A singlet-triplet hole spin qubit in planar Ge

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Spin qubits are considered to be among the most promising candidates for building a quantum processor. Group IV hole spin qubits are particularly interesting owing to their ease of operation and compatibility with Si technology. In addition, Ge offers the option for monolithic superconductor–semiconductor integration. Here, we demonstrate a hole spin qubit operating at fields below 10 mT, the critical field of Al, by exploiting the large out-of-plane hole g-factors in planar Ge and by encoding the qubit into the singlet-triplet states of a double quantum dot. We observe electrically controlled g-factor difference-driven and exchange-driven rotations with tunable frequencies exceeding 100 MHz and dephasing times of 1 μs, which we extend beyond 150 μs using echo techniques. These results demonstrate that Ge hole singlet-triplet qubits are competing with state-of-the-art GaAs and Si singlet-triplet qubits. In addition, their rotation frequencies and coherence are comparable with those of Ge single spin qubits, but singlet-triplet qubits can be operated at much lower fields, emphasizing their potential for on-chip integration with superconducting technologies.

Holes in Ge have emerged as one of the most promising spin qubit candidates because of their particularly strong spin–orbit coupling, which leads to record manipulation speeds and low dephasing rates. In addition, the spin–orbit coupling, together with the low effective mass, relaxes fabrication constraints, and larger quantum dots can be operated as qubits without the need for microstrips and micromagnets. In only three years, single Loss–DiVincenzo qubit, two-qubit and most recently even four-qubit devices have been demonstrated. Here we show that by implementing Ge hole spin qubits in a double quantum dot device, they have the further appealing feature that operation below the critical field of aluminium becomes possible.

In order to realize such a qubit, a strained Ge quantum well (QW) structure with a hole mobility of 1.0 × 105 cm2 V–1 s–1 at a density of 9.7 × 1011 cm–2 was grown by low-energy plasma-enhanced chemical vapour deposition. Starting from a Si wafer, a 10-μm-thick strain-relaxed Si0.75Ge0.25 virtual substrate (VS) is obtained by linearly decreasing the Ge content during the epitaxial growth. The ~20-nm-thick strained Ge QW is then deposited and capped by 20 nm of Si0.7Ge0.3. In Fig. 1a we show the aberration-corrected high-angle annular dark-field (HAADF) scanning transmission electron microscopy image of our heterostructure. The HAADF Z contrast clearly draws the sharp interfaces between the QW and the top and bottom barriers. In addition, X-ray diffraction measurements highlight the lattice matching between the virtual substrate and the QW (Fig. 1b). Holes confined in such a QW are of mainly the heavy-hole (HH) type because compressive strain and confinement move light-holes (LHs) to higher hole energies. The related Kramers doublet of the spin Sz = ±3/2 states therefore resembles an effective spin-1/2 system, |↑⟩ and |↓⟩.

In a singlet-triplet qubit, the logical quantum states are defined in a two-spin-1/2 system with total spin along the quantization axis Sz = 0 (refs. 11,12). This is achieved by confining one spin in each of two tunnel coupled quantum dots, formed by depletion gates (Fig. 1c). We tune our device into the single hole transport regime, as shown by the stability diagram in Fig. 2a where the sensor dot reflected phase signal (ϕref) is displayed as a function of the voltage on gates L and R (Methods and Supplementary Figs. 7 and 8). Each Coulomb blocked region corresponds to a fixed hole occupancy, and is labelled by (NL, NR), with NL (NR) being the equivalent number of holes in the left (right) quantum dot; interdot and dot to-lead charge transitions appear as steep changes in the sensor signal. Fast pulses are applied to the outer barrier gates LB and RB, which eases pulse calibration since the cross capacitance to the opposite dot is negligible. By pulsing in a clockwise manner along the empty–separate–measure (E–S–M) vertices (Fig. 2b), we observe a triangular region leaking into the upper-left Coulomb blocked region. Such a feature identifies the metastable region where the Pauli spin blockade occurs: once initialized in E, the pulse to S loads a charge, and the spins are separated, forming either a spin singlet or a triplet. At the measurement point M within the marked triangle, the spin singlet state leads to tunnel events, while the triplet states remain blocked, which allows spin-to-charge conversion. We repeat the experiment with a counter-clockwise ordering (E–M–S) and no metastable region is observed, as expected (Fig. 2a was acquired while pulsing in the counter-clockwise ordering).

We thus consider the interdot line across the detuning (e) axis of Fig. 2a equivalent to the (2,0) ↔ (1,1) effective charge transitions. The system is tuned along the detuning axis from (2,0) to (1,1) by applying opposite pulses at radio frequency (rf) of amplitude Ve on LB and RB: e = Vref√α2ℓLB + α2ℓRB (Supplementary Fig. 7), where αℓLB (αℓRB) is the rf-lever arm of the left (right) barrier gate. The double quantum dot spectrum for a finite magnetic field B is reported in Fig. 2c (the triplet states T(2,0) lie high up in energy and are not shown; the model Hamiltonian is derived in Supplementary Materials).
Discussion Section 1). We set \( \epsilon = 0 \) at the \((2,0)\) ↔ \((1,1)\) crossing.

Starting from \((2,0)\), increasing \( \epsilon \) mixes \((2,0)\) and \((1,1)\) into two molecular singlets, the ground state \( S_0 \), neglected in the following, which are split at resonance by the tunnel coupling \( 2\sqrt{2}\mu \). The triplet states are almost unaffected by changes in \( \epsilon \). We define the exchange energy \( J \) as the energy difference between \( S = \frac{1}{\sqrt{2}}(|\uparrow\downarrow⟩ - |\downarrow\uparrow⟩) \) and the unpaired triplet \( T_0 = \frac{1}{\sqrt{2}}(|\uparrow\uparrow⟩ + |\downarrow\downarrow⟩) \). At large positive detuning, \( J \) drops due to the decrease of the wavefunction overlap of the two separated holes. Importantly, different \( J \) factors for the left \( g_L \) and the right \( g_R \) dot result in four \((1,1)\) states: two polarized triplets, \( |↑↑⟩ \) and \( |↑↓⟩ \), and two anti-parallel spin states, \( |↑↓⟩ \) and \( |↓↑⟩ \) split by the Zeeman energy difference \( \Delta E_Z = \Delta g\mu_B B \), where \( \Delta g = [g_L - g_R] \), \( \mu_B \) is the Bohr magneton and \( B \) is the magnetic field applied in the out-of-plane direction. However, as noticed later, even at large positive \( \epsilon \), a residual \( J \) persists, which leads to the total energy splitting between \( |↑↓⟩ \) and \( |↓↑⟩ \) being \( E_{\text{tot}} = \sqrt{(J(\epsilon))^2 + (\Delta g\mu_B B)^2} \).

By applying a pulse with varying \( \epsilon \) (Fig. 2d) and stepping the magnetic field, we obtain the plot in Fig. 2e drawing a funnel. The experiment maps out the degeneracy between \( J(\epsilon) \) and \( \Delta g\mu_B B \), where \( E_Z = \frac{1}{2\sqrt{2}} \) is the Zeeman energy of the polarized triplets and \( \Delta g = g_L - g_R \). The doubling of the degeneracy point can be attributed to fast spin–orbit induced \( S - T_0 \) oscillations. At larger detuning, \( S - T_0 \) oscillations become visible.

The effective Hamiltonian of the qubit subsystem is

\[
H = \begin{pmatrix}
-\frac{J(\epsilon)\Delta g\mu_B B}{2} & \Delta g\mu_B B \\
\Delta g\mu_B B & 0
\end{pmatrix}
\]

in the \(|S \rangle \} \) basis, with \( J(\epsilon) \) being the detuning-dependent exchange energy, common to all \( S - T_0 \) qubits. Implementations of \( S - T_0 \) qubits in GaAs typically harvest the local field gradient induced by the nuclear Overhauser field to drive \( S - T_0 \) oscillations. Due to the near absence of nuclear spins in Si, only slow oscillations could be achieved in natural Si/SiGe structures. Hence, micromagnets have been successfully used to enhance and stabilize the magnetic field gradient. In Si metal–oxide–semiconductor devices, \( S - T_0 \) oscillations can be driven by spin–orbit induced \( g \)-factor differences in the two dots, and values of 20 MHz T\(^{-1}\) have been reported. Here, similarly, we realize \( S - T_0 \) oscillations through \( g \)-factor differences. However, we expect a larger \( \Delta g \) since our holes are of mainly HH character. Indeed, as shown below, \( g \)-factor differences exceeding 20 GHz T\(^{-1}\) can be obtained. Pulsing on \( \epsilon \) influences \( J \), and the ratio between \( J \) and \( \Delta g\mu_B B \) determines the rotation axis tilted by an angle \( \theta = \arctan \left( \frac{\Delta g\mu_B B}{J(\epsilon)} \right) \) from the \( Z \) axis. For large detuning, \( \theta \to 90^\circ \), corresponding to \( X \) rotations, while for small detuning, \( \theta \to 0^\circ \), enabling \( Z \) rotations.

A demonstration of coherent \( \Delta g \)-driven rotations at a centre barrier voltage \( V_C = 910 \text{ mV} \) is depicted in Fig. 3c with the pulse sequence shown in Fig. 3b. The system is first initialized in \((2,0)\) in a singlet, then pulsed quickly deep into \((1,1)\), where the holes are separated. Here the state evolves in a plane tilted by \( \theta \) (Fig. 3a,d). After a separation time \( \tau_s \), the system is brought quickly to the measurement point in \((2,0)\) where the Pauli spin blockade enables the distinction of triplet and singlet. Varying \( \tau_s \) produces sinusoidal oscillations with frequency \( f = \frac{1}{2} \sqrt{J(\epsilon)^2 + (\Delta g\mu_B B)^2} \), where \( h \) is the Planck constant. We extract \( \Delta g = 2.04 \pm 0.04 \text{ and } J(\epsilon = 4.5 \text{ meV}) \approx 21 \text{ MHz} \). We attribute the large \( \Delta g \) to the different quantum dot sizes, which directly affect the HH–LH splitting, thus determining the effective \( g \) factor. In addition, the different quantum dot charge occupations can lead to further \( g \)-factor differences. We approach frequencies of 100 MHz at fields as low as 3 mT. We observe similar values of \( \Delta g \) in the range of 1.0 to 2.7 in two additional devices with similar gate geometries (Supplementary Fig. 13). Figure 3f shows the
This is out of the scope of the present work, which focuses on the low-magnetic-field behaviour.

We, furthermore, observe a dependence of $\Delta g$ on the voltage on the centre barrier (CB) (Fig. 3g), confirming electrical control over the $g$ factors. As the voltage is decreased by 50 mV, $\Delta g$ varies from $\sim 1.5$ to more than 2.2, which conversely increases the frequency of $X$ rotations. Concurrently we measure a similar trend in $T_2^*$ reported at $B = 1$ mT in Fig. 3h; as the centre barrier is lowered, the coherence of the qubit is enhanced. The origin and consequences of this observation are discussed later.

Next, we demonstrate full access to the Bloch sphere, achieved by $Z$ rotations leveraging the exchange interaction. We change the pulse sequence (Fig. 4b) such that after initialization in a singlet, the system is pulsed to large detuning, but is maintained in this position only for $t = t_{\pi/2}$, corresponding to a $\pi/2$ rotation, bringing the system close to $|1\rangle$. Now we let the state evolve for a time $\tau$, at a smaller detuning, increasing $J$ and changing the rotation angle $\theta$ (Fig. 4d), before applying another $\pi/2$ rotation at high detuning and pulsing back to read-out. The state evolution on the Bloch sphere in Fig. 4a shows that full access to the qubit space can be obtained by a combination of appropriately timed pulses.

The resulting oscillation pattern is depicted in Fig. 4c. From the inferred frequency, we find the dependence of $J$ on $\epsilon$ and extract $t_{\pi}/h = 3.64$ GHz as a free fitting parameter. The extracted values of $J$ are plotted in Fig. 4e with the blue markers obtained from the exchange oscillation frequency. The green dots, on the other hand, correspond to $J(\epsilon) = E_2^* = \frac{2 e \delta\epsilon}{h}$ extracted from the tunnel experiment (Fig. 2e). We find that the two sets of data points coincide when $\Delta g = 11.0$. Together with the $g$-factor difference already reported, we obtain the two out-of-plane $g$ factors to be 4.5 and 6.5, comparable to previous studies. In Fig. 4f-g we plot $P_1$ as a function of separation time at different values of $\epsilon$. $P_1$ now oscillates between 0 and 1 due to the combination of $\pi/2$ pulses and free evolution time at lower detuning. From the fits (black solid lines) at different detunings, we extract $T_2^*$ as a function of $\epsilon$ (Fig. 4i). For low $\epsilon$ the coherence time is shorter than 10 ns, while it increases for larger $\epsilon$ and saturates at around 2 ns. This is explained by a simple noise model given by $\frac{1}{T_2^*} = \frac{\pi}{4} \left( \frac{\Delta E_{\text{zrms}}}{E_{\text{tot}}} \right)^2 + \frac{1}{T_2^*} \frac{\delta E}{E_{\text{tot}}} \frac{\Delta E_{\text{zrms}}}{E_{\text{tot}}}^2$, where $T_2^*$ depends on electric noise affecting $J$ and a combination of electric and magnetic noise affecting $\Delta E_{\text{zrms}}$

\[
\frac{1}{T_2^*} = \frac{\pi}{4} \left( \frac{\Delta E_{\text{zrms}}}{E_{\text{tot}}} \right)^2 + \frac{1}{T_2^*} \frac{\delta E}{E_{\text{tot}}} \frac{\Delta E_{\text{zrms}}}{E_{\text{tot}}}^2,
\]

\[
\frac{1}{T_2^*} = \frac{\pi}{4} \left( \frac{\Delta E_{\text{zrms}}}{E_{\text{tot}}} \right)^2 + \frac{1}{T_2^*} \frac{\delta E}{E_{\text{tot}}} \frac{\Delta E_{\text{zrms}}}{E_{\text{tot}}}^2,
\]

where $\delta E_{\text{zrms}}$ is the root mean square (r.m.s.) noise on detuning, and $\delta E_{\text{zrms}}^2$ describes the combination of electric noise on $\Delta g$ and magnetic noise affecting $B$. We assume $\Delta E_{\text{zrms}} \approx 0$ as we observe almost no change in $\Delta g$ with detuning (Supplementary Fig. 9). From the fit (dark red solid line), we find $\delta E_{\text{zrms}} = 7.59 \pm 0.35$ meV, in line with comparable experiments. In Fig. 4i, $\delta E_{\text{zrms}} = 1.78 \pm 0.10$ meV. Although $\delta E_{\text{zrms}}$ is much smaller than $\delta E_{\text{zrms}}$, we find that at large detuning, coherence is still limited by noise $\Delta E_{\text{zrms}}$ because $\frac{\delta E}{E_{\text{tot}}} \rightarrow 0$ (red and violet dashed lines in Fig. 4i). We attribute the magnetic noise to randomly fluctuating hyperfine fields caused by spin-carrying isotopes in natural Ge, but a distinction from charge noise affecting $\Delta g$ cannot be made here.

Equation (2) also gives insights into the trends observed in Fig. 3f-h. With $B$ we now affect $\Delta E_{\text{zrms}}$, and, thereby, its contribution to the total energy. The higher the ratio $\Delta E_{\text{zrms}}/E_{\text{tot}}$, the more the coherence is limited by this term, as confirmed by the drop of $T_2^*$ with magnetic field in Fig. 3f. Similarly one would expect that by increasing $\Delta g$, $T_2^*$ should be lower. But, as shown in Fig. 4h, the increasing $g$-factor difference is accompanied by an increase of the tunnel coupling by 2 GHz. Hence, the increase of $J$ is larger than the increase of $\Delta E_{\text{zrms}}$ at lower $V_{\text{CB}}$, and $\Delta E_{\text{zrms}}$ is reduced, leading
Fig. 3 | $\Delta g$-driven rotations. a, State evolution on the Bloch sphere. X rotations are controlled by $\Delta g$ and the applied magnetic field. The ideal rotation axis is depicted as a dark red arrow. The dashed purple trajectory corresponds to a perfect X rotation while the effective rotation axis is tilted by an angle $\theta$ from the Z axis due to a finite residual J (orange arrow pointing along the Z axis) resulting in the state evolution depicted by the solid purple curve. b, Pulse sequence used for performing the $\Delta g$-driven rotations. The colour gradient is as in Fig. 2d. After initialization in a singlet, the separation time $\tau_s$ is varied (dashed lines) while the amplitude is $\phi = 4.5$ meV. The system is then diabatically pulsed back to the measurement point. c, The $\Delta g$-driven oscillations as a function of magnetic field and separation time at $V_{cb} = 910$ mV. The average of each column has been subtracted to account for variations in the reflectometry signal caused by the magnetic field. A low (high) signal corresponds to a higher singlet (triplet) probability. Each point is integrated for 100 ms under continuous pulsing (Supplementary Fig. 17). d, $\theta = \arctan \left( \frac{\Delta \phi B_0}{2 \Delta g B_0} \right)$ versus magnetic field. The effective oscillation axis is magnetic-field dependent and approaches $80^\circ$ for $B = 5$ mT. e, Frequency of $\Delta g$-driven oscillations as a function of magnetic field (dark red dots). The black line is a fit to $f = \frac{1}{2} \sqrt{P^2 + (\Delta g B_0)^2}$ where we extract a $g$-factor difference $\Delta g = 2.04 \pm 0.04$ and a residual exchange interaction $J \approx 4.5$ meV $= 20 \pm 1$ MHz. We reach frequencies of 100 MHz at fields as low as 3 mT. f, Singlet probability $P_s$ as a function of $r_g$ at different $B$ fields for $V_{cb} = 910$ mV extracted through averaged single shot measurements (Supplementary Figs. 17 and 18). The solid lines are a fit to $P_s = A \cos(2\pi r_g \omega + \phi) \exp\left(-\left(1/T_2^*\right)^2\right) + C$, where $A$ is the oscillation amplitude and $C$ is an offset. Because of the tilted angle, $P_s$ oscillates only between 0.5 and 1.0. Moreover, we observe a further decrease in visibility at higher magnetic fields due to decay mechanisms during the read-out process$^{23}$. The extracted $T_2^*$ shows a magnetic field dependence explainable by equation (2). g, The $g$-factor difference as a function of the centre barrier voltage $V_{cb}$. By opening the centre barrier, the $g$-factor difference increases from 1.50 to 2.25. Error bars correspond to one standard deviation. h, $T_2^*$ versus $V_{cb}$. A near doubling in coherence time with lower centre barrier voltage is a consequence of an increased tunnel coupling (Fig. 4h) as explained in the main text. The error bars represent the confidence interval of the fits.

to a longer $T_2^*$. While $V_{cb}$ affects both $t_c$ and $\Delta g$, we see that $V_{lb}$ and $V_{rb}$, where $V_{lb}(rb)$ is the voltage on gate LB (RB), affect mostly $t_c$ and leave $\Delta g$ unaltered (Supplementary Fig. 10). This exceptional tunability enables electrical engineering of the potential landscape to favour fast operations without negatively affecting the coherence times, thus enhancing the quality factor of this qubit. We find a quality factor $Q = f \times T_2^*$ that increases with magnetic field, reaching $Q = 52$ at 3 mT (Supplementary Fig. 15). While the longest $T_2^*$ reported here is already comparable to electron single-shot qubits in natural Si (ref. 17), a reduction in the
Fig. 4 | Exchange-rotations at $B = 1 \text{ mT}$ and $V_{\text{cb}} = 910 \text{ mV}$. a, State evolution on the Bloch sphere. The purple arrows represent $\pi$ pulses applied at maximum detuning, while the red trajectory corresponds to the free evolution at smaller $\epsilon$. b, Pulse sequence used to probe $Z$ rotations. A $\pi^\text{z}$ pulse prepares the state close to the equator of the Bloch sphere, where it subsequently precesses under the influence of $J$ at different $\epsilon$ as schematically indicated by the dashed lines. Another $\pi^\text{z}$ pulse maps the final state on the qubit basis for read-out. The colour gradient is as in Fig. 2d. c, $Z$ rotations as a function of $\tau_z$ and $\epsilon$. The acquisition method is the same as in Fig. 3c. d, Rotation angle $\theta$ as a function of $\epsilon$ for $B = 1 \text{ mT}$ and $J$ extracted from c, e, $J/h = \sqrt{f(\epsilon)^2 - (\Delta g B/\hbar)^2}$ as a function of $\epsilon$ as extracted from the oscillation frequency $f$ (blue markers). Green dots correspond to the spin funnel (Fig. 2e) condition $J(\epsilon) = E_2^f$ with $\Sigma g = 11$, and the red dashed line is the best fit to $J(\epsilon) = \left| \frac{\delta}{2} - \sqrt{\frac{\epsilon^2}{4} + 2\epsilon^2} \right|$. f–g, $P_\text{r}$ as a function of $\tau_z$ for different $\epsilon$. An offset of +1 was added to $P_\text{r}$ for clarity. The pulse sequence adopted here increases the amplitude of oscillations as compared to Fig. 3l enabling full access to the Bloch sphere. At very low $\epsilon$, we observe the signal to chirp towards the correct frequency as a direct consequence of a finite pulse rise time. As a result, the coherence time is overestimated. h, Tunnel coupling $t_z/\hbar$ as a function of $V_{\text{cb}}$, demonstrating good control over the tunnel barrier between the two quantum dots. i, $T_2^\text{e}$ as a function of $\epsilon$. The dark red solid line is a fit to equation (2). We find $\delta g_{\text{spin}} = 7.59 \pm 0.49 \text{ meV}$, in line with comparable experiments, and $\delta E_{\text{pin}} = 1.78 \pm 0.01 \text{ meV}$, smaller by a factor of two than in a comparable natural Si qubit. The bright red (violet) dashed line represents the noise on $J$ ($\Delta E_2$). For low detuning, clearly detuning charge noise on $J$ dominates. At higher $\epsilon$, the sum of electric noise acting on $\Delta g$ and magnetic noise acting on $B$ limit coherence. Error bars correspond to one standard deviation.

magnetic noise contribution by isotopic purification could further improve qubit dephasing and quality\cite{11,19}. We now focus on extending the coherence of the qubit by applying refocusing pulses similar to those developed in nuclear magnetic resonance (NMR) experiments. We investigate the high $\epsilon$ region, where charge noise on detuning is lowest. Exchange pulses at $\epsilon = 0.64 \text{ meV}$ are adopted as refocusing pulses. We note, however, that to obtain a perfect correcting pulse, it would be necessary to implement a more complex pulse scheme\cite{8}. We choose convenient $\tau_z$ values ($\tau_z = (2n + \frac{1}{2})\tau_{\text{echo}}$, with $n$ being an integer and $\tau_{\text{echo}}$ being the time needed for a $\pi$ rotation along the $x$-axis) such that, if no decoherence has occurred, the system will always be found in the same state after $\tau_z$. The refocusing pulse is then calibrated to apply a $\pi$ pulse that brings the state on the same trajectory as before the refocusing pulse (Fig. 5a and Supplementary Fig. 16). The free evolution time after the last refocusing pulse $\tau_{\text{ref}}$ is varied in length from $\tau_z - \delta t$ to $\tau_z + \delta t$ (Fig. 5b,c, and we observe the amplitude of the resulting oscillations (Fig. 5e). Also, we increase the number of applied pulses.
from \( n_s = 2 \) to \( n_s = 512 \), thereby increasing the total free evolution time of the qubit and performing a Carr–Purcell–Meiboom–Gill echo. The decay is fitted to a Gaussian decay, and we extract the echo coherence time \( T_{E}^{\text{echo}} \) of 4.5 \( \mu \)s for \( n_s = 2 \) and \( T_{E}^{\text{echo}} \) of 158 \( \mu \)s for \( n_s = 512 \). Furthermore, we observe a power law dependence of \( T_{E}^{\text{echo}} \) as a function of the number of refocusing pulses and find \( T_{E}^{\text{echo}} \approx n_s^\beta \) with a constant \( \beta = 0.56 \), suggesting a limitation by low frequency 1/\( f \) noise\(^5\). We note that for \( n_s < 32 \), we extract \( \beta = 0.72 \), being a signature of quasi-static noise with spectral density \( \sim 1/ f \).

In conclusion, we have shown coherent two-axis control of a hole singlet-triplet qubit in Ge with an inhomogeneous dephasing time of 1 \( \mu \)s at 0.5 mT. We have taken advantage of an intrinsic property of HH states in Ge, namely their large and electrically tunable out-of-plane \( g \) factors. We achieved electrically driven \( \Delta g \) rotations of 150 MHz at fields of only 5 mT. Compared to \( \Delta g \)-driven singlet-triplet qubits in isotopically purified Si metal–oxide–semiconductor structures\(^{6,7,8,9,10}\), we find a \( g \)-factor difference that is three orders of magnitude larger. Moreover, we demonstrate an electrical tunability of the \( g \)-factor difference ranging from 50% to more than 200% over a gate range of 50 mV in different devices. The large \( g \)-factor differences were confirmed in two additional devices underlining the reproducibility of the Ge platform. While varying \( g \) factors might be an obstacle to scale-up, the fast developing field of automated tuning\(^{11} \) will be an asset in future experimental designs where \( g \) factors can be tailored in situ to the specific requirements. Echo sequences revealed a noise spectral density dominated largely by low frequency 1/\( f \) noise. The results and progress of singlet-triplet qubits, especially in the GaAs platform, will largely be applicable in Ge as well. Real-time Hamiltonian estimation\(^{12} \) can boost \( T_2^* \); a deeper understanding of the noise mechanisms might result in prolonging coherence even further\(^13 \); and feedback-controlled gate operation could push gate fidelity beyond the threshold for fault-tolerant computation\(^14 \).

In the future, a latched or shelved read-out could circumvent the decay of \( T_1 \) to the singlet during read-out, opening the exploration of the qubit’s behaviour at slightly higher magnetic fields where the \( \Delta g \) rotation frequencies could surpass the highest electron-dipole spin-resonance Rabi frequencies reported so far\(^3 \), without suffering from reduced dephasing times. Furthermore, by moving towards symmetric operation or resonant driving, the quality of exchange oscillations can be increased, since the qubit is operated at an optimal working point\(^{15,16,17} \); the operation of Ge qubits at very low fields can further improve their prospects in terms of scalability and high-fidelity fast read-out, as it will facilitate their integration with superconducting resonators and superconducting quantum interference devices\(^{18,19,20,21} \). The long coherence times combined with fast and simple operations at extremely low magnetic fields make this qubit an optimal candidate for integration into a large-scale quantum processor.

Online content

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SiGe spacers tend to display a small residual tensile strain by thermal chemical vapour deposition on reverse-graded buffers, the buffer and large out-of-plane factors differences. In the case of Ge QWs grown factors and low-energy plasma-enhanced chemical vapour deposition instead of thermal dislocation density of about 5×10^6 cm⁻². The graded virtual substrate typically presents a threading structure is a graded region approximately 100 nm thick in which the Ge content was increased linearly from pure Si to an idependent strain of the SiGe structure. The substrate temperature was reduced from 760 to 550°C with increasing Ge content. The buffer was completed with a 2 μm region at a constant composition of Si₀.₃Ge₀.₇. This part was concluded in about 30 min, with a growth rate of 5–10 nm s⁻¹ due to the efficient dissociation of the precursor gas molecules by the high-density plasma. The graded virtual substrate typically presents a threading dislocation density of about 7×10⁹ cm⁻² (ref. 44). The substrate temperature and plasma density was then reduced without interrupting the growth. The undoped Si₀.₃Ge₀.₇/Ge/Si₀.₃Ge₀.₇ QW stack was grown at 350°C and a growth rate of about 0.5 nm s⁻¹ to limit Si intermixing and interface diffusion. A 2 nm Si cap was deposited after a short (60 s) interruption to facilitate the formation of the native oxide (the interruption reduces Ge contamination in the Si cap from residual precursor gases in the growth chamber). Secondary ion mass spectroscopy analysis indicates that boron levels are below the detection limit of 10⁶ cm⁻³ to a depth of at least 200 nm.

**Device fabrication.** The samples were processed in the Institute of Science and Technology Austria Nanofabrication Facility. A 6×6 mm² chip is cut out from a four inch wafer and cleaned before further processing. The ohmic contacts are first patterned in a 100 keV electron beam lithography system; then a few nanometres of native oxide and the SiGe spacer is milled down by argon bombardment, and subsequently a layer of 60 nm Pt is deposited in situ at an angle of 5°, to obtain reproducible contacts. No additional intentional annealing is performed. A mesa of 90 nm is etched in a reactive ion etching step. The native SiOₓ is removed by a 10 s dip in buffered HF before the gate oxide is deposited. The oxide is a 20 nm atomic-layer-deposited aluminium oxide (Al₂O₃) grown at 300°C, which unintentionally anneals the ohmic contacts, resulting in a low-resistance contact to the carriers in the QW. The top gates are first patterned via electron beam lithography and then a Ti/Pt 3/27 nm layer is deposited in an electron beam evaporator. The thinnest gates are 30 nm wide and 30 nm apart. An additional thick gate metal layer is subsequently written and deposited and serves to overcome the mesa step and allow wire bonding of the sample without shorting the gates together. Quantum dots are formed by means of depletion gates (Fig. 1c). The lower gates (LB, C, R, RB) form a double quantum dot, and the upper gates are connected to a charge sensor (CS) dot. The separation gates in the middle are tuned to maximize the CS sensitivity to charge transitions in the double quantum dot. An inductance-capacitance (LC) circuit connected to a CS ohmic contact allows fast read-out through microwave reflectometry. LB and RB are further connected to fast gate lines, enabling fast control of the energy levels in the double quantum dot.

**Data availability**

All data included in this work are available from the Institute of Science and Technology Austria repository.

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**Author contributions**

D.J. fabricated the sample and performed the experiments and data analysis. D.J., A.H. and I.P. developed the fabrication recipe. D.J., A.H., O.S. and M. Borovkov performed precharacterizing measurements on equivalent samples. J.S.-M. and G.K. fabricated the two additional devices discussed in the Supplementary Information. J.K. performed the experiments on those additional devices. D.C. and A.B. designed the SiGe heterostructure. A.B. performed the growth, supervised by G.I.; D.C. performed the X-ray diffraction measurements and simulations. G.T. performed Hall effect measurements, supervised by D.C.; P.M.M. derived the theoretical model. M. Botifoll and J.A. performed the atomic resolution scanning transmission electron microscopy structural and electron energy-loss spectroscopy compositional related characterization and calculated the strain by using geometrical phase analysis. D.J., A.H., J.K., A.C., F.M., J.S.-M. and G.K. discussed the qubit data. D.J. and G.K. wrote the manuscript with input from all the authors. G.I. and G.K. initiated and supervised the project.

**Competing interests**

The authors declare no competing interests.

**Additional information**

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