Spins in semiconductor quantum dots constitute a promising platform for scalable quantum information processing. Coupling them strongly to the photonic modes of superconducting microwave resonators would enable fast non-demolition readout and long-range, on-chip connectivity, well beyond nearest-neighbour quantum interactions. Here we demonstrate strong coupling between a microwave photon in a superconducting resonator and a hole spin in a silicon-based double quantum dot issued from a foundry-compatible metal–oxide–semiconductor fabrication process. By leveraging the strong spin–orbit interaction intrinsically present in the valence band of silicon, we achieve a spin–photon coupling rate as high as 330 MHz, largely exceeding the combined spin–photon decoherence rate. This result, together with the recently demonstrated long coherence of hole spins in silicon, opens a new realistic pathway to the development of circuit quantum electrodynamics with spins in semiconductor quantum dots.

Cavity quantum electrodynamics (QED) deals with the interaction between the quantum degrees of freedom of an atom and the electromagnetic modes of a cavity1. The extension of this concept to superconducting quantum circuits has led to the development of circuit QED, opening new opportunities for the study of light–matter interaction and fostering the progress of solid-state quantum processors based on superconducting qubits2–4. In the same footsteps, a variety of alternative realizations have been explored using different types of quantum system as artificial atoms5. Hybrid systems made of quantum dots coupled to superconducting microwave resonators are a prominent example6–12. Of particular interest are silicon-based quantum dots owing to their ability to host long-coherence qubits encoded in a spin degree of freedom. Silicon-based spin qubits have made remarkable progress, reaching high fidelities in both one- and two-qubit gate operations, the latter being enabled by tunnelling-mediated exchange interaction between neighbouring qubits13. The co-integration with superconducting cavities acting as quantum buses would allow for long-range connectivity, largely facilitating the scalability of silicon spin qubits14–16.

As the spin does not directly couple to the cavity electric field, a spin–charge hybridization mechanism is needed to achieve coherent spin–photon interfaces. For electrons in Si/SiGe double quantum dots (DQDs), spin–photon coupling rates of a few tens of megahertz have been demonstrated with the help of a synthetic spin–orbit (SO) interaction created by nearby micromagnets16–19. The reported coupling rates are several times larger than the spin dephasing rate, thereby enabling coherent spin–photon coupling16,17 and cavity-mediated interaction between spins in distant DQDs18,19. Yet, to fully profit from circuit QED tools, including long-range, high-fidelity two-qubit operations and quantum non-demolition readout, a much stronger coupling strength is required, which necessitates more efficient coupling schemes.

In this work, we turn to a hole spin in a silicon nanowire metal–oxide–semiconductor (MOS) DQD to exploit the strong intrinsic SO interaction created by nearby micromagnets20–22, whose potential for circuit QED23–25 has remained unexplored. In our device geometry, the quasi-one-dimensional hole confinement enhances this SO interaction27 such that the SO length $\ell_{so}$, that is, the distance over which a
Spin rotates by $\pi$ due to SO interaction, is reduced to a few tens of nanometres, comparable with the DQD spatial extension $d$. The presence of such a strong SO interaction dramatically modifies the DQD energy levels, resulting in the formation of a flopping-mode SO qubit$^{23}$ whose energy is well separated from the other excitations of the DQD system. Here we demonstrate that this spin qubit strongly interacts with the quantized field of a high-impedance superconducting microcavity. We observe a spin–photon coupling rate as large as $\simeq 1\text{MHz}$, in line with recent predictions$^{25,26}$. A high-impedance (2.5 kΩ) microwave cavity is then patterned in the NbN film, along with a 50 Ω microwave feed line, ground planes and fanout lines (Fig. 1b). Besides being well suited for future large-scale integration, the used foundry-compatible MOS technology comes with large gate capacitances, resulting in tight electrostatic control and hence strong coupling to the electric-field component of the cavity mode.

Negative voltages applied to G1 and G2 accumulate holes in a DQD—so-called holes in the valence band of silicon$^{1}$—controls the spin–charge mixing in the DQD (as discussed later).

We probe the microwave response of this hybrid system in transmission at a temperature of 8 mK and at powers corresponding to less than one photon on average in the cavity ($n_{\text{avg}} = 0.1$), which is assumed to be in its ground state. We first characterize the bare-cavity response by sweeping the probe frequency $f_p = \omega_p/2\pi$ across the resonance frequency and keeping the charges in the DQD fixed. This way, we extract the bare-cavity resonance frequency $f_c = \omega_c/2\pi = 5.428$ GHz and cavity decay rate $\kappa/2\pi = 14$ MHz. To characterize the charge–photon coupling strength $g_c$, we monitor the transmission at frequency $f_c$ as $V_{G1}$ and $V_{G2}$ are varied (Fig. 1c). In the dark-blue region (Fig. 1c), the levels of the two dots are aligned so that a hole oscillates between the dots in response to the cavity electric field. We next probe the transmission as a function of $\phi$ to the nanowire axis.

Hole-spin-based circuit QED architecture

The spin circuit QED architecture designed here features a hole confined in a silicon DQD device interacting with a single microwave photon trapped in a superconducting cavity. The DQD is hosted in a natural silicon nanowire MOS transistor whose channel is controlled by four Ω-shape gates crossing the nanowire (Fig. 1a). The back-end-of-line fabrication of the silicon chip is interrupted to replace the first metallic interconnect layer with a 10-nm-thick niobium nitride (NbN) layer with large kinetic inductance$^{26,29}$ and magnetic-field resilience$^{29}$. A high-impedance cavity is then patterned in the NbN film, along with a 50 Ω microwave feed line, ground planes and fanout lines (Fig. 1b).
of $f_\text{c}$ and energy detuning $\epsilon$ between the two dots (Fig. 1d). This reveals a dispersive downshift in cavity resonance near $\epsilon = 0$, due to the electric-dipole interaction with the DQD hole charge.\(^{10,11}\) with energy $\hbar \omega_\text{c} = \sqrt{\epsilon^2 + 4t_\text{c}^2} > \hbar \omega_\text{n}$, where $h = h/2\pi$ is the reduced Planck constant (Fig. 1e,f). From the temperature dependence of this dispersive shift, we extract a charge–phonon coupling strength $g_\text{c}/2\pi = 513\text{ MHz}$ together with an interdot tunnel coupling $t_\text{c}/\hbar = 9.57\text{ GHz}$ (Supplementary Section IV).

**Strong hole spin–photon coupling**

An in-plane magnetic field $B$ lifts the spin degeneracy of the DQD charge states (Fig. 2a, inset). The two lowest spin-polarized states define an SO flopping-mode qubit.\(^{16,17}\) We probe the spin–photon interaction at $\epsilon = 0$, where the electric dipole of the hole in the DQD is the maximum. When the Zeeman spin splitting $\hbar \omega_\text{Z}$ matches the resonance frequency of the cavity ($f_\text{c} = f_\text{c} = \omega_\text{c}/2\pi$), spin–photon hybridization results in an avoided crossing that splits the cavity response into two branches separated by the vacuum Rabi-mode splitting.\(^{2,3}\) A representative measurement of this avoided crossing is shown in Fig. 2a, where the normalized transmission is plotted as a function of $f_\text{c}$ and $B = |B|$ at $\phi = 45^\circ$ with respect to the nanowire axis. The linecut at resonance shows two distinct dips separated by a vacuum Rabi-mode splitting $2g_\text{c}/2\pi = 184\text{ MHz}$, where $g_\text{c}$ is the spin–phonon coupling. The linewidth of these dips yields the decoherence rate of the hybridized spin–photon states $\gamma_\text{c}/2\pi = 7\text{ MHz}$, where $\gamma_\text{c}$ is the spin decoherence rate. From $\kappa/2\pi = 14\text{ MHz}$, we extract $\gamma_\text{c}/2\pi = 7\text{ MHz}$. The fact that $g_\text{c} \gg k, \gamma_\text{c}$ demonstrates a strong coupling between the hole spin and photon in the cavity.

**Spin–photon coupling versus magnetic-field orientation**

Varying the orientation of the in-plane magnetic field reveals pronounced anisotropy in the vacuum Rabi-mode splitting with a measured maximum (minimum) $g_\text{c}/2\pi$ of $330\text{ MHz (10 MHz)}$ at $\phi = 3^\circ (\phi = 79^\circ)$ (Fig. 3a,b). This large modulation in the spin–photon coupling strength results from the interplay between the Zeeman effect and SO interaction, leading to $g_\text{c} = g_\text{g} \left( |gB| \times (gB) \right)$, where $g_\text{g}$ is the average gyromagnetic $g$-matrix of the two dots and $B_\text{eff}$ is the effective SO field (note that $|gB|$ is related to the SO length (Supplementary Section IX)). The spin–photon coupling is, thus, expected to vanish when $B$ is parallel to $B_\text{eff}$ and is the maximum at $\phi = 90^\circ$, where the spin Larmor vector $gB$ is approximately perpendicular to $B_\text{eff}$. The $g$-matrix of holes embodies an anisotropic Zeeman splitting $E_\text{z} = \mu_\text{B} g_\text{B}$, in contrast to electrons ($E_\text{z} = 2 \mu_\text{B} B$) (ref. 19). In the present case, both dots show similar $g$-matrix anisotropies, with $E_\text{z} = 1.3 \mu_\text{B} B$ when $B$ is along the $x$ axis and $E_\text{z} = 2 \mu_\text{B} B$ when $B$ is along the $y$ axis. As shown in Fig. 3b, $g_\text{c}/2\pi$ gets almost entirely suppressed around $\phi = 75^\circ$. Since $g_\text{c}$ does not vanish completely, $B_\text{eff}$ must have a small out-of-plane component. Overall, however, we can conclude that the orientation of $B_\text{eff}$ is rather close to the $y$ axis. As discussed in Supplementary Section IX, the orientation of $B_\text{eff}$ is primarily determined by the device geometry. It is expected to be perpendicular to the interdot tunnelling direction ($\pm x$) as well as to the average electric-field direction in the interdot barrier region ($\pm z$ axis). Due to the anisotropy of the $g$-matrix, the magnetic-field orientation maximizing $g_\text{c}$ is not exactly orthogonal to $B_\text{eff}$.

The quality of the spin–photon interface can be further quantified by the ratio between the coupling strength and decoherence rate, that is, $2g_\text{c}/(\gamma_\text{c} + \kappa/2\pi)$ (Fig. 3c). This ratio reaches up to 27 for $\phi = 0^\circ$, which—along with a cooperativity $C = 4g_\text{c}/(\gamma_\text{c} + 1.600)$ (ref. 18)—highlights an extremely strong light–matter interaction. To evaluate the relative impact of spin dephasing and cavity decay rates, we can extract the angular dependence of $\gamma_\text{c}/2\pi$ (Supplementary Section VIII). We find values ranging from 2.5 to 17.0 MHz. Small (large) spin dephasing generally coincides with a small (large) spin–charge mixing. Over a large angular range centred around $\phi = 75^\circ$, $\gamma_\text{c}$ remains rather small and the quality of the spin–photon interface is mostly limited by the relatively large cavity decay rate $\kappa/2\pi = 14\text{ MHz}$.

A model for the spin–photon coupling of a hole DQD with SO interaction is presented in Supplementary Sections IX and X. With distinct anisotropic Zeeman response for the two quantum dots and spin-dependent tunnel couplings, this model captures $g_\text{c}$ at all the magnetic-field orientations (Fig. 3b). The effect of spin–charge mixing on the DQD energy diagram is illustrated in Fig. 3d,e for two different magnetic-field orientations. In particular, Fig. 3e highlights the large renormalization of energy levels in a DQD due to a strong SO interaction. Our model also catches the decrease in coherence by spin–charge mixing as the magnetic-field orientation approaches $\phi = 90^\circ$ (Fig. 3c) and Supplementary Section XII. In Fig. 3c, the black solid line is calculated assuming a charge qubit decoherence rate $\gamma_\text{c}/2\pi = 9.9\text{ MHz}$, due to the dominant detuning noise, on top of an isotropic bare spin decoherence rate $\gamma_\text{c}/2\pi = 3.4\text{ MHz}$, collecting other mechanisms such as nuclear-spin noise,\(^{15}\) electrically induced $g$-factor fluctuations\(^{16,17}\) or phonon-mediated spin relaxation.\(^{18}\)

**Spin–photon coupling in the single-dot limit**

The spin–photon couplings reported so far benefit from the large electric-dipole moment of the DQD at $\epsilon = 0$. In the following, we explore the interaction between the spin and cavity when the hole is localized in a single quantum dot. As previously reported for electrons,\(^{15}\) this interaction quickly vanishes for increasing detuning since charge localization in a single dot quenches the electric-dipole moment. This is evidenced by the detuning dependence of the charge–cavity coupling $g_{\text{c}}^{\text{eff}} = 2g_\text{c}t_\text{c}/\sqrt{4t_\text{c}^2 + \epsilon^2}$. Figure 4a shows the transmission when the microwave cavity is resonant with the spin splitting of the hole confined in the right dot ($\epsilon = -t_\text{c}$). Despite the considerable reduction in the hole-dipole moment, two dips are still visible in the transmission, implying $\gamma_\text{c} < g_\text{c}$. Their separation reveals spin–photon coupling $g_\text{c}/2\pi = 1\text{ MHz}$. Since $g_\text{c} < k$ (the so-called bad-cavity limit), a clear vacuum
Fig. 3 | Spin–photon coupling versus magnetic-field orientation. a, Normalized transmission as a function of probe frequency $f_p$ for various magnetic-field orientations $\phi$. All the curves are measured at resonance (that is, for $f_e = f_p$) and all of them show a clear vacuum Rabi-mode splitting. The curves are vertically offset and centred around the cavity resonance frequency $f_c$ for clarity. The dashed lines are fits to a superposition of two Lorentzians. For each vacuum Rabi-mode cut shown here, a full map of the avoided crossing (similar to Fig. 2) is shown in Supplementary Fig. 8. b–e, Angular dependence of spin–photon coupling $g$. (b) and $2g_c/(\gamma + \kappa/2)$ (c). The experimental data are in excellent agreement with the theory (solid line). The grey-shaded area outlines the magnetic-field orientations where the spin–photon resonance is achieved for magnetic fields larger than 1 T, which are inaccessible in our experimental setup.

Rabi-mode splitting cannot be resolved. To further support the existence of a single-dot spin–photon interaction, we measure $g_c$ as a function of $\epsilon$ (Fig. 4b). We find that $g_c$ drops by more than two orders of magnitude when increasing $|\epsilon|$, but tends to saturate once the hole is fully localized in the right dot. This limit presents potential interest since the use of single dots would simplify the device architecture, reduce the number of control parameters, allow for longer hole-spin coherence, and enable alternative and possibly more efficient spin–photon architectures. In addition, significant progress can be expected from the implementation of spin–photon coupling schemes relying on operational sweet spots, where decoherence is reduced and efficient electrical control is preserved.

Impact of strong SO interaction
The large spin–photon cooperativity observed in our experiment can be seen as a combined effect of an efficient charge–photon coupling, favoured by the MOS device layout, as well as the intrinsic SO interaction of holes, much stronger than the synthetic SO interaction of electrons in silicon. We note that the maximum $g_c$ is generally achieved when $2t_e = \hbar\omega_e = E_g$. Under this condition, the excited states $\{\uparrow\downarrow\}$ and $\{\uparrow\uparrow\}$ completely mix and $g_c$ approaches $g_s$, regardless of the SO interaction strength. In the case of electrons in silicon, however, the weak SO interaction ($\ell_{so} \gg d$) cannot keep these two states sufficiently apart from each other to prevent unwanted excitations to the second excited state $\{\downarrow\uparrow\}$ (ref. 8). This makes the limit of strong spin–charge mixing impractical, thereby preventing the achievement of maximum spin–photon coupling. The hole system studied here does not suffer from this limitation. The level repulsion induced by the stronger SO interaction ($\ell_{so} = d$) keeps the $\{\uparrow\downarrow\}$ and $\{\uparrow\uparrow\}$ states apart (Fig. 3e). Noticeably, the zero-detuning excitation energy of the flipping-mode qubit gets significantly reduced compared with Zeeman splitting in the absence of SO interaction. This implies that the spin–photon resonance condition occurs for $2t_e > \hbar\omega_e$. In this regime, the flopping-mode qubit is well isolated from the higher-energy levels and strongly coupled to the cavity mode. Moreover, operating the DQD at larger tunnel coupling ($2t_e = 4\hbar\omega_e$, Fig. 3e) reduces the impact of charge noise on the detuning energy, thereby enabling large cooperativity.

Conclusions
Looking further ahead, we foresee ample room to improve the spin–photon interface. On an engineering level, largely reduced resonator losses with $\kappa/2\pi < 1$ MHz should be readily feasible and, further improvements harnessing the advanced MOS fabrication platform to integrate a multilayer superconducting back-end-of-line would allow for a well-controlled microwave environment. We would also like to emphasize that in our hole system, the spin–photon coupling strength is ruled by the geometry and electrostatic design of the DQD (Supplementary Section IX) and controlled by the amplitude and direction of the externally applied magnetic field. We expect this should limit the impact of device-to-device variability and facilitate the development of large-scale quantum networks.

Our work promotes holes in Si MOS devices as a powerful playground for the development of spin circuit QED. As opposed to electrons, holes benefit from an intrinsically strong and versatile SO interaction ($\ell_{so} = d$), which is available without the need of micromagnets. This not only simplifies the device architecture but also allows to engineer an SO two-level system, which remains well isolated from other energy levels of the DQD. Our demonstration of such a hole-flopping-mode spin qubit, with an unprecedented spin–photon cooperativity of 1,600, opens the door to a wide range of circuit QED implementations of fundamental and practical interest (for example, two-qubit operations between distant hole spins with fidelities as high as 90% already seem within reach).
Fig. 4 | Spin–photon coupling in the single-dot limit. a, Normalized transmission as a function of $f_s$ and $B$ at $\varepsilon/h = -116$ GHz and $\phi = 11.25^\circ$ (left). Frequency linecut at the position of the arrows in the left panel (right). The spin transition is clearly visible and a double-dip structure split by ~2 MHz is observed at resonance. b, $g_s/2\pi$ as a function of energy detuning $\varepsilon$. The DQD flopping mode model (black solid line) effectively captures the large spin–photon coupling around $\varepsilon = 0$, but underestimates $g_s$ by more than one order of magnitude at large $\varepsilon$, where the hole is confined in a single quantum dot (SQD) on the left ($\varepsilon > 0$) or on the right ($\varepsilon < 0$). The measured $g_s$ is very well reproduced by adding an asymptotic single-dot spin–photon coupling $g_s^{(1)}(\varepsilon)$ in each quantum dot. We find $g_s^{(1)}(\varepsilon) / 2\pi \approx 1.16$ MHz and $g_s^{(2)}(\varepsilon) / 2\pi \approx 0.66$ MHz, as shown by the red solid line. The grey-shaded area outlines the energy detuning range where the spin–photon resonance is achieved for magnetic fields larger than 1 T, which are inaccessible in our experimental setup. The error bars represent the standard deviation from the fitting.

Online content
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Author contributions
C.X.Y. fabricated the NbN circuitry with help from S.Z. C.X.Y. and S.Z. performed the measurements. S.Z. analysed the data with inputs from C.X.Y., J.C.A.-U., É.D. and R.M. J.C.A.-U. developed the theoretical model with help from V.P.M., M.F. and Y.-M.N. S.Z., R.M., J.C.A.-U., S.D.F. and Y.-M.N. co-wrote the manuscript with inputs from all the authors. N.R., H.N, T.B., M.V. and B.B. were responsible for the front-end fabrication of the device. R.M. initiated the project.

Competing interests
M.V. is co-founder and CEO of squance.

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