are generated by the PSG8267D using the internal IQ-mixer, driven by two output channels of the AWG. Both qubits can be addressed by setting the vector source to an intermediate frequency of typically $f_{\text{int}} = 2.56$ GHz, and IQ-mixing this with a (co)sine wave generated on channels 3 and 4 of the AWG. Because the on/off ratio of the IQ-modulation of our vector source is only 40 dB and small residual output power may lead to added infidelity, we use digital pulse modulation in series with the IQ-modulation. The pulse modulation is driven by the AWG and is turned on 15 ns before the first pulse and turned off 7 ns after the last pulse in the sequence, resulting in a total suppression of 120 db when the source is off.

We typically apply a source–drain bias voltage of $V_{\text{SD}} = 0.3$ mV and measure the current through the device using an in-house-built transimpedance amplifier, after which the signal is low-pass filtered at 10 kHz and measured using an Stanford Research SR830 lock-in amplifier, as described in Methods section ‘Sequence details’ below.

Virtual gates
To allow independent control over the tunnel coupling and the charging occupation of the double dot system, we make use of virtual gates. When changing the different barrier gate voltages, linear corrections are applied to the device’s plunger gates to correct for the cross-capacitance between the different gates. These coefficients are obtained from the relative slopes of the charge-addition lines with respect to the different device gates and normalized to the respective plunger gate coefficient. We write

$$ \begin{bmatrix} V_{P1} \\ V_{P2} \end{bmatrix} = \begin{bmatrix} a_{P1,P1} & a_{P2,P1} & a_{BC,P1} & a_{BR1,P1} & a_{BR2,P1} \\ a_{P1,P2} & a_{P2,P2} & a_{BC,P2} & a_{BR1,P2} & a_{BR2,P2} \end{bmatrix} \begin{bmatrix} P1 \\ P2 \\ BC \\ BR1 \\ BR2 \end{bmatrix} \delta $$

with VP1 and VP2 the virtual plunger gates, and P1, P2, BC, BR1 and BR2 the different physical device gates as indicated in Fig. 1a. The virtual gate matrix describes the different couplings and is given by

$$ \begin{bmatrix} a_{P1,P1} & a_{P2,P1} & a_{BC,P1} & a_{BR1,P1} & a_{BR2,P1} \\ a_{P1,P2} & a_{P2,P2} & a_{BC,P2} & a_{BR1,P2} & a_{BR2,P2} \end{bmatrix} = \begin{bmatrix} 1 & 0.8 & 0.35 & 0 \\ 0 & 1 & 0.8 & 0.4 \end{bmatrix} $$

We do not correct for the crosstalk between the two plunger gates, such that $a_{P1,P2} = a_{P2,P1} = 0$. The crosstalk between the quantum dot and the reservoir barrier of the other dot is negligible because of their physical separation. Furthermore, it can be observed that the coupling of the centre barrier to both dots is approximately twice as strong as the reservoir barriers as a direct effect of its increased size.

Sequence details
To improve the quality of the transport measurements, we establish a lock-in measurement scheme in which the measurement of interest is alternated with a reference measurement to account for slow variations in the transport current through the device, as well as temperature-dependent drifts in our transimpedance amplifier, as is illustrated in Extended Data Fig. 1. The measurement cycle, consisting of the readout as well as the manipulation phase, typically has a length of $t_{\text{ cyc} \text{ le}} = 1$ μs. With the AWG, we generate a waveform that repeats the measurement cycle $N$ times, followed by $N$ repetitions of a similar reference measurement, with $N$ chosen such that these cycles alternate at a lock-in frequency of $f_{\text{lock-in}} = 89.75$ Hz. The measured transport current is then demodulated by a lock-in amplifier, using a reference signal generated by the AWG. As a result, the lock-in output signal will be directly related to the difference in transport current between the measurement and the reference cycle. During the readout, no differential current is observed when the qubits are in their $|\uparrow\downarrow\rangle$ ground state, whereas a signal of typically $\Delta I_{\text{ref}} = 0.3$ pA is measured for all other spin configurations and a total cycle length of $t_{\text{cyc}} = 900$ ns. This is in good agreement with a bias current $\Delta I = 2e/t_{\text{cyc}} = 0.4$ pA, as expected for the random loading of a hole spin.

For a Rabi experiment, the measurement cycle contains a single microwave pulse of duration $t_{\text{p}}$, whereas the reference cycle has no pulses. In the case of a Ramsey experiment, both the measurement and reference cycle contain a π/2 pulse, a wait $\tau$ and a final π/2 pulse, but in the reference cycle the final π/2 pulse is phase-shifted by $\phi = \pi$. This will result in an opposite projection for the two measurements and thereby maximum differential signal. For the randomized benchmarking, a similar scheme is used (see Fig. 2a), where the reference pulse in the measurement cycle is chosen to project to the spin-up state, while the recovery pulse in the reference cycle projects to the spin-down state, resulting in an exponential decay towards $\Delta I_{\text{ref}} = 0$. Each data point is averaged over approximately $10^3$ repetitions of 1,500 randomly drawn gate sequences. Finally, for the exchange measurements, we alternate a measurement cycle where we apply a $\pi$ and $-\pi$ pulse to Q1 (Q2) before and after the probing pulse respectively, with a reference cycle where Q1 (Q2) is not pulsed. When the probing pulse is off-resonant with both resonance frequencies, the measurement cycle gives effectively no rotation of Q1 (Q2) and the reference cycle does not result in any rotation. As a result the demodulated signal will be zero. When the probing pulse frequency is on resonance with the unprepared resonance frequency $f_1 (f_2)$, the measurement cycle will still be an effective zero rotation on Q1 (Q2) due to the selective driving of $f_1 (f_2)$ and thus give no signal. The reference cycle will now result in a π rotation on Q2 (Q1) and will therefore give a signal, resulting in a negative demodulated signal. In the case where the probing pulse is resonant with the prepared resonance line $f_1 (f_2)$, the measurement cycle will generate a signal whereas the reference cycle will give no signal, thus resulting in a positive demodulated signal. All different pulse cycle configurations and the respective qubit projections are illustrated in Extended Data Fig. 2b.

Phase corrections for pulsing
We observe a shift of the resonance frequency of the qubits as a function of the microwave driving power. We attribute this to a rectification of the microwave signal, resulting in a d.c. voltage pulse which can modulate the resonance frequency through the SOC and exchange interaction. As a result of the shift during the pulsing, each qubit picks up a phase when it is idling, as well as an additional phase due to the pulses on the other qubit. We can calibrate these frequency shifts and correct all following pulses to counteract this phase shift.

To probe the effect of all possible pulses on all possible resonances, we perform an extended Ramsey experiment. We prepare a pulse sequence
consisting of two $\pi/2$ pulses with a test gate (each of the four resonance lines, as well as idling) and $\pi$ phase-shifted test gate in between, as indicated in Extended Data Fig. 9. For the experiment on $f_2$ and $f_4$, we add an additional preparation and projection pulse at the start and end respectively, as indicated in grey in Extended Data Fig. 9. The back-and-forth rotation on the test gate cancels any driving effects, as well as the $\theta/2$ phase picked up due to the conditional rotation, and leaves us with only the detuning phase. We now plot the transport current $\Delta I_{SD}$ as a function of the phase $\phi$ of the second $\pi/2$-pulse, as well as the length of the test gate. As a result of the frequency shift caused by the test gate, we observe a phase shift increasing linearly with the length of the test gate. We fit this phase shift for each gate, and we apply a correction to all following gates. Extended Data Fig. 9 shows the phase evolution for all test gates on all four resonance lines, both without corrections (Extended Data Fig. 9a), as well as with corrections applied (Extended Data Fig. 9b).

**Data availability**

All data underlying this study are available from the 4TU ResearchData repository at https://doi.org/10.4121/uuid:95bc1f2e-0218-4c55-8e5b-2b59e8fcc5e6.

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**Author contributions**

N.W.H. and D.P.F. performed the experiments. N.W.H. fabricated the device. A.S. and G.S. supplied the heterostructures. N.W.H., D.P.F. and M.V. wrote the manuscript with the input of all other authors. M.V. conceived and supervised the project.

**Competing interests**

The authors declare no competing interests.

**Additional information**

**Supplementary information** is available for this paper at https://doi.org/10.1038/s41586-019-1919-3.

**Correspondence and requests for materials** should be addressed to M.V.

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Extended Data Fig. 1 | Instrumentation set-up for the lock-in transport measurements. Illustration of the set-up and relevant signals for the lock-in transport measurements. The AWG is used to generate alternating pulse cycles consisting of a repeated measurement and a repeated reference. The signal is demodulated in a lock-in amplifier to give a direct measure of the difference between the two measurements and subtract slow variations in the transport signal.
Extended Data Fig. 2 | Pulse cycles used for the transport measurements.

a Pulse cycles used for the randomized benchmarking experiments. The measurement pulse cycle consists of \( m \) gates randomly drawn from the Clifford group \( C_{\text{rand}} \) and a final Clifford gate projecting the qubit onto the spin-up state. The reference pulse cycle consists of the same \( m \) Clifford gates and a different final Clifford gate projecting the qubit onto the spin-down state. Each cycle is repeated \( N \) times, and a series of typically \( k = 50 \) independent randomly drawn measurement and reference pulse cycles are alternated. These \( k = 50 \) different draws are thus hardware-averaged on the lock-in amplifier, and the entire experiment is repeated and averaged 30 times, yielding a total approximate \( 10^5 \) repetitions of 1,500 different randomly drawn Clifford sequences of length \( m \). An example of the qubit evolution for each pulse cycle is plotted on the Bloch sphere below.

b Pulse cycles used for the exchange mapping experiments. The measurement pulse cycle consists of a broad preparation and restoring pulse at frequency \( f_3 \) (\( f_1 \)), around a probing pulse at frequency \( f_{\text{prb}} \). The reference pulse cycle consists solely of the probing pulse at \( f_{\text{prb}} \). The qubit evolutions for the different resonance conditions are plotted on the Bloch sphere and illustrate the different signals measured in Fig. 4c, d.
Extended Data Fig. 3 | Demonstration of qubit operation at a second hole occupancy. **a**, Charge stability diagram showing the \((m, n)\) hole occupancy used during all experiments in the main text, as well as the \((m, n + 1)\) occupancy for which we observe PSB as well. For an unpolarized filling of the quantum dots, one expects an alternating suppression of the transport current due to PSB, as spin blockade occurs only when an orbital level is fully occupied. However, the spin-filling for holes is known to be highly polarized\(^3\), and therefore PSB can occur in sequential quantum dot fillings. **b**, Coherent Rabi oscillations measured in the \((m, n + 1)\) occupancy. A slight linear offset is observed for Q1, which can be attributed to the microwave power. We note that, for the same microwave power, the Rabi frequency of Q2 in the \((m, n + 1)\) occupancy is increased substantially compared to the \((m, n)\) filling. We attribute this to the hole being in a different orbital, where the effective SOC may be different.
Extended Data Fig. 4 | Qubit resonance frequencies as a function of magnetic field. Colour plot indicating the transport current $\Delta I$ through the double dot system, as a function of external magnetic field $B_0$ and the frequency $f$ of the applied microwave signal. We have numerically subtracted the mean of each row and column in each of the three individual colour plots, to account for the slow drifts in transport current, as well as the line resonances in our fridge cabling. The two bright lines indicate an increase in the transport current due to the microwave rotating either spin and thus lifting PSB.
Extended Data Fig. 5 | Temporal dependence of the resonance frequency.
We track the resonance frequency of both Q1 and Q2 over the time of approximately 110 h. We observe that the qubit frequency remains remarkably stable over this period, but do observe discrete, uncorrelated steps in the resonance frequency of both qubits. The resonance frequency of Q1 only shows steps of \( \Delta f \approx 2 \) MHz between two distinct levels, whereas for Q2 we observe steps of \( \Delta f \approx 1 \) MHz and \( \Delta f \approx 2 \) MHz, between three different levels, as also becomes apparent from the histogram. The origin of these steps could be, for example, the slow loading and unloading of charge traps, which manipulate the qubit resonance frequency through the change in electric field, or hyperfine coupling to a nearby nuclear spin.
Extended Data Fig. 6 | Magnetic field dependence of the driving speed of Q1.

(a, b) Rabi frequency dependence on the applied microwave power $P$ in arbitrary units, for $B_0 = 0.5$ T (a) and $B_0 = 1.65$ T (b). Multiple mechanisms can be at play for the EDSR driving of the spins and these are typically all linearly dependent on $B_0$. As a result of this, considerably higher driving frequencies can be reached at higher magnetic fields. We note that the exact microwave power cannot be compared between the two measurements, owing to the strong frequency dependence of the attenuation of our fridge lines.
Extended Data Fig. 7 | Relaxation, dephasing and coherence times. We perform a Ramsey experiment, in which two \( \pi/2 \) pulses are separated by time \( \tau \), during which the qubit will evolve as a result of the implemented detuning. We fit the decay of the observed oscillations to
\[
I(\tau) = a \cos(2\pi \Delta f \tau + \phi) \exp[-(\tau/T^*_{\Delta})^2],
\]
with \( a \) a scaling factor, \( \Delta f \) the detuning and \( \phi \) a phase offset, and find a spin coherence time of \( T^*_{\Delta} = 833\,\text{ns} \) and \( T^*_{\Delta} = 419\,\text{ns} \) and decay coefficients of \( \alpha^*_1 = 1.2\pm0.2 \) and \( \alpha^*_2 = 1.5\pm0.2 \), for Q1 and Q2, respectively. The spin coherence can be extended by performing a Hahn echoing sequence, consisting of \( \pi/2, \pi \) and \( \pi/2 \) pulses separated by waiting times \( \tau \). Fitting the observed decay as a function of the total waiting time \( 2\tau \) to a power law \( I(\tau) = a \exp[-(2\tau/T^*_{H})^2] \), we find extended coherence times of \( T^*_{H} = 1.9\,\mu\text{s} \) and \( T^*_{H} = 0.8\,\mu\text{s} \) and decay coefficients of \( \alpha^*_1 = 1.5\pm0.1 \) and \( \alpha^*_2 = 2.5\pm0.3 \), for Q1 and Q2, respectively. Finally, we perform a measurement of the spin lifetime by applying a single \( \pi \) pulse, after which we wait for a time \( \tau \). We fit the decay to \( I(\tau) = \exp[-(\tau/T_1) \right) \) and find lifetimes of \( T_1 = 9\,\mu\text{s} \) and \( T_1 = 3\,\mu\text{s} \).
Extended Data Fig. 8 | Relaxation time \(T\) as a function of gate voltage on the tunnel barriers between dot and reservoir. a, b. The relaxation time \(T\) of the dots increases approximately exponentially as a function of the respective dot-reservoir gate voltage, for Q1 (a) as well as for Q2 (b). The relaxation time of Q1 increases exponentially from \(T < 1\) μs to \(T > 10\) μs, and a similar scaling is observed for Q2. For even smaller dot-reservoir couplings, the transport signal drops below our measurement limit, but switching to charge sensing could allow a further increase in \(T\).
Extended Data Fig. 9 | Phase corrections on the qubits. 

**a**, Extended Ramsey experiment on each of the four resonance line, using five different test gates between the π/2 pulses to observe the effect on the resonance frequency. A linear phase shift as a function of test gate pulse length $\tau$ can be observed for some lines, indicating a frequency shift during the pulsing. **b**, We compensate for this effect by performing a software update of $\delta \phi = \delta \nu \tau$ to each additional pulse, with $\delta \nu$ the frequency shift of the qubit as a result of the microwave signal.