

Document Version

Final published version

Citation (APA)

Jirovec, D., Reale, S., Fariña, P. C., Ventura-Meinersen, C., Nguyen, M. P., Zhang, X., Oosterhout, S. D., Scappucci, G., Veldhorst, M., Rimbach-Russ, M., Bosco, S., & Vandersypen, L. M. K. (2026). Many-body interferometry with semiconductor spins. *Science*, 392(6794), 183-187. <https://doi.org/10.1126/science.aed4177>

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QUANTUM SIMULATION

Many-body interferometry with semiconductor spins

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Quantum simulators enable studies of many-body phenomena, which are intractable with classical hardware. The manipulation of electronic spin states in devices based on semiconductor quantum dots promises precise electrical control and scalability advantages, but accessing many-body phenomena has so far been restricted by challenges in nanofabrication and simultaneous control of multiple interactions. In this study, we performed spectroscopy of up to eight interacting spins using a 2-x-4 array of gate-defined germanium quantum dots. The spectroscopy protocol is based on Ramsey interferometry and adiabatic mapping of many-body eigenstates to single-spin eigenstates, enabling complete energy spectrum reconstruction. As the interaction strength exceeds magnetic disorder, we observed signatures of the crossover from localization to a chaotic phase marking a step toward the observation of many-body phenomena in quantum dot systems.

Interacting many-body quantum systems are at the heart of some of the most notable phenomena in physics, ranging from high-temperature superconductors (1) to spin liquids (2). However, as system size grows, the systems can become exponentially hard to simulate using classical computations. Analog quantum simulation platforms, such as cold atoms in quantum gas microscopes (3), atomic tweezer arrays (4), superconducting circuits (5, 6), and ion trap processors (7, 8), have proven invaluable tools for studying such systems by mimicking many-body Hamiltonians in controlled laboratory settings, providing a clear path to insights into synthetic quantum matter. Compared with these platforms, arrays of gate-defined semiconductor quantum dots offer several advantages (9). Unlike most other platforms, they combine single-site addressability with native implementation of the Fermi-Hubbard model, leading to the observation of many unusual phenomena (10–14). Furthermore, at half-filling, the Fermi-Hubbard Hamiltonian exhibits particle-hole symmetry and reduces to an effective Heisenberg spin model, enabling quantum dots to function as high-performance spin qubits with demonstrated millisecond-scale coherence times and single-qubit gate fidelities exceeding 99% (15, 16). In the context of quantum simulations, the resulting Heisenberg interactions between localized spins have been used to explore interesting quantum phases at small scales (17–20).

Many exotic phases leave their fingerprint in the energy spectrum, through the emergence of energy gaps or changes in the level statistics (5, 21, 22). For some systems, the energy spectrum can be obtained on practical timescales with classical methods, therefore also serving as a convenient validation tool (8). Measuring the many-body energy spectrum is, however, not trivial, as traditional spectroscopic methods are often limited to low-energy excitations and are subject to selection rules that further narrow their application window.

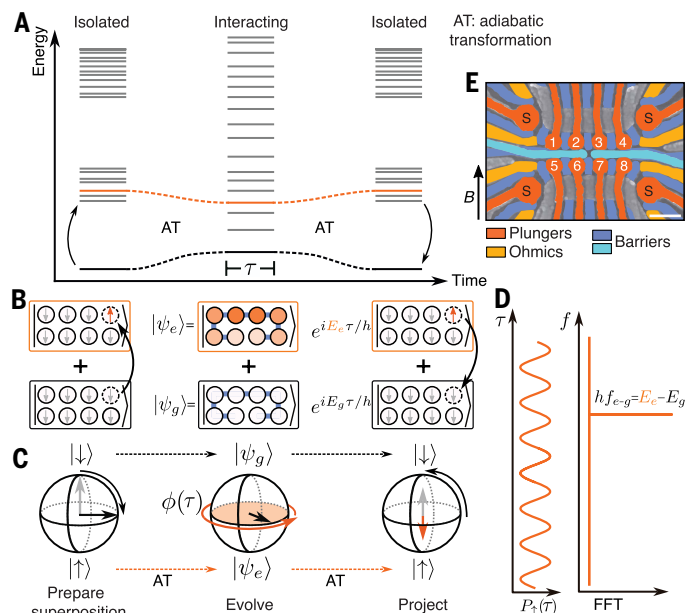


Fig. 1. Many-body Ramsey interferometry protocol. (A) Eigenenergies of the Heisenberg Hamiltonian with on-site disorder as a function of time. The eigenstates in the isolated and interacting regimes are connected by adiabatic transformations (ATs). Waiting for a time τ in the interacting regime allows for the superposition states to acquire a relative phase. (B) In the isolated regime, one spin can be initialized in a superposition state of $\uparrow$ and $\downarrow$. When interactions are turned on, $|\downarrow\rangle$ ($|\uparrow\rangle$) is adiabatically transformed to $|\psi_g\rangle$ ($|\psi_e\rangle$), and the initial spin excitation can delocalize over the array. The orange shades encode the probability of finding the spin excitation on a specific site. After a time τ and returning to the isolated basis, the states will have acquired a phase $E_g\tau/h$ and $E_e\tau/h$, respectively. (C) Bloch-sphere representation of the many-body Ramsey protocol. The relevant phase accumulation occurs in the interacting regime. At the end, the accumulated phase is projected back onto the readout basis. (D) The $\uparrow$ -probability of the qubit initialized in a superposition will oscillate with a frequency $hf = E_e - E_g$, allowing extraction of the many-body energy level spacings. (E) False-colored scanning electron microscope image of a nominally identical device to the one measured. Four charge sensors (S) in the corners of the device were used for radio frequency-reflectometry charge sensing. The plunger gates (orange) p_i and the barrier gates (blue and teal) b_{ij} confined eight unpaired charges in a 2-x-4 arrangement of quantum dots. An external magnetic field $B = 10$ mT was applied, as indicated. Scale bar, 200 nm.

In this work, we leveraged the strengths of spin qubits by utilizing a 2-x-4 array of hole spin qubits in Ge/SiGe to implement many-body Ramsey interferometry (23). This protocol extends the well-known Ramsey interferometry, used to characterize single-qubit energy-level splittings, to many-body states by adiabatically turning on and off interactions (24–26). With this approach (Fig. 1A), we initialized and disentangled superposition states and performed readout in the isolated, single-qubit (computational) basis. By adiabatically connecting the isolated spin states to interacting many-body states, the relevant phase accumulation of the Ramsey protocol occurred in the many-body sector (Fig. 1, A to C). The acquired phase was therefore a direct measure of the many-body energy level spacing (Fig. 1D).

We applied this method to the reconstruction of the energy spectrum of extended spin chains and showcased access to higher excited states. From the extracted energy levels, we inferred global system properties that shed light on the transition from local observables to information about extended many-body states.

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Experimental device and model

The device was fabricated on a Ge/SiGe heterostructure with a quantum well buried 50 nm beneath the surface hosting a two-dimensional hole gas. Once cooled in a dilution refrigerator, voltages applied to metallic electrodes (gates) defined the potential landscape in which we trapped eight unpaired hole spins in a 2-x-4 arrangement of quantum dots (Fig. 1E) (27).

With a single unpaired spin in each quantum dot, the system is well described by a Heisenberg Hamiltonian with on-site disorder and a spin-orbit-induced spin-flip term (28–31):

$$H = \frac{1}{2} \sum_i (E_{z,i} \sigma_{i,z} + \Delta_{\text{SO},i} \sigma_{i,x}) + \frac{1}{4} \sum_{(i,j)} J_{ij} \vec{\sigma}_i \cdot R_z(2\gamma_{ij}) \vec{\sigma}_j \quad (1)$$

where $E_{z,i} = g_i \mu_B B$ is the Zeeman energy, and $\Delta_{\text{SO},i}$ encodes the spin-orbit-induced avoided crossing. Here, $B = 10$ mT is the external in-plane magnetic field, μ_B is the Bohr magneton, g_i is the effective g -factor of the spin in dot i , $\vec{\sigma} = (\sigma_x, \sigma_y, \sigma_z)$ is the vector of Pauli matrices, and J_{ij} is the nearest-neighbor exchange interaction. The term $\Delta_{\text{SO},i}$ originates from the different orientations of the principal in-plane g -tensor components, which can be substantial for holes in planar Ge (32) as well as intrinsic spin-orbit contributions (29, 33). To capture the resulting anisotropy in the exchange interaction, $R_z(2\gamma_{ij})$ denotes a rotation about the z axis by an angle $2\gamma_{ij}$ (27, 30). We note that equivalent formulations of Eq. 1 utilize isotropic exchange but explicitly introduce quantization axis tilts through site dependent g -tensors (32, 34, 35).

We can develop intuition for the resulting system dynamics by using the magnon quasiparticle representation, where the exchange interactions J_{ij} constitute nearest-neighbor hopping amplitudes that enable coherent spin excitation transport across the lattice. The locally different g -factors introduce disorder through site-dependent Zeeman energies $\Delta E_{z,ij} = (g_i - g_j) \mu_B B = \Delta g_{ij} \mu_B B$, acting as random on-site potentials that compete with the coherent hopping dynamics (20). We found effective g -factors in the range of 0.3 to 0.47, such that $\Delta E_{z,ij} \ll E_{z,i}, E_{z,j}$ for all adjacent pairs ij .

In the regime where the exchange coupling strengths satisfy $J_{ij} \ll \Delta E_{z,ij} \ll E_{z,i}, E_{z,j}$, spin excitations remain localized to their initialization sites owing to the dominant single-particle Zeeman energies. The Hamiltonian eigenstates organize into well-separated spin manifolds characterized by the z component of total angular momentum, the total amount of $|\uparrow\rangle$ sites, with the n -th manifold spanning all states containing exactly $n|\uparrow\rangle$ spins (Fig. 1, A and B). Within this localized regime, energy level correlations are absent (5, 22), as the large Zeeman energy differences suppress intersite quantum tunneling, rendering each spin effectively isolated. As the exchange coupling J_{ij} increases relative to the Zeeman disorder $\Delta E_{z,ij}$, quantum tunneling enables spin excitation delocalization across the array, generating coherent spin-wave eigenmodes (36) and correlations among the energy levels (5). We also note that, in the strongly interacting regime where J_{ij} approaches the characteristic Zeeman energies $E_{z,i}$ and

$E_{z,j}$, the Δ_{SO} term induces hybridization between distinct spin manifolds, breaking the spin conservation of isotropic Heisenberg models. The strong spin-orbit coupling of holes in Ge, therefore, enables the study of unconventional and potentially nontrivial phenomena still largely unexplored by theory (31).

Interferometry of noninteracting spins and two-spin chains

As a starting point, we demonstrated interferometry of noninteracting spins. Standard spin-qubit Ramsey protocols rely on microwave-driven single-spin qubit rotations; however, the small Zeeman splitting at low magnetic fields, typical of our experiments, renders such resonant driving impractically slow compared with relevant decoherence timescales in our device ($T_2^* \approx 2$ μ s) (37, 38). For this reason, we initialized superpositions of spin states by leveraging the spin-orbit-induced avoided crossing (39, 40). In Fig. 2A, we plot the energy levels of a two-spin system as a function of time when pulsing the energy detuning between adjacent dots ϵ , as shown. The approximate eigenstates, symbolized by a tilde ($\sim$), are also labeled. In the following, we omit the tildes, implicitly referring to the approximate Hamiltonian eigenstates. At $t = 0$, we initialized $|S\rangle$ at $|\epsilon| \gg 0$, where both charges reside in the same dot [(2,0) or (0,2) charge configuration]). We then ramped ϵ to zero, passing semiadiabatically through the avoided crossing (dashed green circles) (41). This Landau-Zener passage initialized a superposition state $|\psi\rangle = \frac{1}{\sqrt{2}} (P_{LZ} |S\rangle + (1 - P_{LZ}) |T_{-}\rangle)$ with $P_{LZ} = e^{-2\pi\Delta_{\text{SO}}^2/\hbar v}$ being the Landau-Zener probability (40, 42, 43), and $\hbar$, the reduced Plank constant. Here, $v = \frac{dE}{d\tau_{\text{in}}}$ is the ramp speed, E is the energy difference of the states involved in the transition, and τ_{in} is the ramp-in time. After the

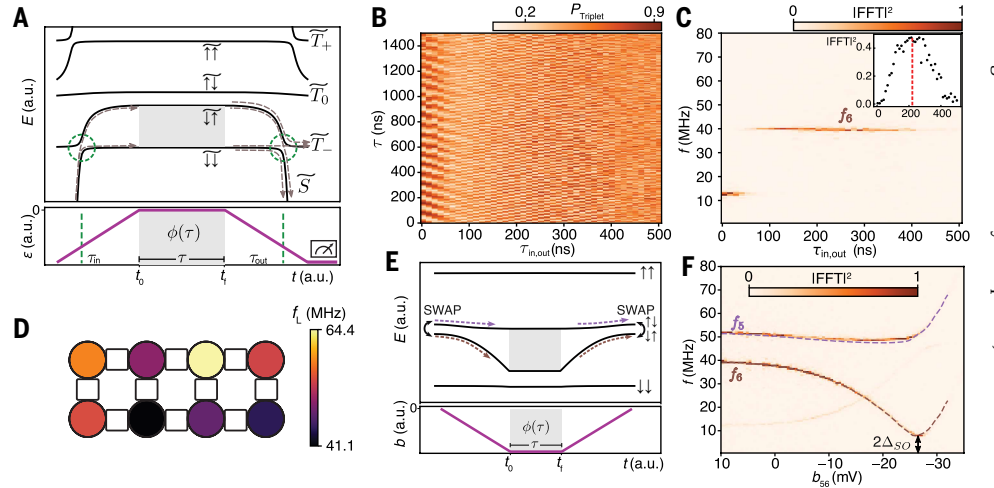


Fig. 2. Interferometry implementation. (A) (Top) Energy level diagram of a two-spin system as a function of the lab time. (Bottom) The corresponding pulse on the detuning ϵ . The gray dashed lines indicate the state trajectories resulting from the ramp in detuning ϵ , passing through the spin-orbit-induced avoided crossing (green dashed circles). After letting the system acquire relative phase ϕ for a time $\tau = t_f - t_0$, another ramp back toward the readout point closed the interference loop of the two trajectories. (B) Probability of finding the final state in a triplet as a function of ramp time τ_{ramp} and free evolution time τ for spin pair 5-6. (C) Squared norm of FFT of (B). The inset shows the FFT amplitude at f_6 as a function of $\tau_{\text{in,out}}$. The red dashed line marks the ramp time corresponding to $P_{LZ} = 0.5$. (D) Extracted Larmor frequencies f_L defining the magnetic disorder landscape. (E) Energy spectrum as a function of time and barrier voltage (as sketched in the bottom panel). After the initialization through the avoided crossing, we could choose to apply a local SWAP operation. If no SWAP (a SWAP) was applied, then we followed the lower (higher) branch. After adiabatically activating the exchange interaction, we let the system acquire a relative phase. The adiabatic turn-off was followed by no SWAP (a SWAP) in case we selected the lower (higher) branch. Subsequently, we closed the interference loop, like in (A). (F) Sum of FFTs of oscillations as a function of the interqubit barrier b_{56} using $\tau_{\text{ramp}} = 200$ ns. When turning interactions on and off without SWAPs applied, we observed the lower frequency. The brown and purple dashed lines are the fitted model (Eq. 1) from which we extracted the exchange interaction strength at each point. The minimum of f_6 corresponds to $2\Delta_{\text{SO}}$, a.u., arbitrary units.

avoided crossing, the detuning-ramp transformed $|S\rangle$ into $|\downarrow\uparrow\rangle$ adiabatically (with respect to ΔE_Z), which effectively initialized the state $|\psi\rangle = |\downarrow\rangle \otimes \frac{1}{\sqrt{2}}(P_{LZ}|\uparrow\rangle + (1-P_{LZ})|\downarrow\rangle)$, i.e., where only one of the spins in the pair is in a superposition. We then let the system acquire a relative phase $\phi(\tau) = \frac{(E_{11}-E_{11})\tau}{\hbar}$. Another semiadiabatic passage through the

avoided crossing with a ramp time of $\tau_{\text{out}} = \tau_{\text{in}}$ rotated the state back into the readout basis, closing the interference loop. Lastly, we read out the resulting triplet probability [(27) and supplementary text]. By choosing $\tau_{\text{in,out}}$ such that $P_{LZ} = 0.5$, this sequence is analogous to a traditional Ramsey experiment, and the oscillations that we observed are known as Landau-Zener-Stückelberg-Majorana oscillations (40).

Figure 2B plots the probability of measuring a triplet state after an experiment, as described in Fig. 2A, as a function of $\tau_{\text{in,out}}$ and waiting time in $(1,1)\tau$ for the spin pair 5-6. Figure 2C plots the squared norm of the fast Fourier transform (FFT) of Fig. 2B, where we can discern two main frequency components. For short ramp times, we observed oscillations with a frequency set by the Zeeman energy difference between the two spins: $2\pi\hbar f_{5-6} = \Delta g\mu_B B$, known as $S-T_0$ oscillations. At longer ramp times, the low-frequency component disappears, indicating adiabaticity with respect to ΔE_Z . However, we saw a higher-frequency component appearing with a maximum FFT amplitude around 220 ns (inset, Fig. 2C), corresponding to the energy difference of the states $|\downarrow_5\downarrow_6\rangle \leftrightarrow |\downarrow_5\uparrow_6\rangle$. The maximum visibility corresponds to $P_{LZ} \approx 0.5$, and in the localized regime, the oscillation frequency f_6 corresponds to the Larmor frequency of spin 6, $f_6 = g_6\mu_B B / 2\pi\hbar$, in this case (27). For $\tau_{\text{in,out}} > 400$ ns, this frequency component also tends to fade away, which indicates adiabaticity, now with respect to ΔS_0 .

To extract the other Larmor frequency in the pair, corresponding to spin 5, we can perform a local SWAP before and after the relative phase accumulation (27). Figure 2D summarizes the extracted Larmor frequencies for all the spins in the array, which determines the on-site magnetic disorder landscape for the system Hamiltonian.

For the many-body interferometry protocol, we adiabatically turned on the exchange and reconstructed the two-qubit energy spectrum with the pulse scheme depicted in Fig. 2E. Similar to Fig. 2A, we schematically plotted the energy spectrum as a function of time, but now only after the initialization and before the readout part. In this case, only the barrier voltage b was varied, maintaining $\epsilon = 0$. Note that barrier voltages are virtual, meaning that b_{ij} only affects J_{ij} to good approximation [(44); see details in (27)]. If we did (did not) apply a SWAP prior to and after the phase accumulation, we followed the upper (lower) energy branch. Figure 2F displays the sum of the FFTs of two experiments, one with and one without SWAPs. Without (with) SWAPs, we recorded f_6 (f_5) as a function of b_{56} . For each barrier value, we recorded oscillations for up to 1 μ s and extracted the highest peak position from the FFT through a Gaussian fit. We then fit the extracted frequencies to the full Hamiltonian (Eq. 1) for two spins to obtain J_{56} for each b_{56} value. In our fitting protocol, we generally assumed that the g -factors are not affected by barrier gate voltages (27). The size of the avoided crossing can be directly measured from the minimum frequency of f_6 . As we did not have access to different magnetic field orientations, we were not able to fully determine the angle γ_{56} . However, with the constraint of a finite minimum f_6 , throughout this work, we took $\gamma_{ij} = \pi/2$. The dashed lines overlaid in Fig. 2F show f_5 and f_6 using the extracted J_{56} for each b_{56} . The additional, faintly visible low-frequency line corresponds to the energy difference between f_5 and f_6 (hence, to the $S-T_0$ oscillation frequency), and it can be attributed to an imperfect SWAP (27).

We repeated the experiments in Fig. 2 for most qubit pairs and extracted the calibrated initialization and readout ramp times, SWAP times, and exchange values as a function of gate voltage (supplementary text).

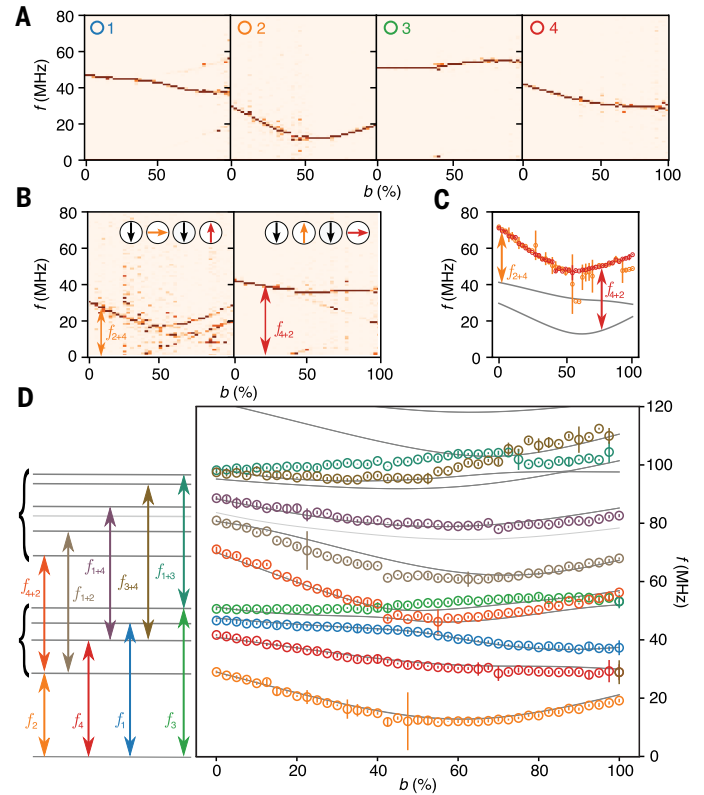


Fig. 3. Interferometry of extended spin chains. (A) FFTs of oscillations for spins in a four-spin chain as a function of global barrier gate voltage, in percent, where $b = 0\%$ corresponds to $J \ll \Delta E_Z$, and $b = 100\%$ corresponds to $J > \Delta E_Z$. The panels are ordered to reflect the position of the spin initialized in a superposition in the array. The four frequencies f_i allowed reconstruction of the first spin manifold. (B) FFTs of oscillations for the initial states are shown in the inset. On the left (right), we initialized spin 2 (4) in a superposition, whereas spin 4 (2) was initialized in $|\uparrow\uparrow\rangle$. (C) To reconstruct the energy levels of the second manifold, we added f_4 (f_2) measured in (B), represented by the solid gray lines, to the corresponding frequencies measured in (B). We saw that the data points lie almost perfectly on top of each other and returned the frequency of the state $|\downarrow_1\uparrow_2\downarrow_3\uparrow_4\rangle$ from two complementary measurements. Error bars represent the full width at half maximum of the fitted FFT peaks. (D) Reconstructed spectrum of the four-spin chain 1-2-3-4. When redundant frequencies were measured, we chose the one with the fewest artifacts. The gray lines show the spectrum obtained from the fitted model.

Interferometry of extended spin chains

The adiabatic turn-on of exchange can be simultaneously applied and uniformly implemented along extended spin chains by appropriately varying all the corresponding barrier gate voltages. To reconstruct the energy spectrum, we repeated the same interferometry experiment with all the different combinations of initial states. Mapping out the first spin manifold containing a single spin excitation, we prepared spin i in a superposition state $|\rightarrow\rangle = (|\uparrow\rangle + |\downarrow\rangle)/\sqrt{2}$ and measured f_i as a function of exchange. For the second spin manifold, we prepared spin j in $|\uparrow\rangle$ and spin i in a superposition. The total frequency was then the measured frequency f_i summed with the previously obtained f_j from the lower manifold. Similar arguments hold for higher spin manifolds. We adopted the convention that in $f_{i+j,k,l}$ the first index corresponds to the spin initialized in a superposition, whereas the other indexes correspond to the spins that are $|\uparrow\rangle$.

As an example for a chain of four spins, we plotted the extracted frequencies as a function of global barrier voltage of a chain formed by spins 1-2-3-4 for the first manifold in Fig. 3A. In Fig. 3B, we showcase measurements to reach the second manifold. In the left panel, we initialized the state $|\downarrow_1 \rightarrow \downarrow_3 \uparrow_4\rangle$, measuring f_{2+4} . In the right panel, we instead initialized $|\downarrow_1 \uparrow_2 \downarrow_3 \rightarrow_4\rangle$ and measured f_{4+2} . Figure 3C plots f_{2+4} (orange) and f_{4+2} (red) after adding the corresponding frequencies measured for the first spin manifold, represented by the solid gray lines. We clearly saw that the effective frequency that we extracted for the two initial states was the same, allowing us to determine the actual frequency corresponding to the state $|\downarrow_1 \uparrow_2 \downarrow_3 \uparrow_4\rangle$ in two ways. This redundancy can become a useful verification tool in the energy level reconstruction.

In Fig. 3D, we present the experimentally reconstructed energy spectra for both the first and second spin manifolds of this four-spin chain. The measured transition frequencies (circles) are color coded according to the double arrows in the schematic energy level diagram (left). The gray lines are the fitted model (Eq. 1) to the data. The values for the g -factors, Δ_{SO} for each pair, and the functional form of J_{ij} versus b_{ij} were close to those extracted from the data on pairs of spins (see supplementary text for further details). One spectral transition (light gray) remained experimentally unresolved owing to insufficient SWAP gate fidelity.

Inferring global system properties

Adding all the spins in the array, we performed comprehensive interferometry of an eight-spin chain (dots 6-5-1-2-3-4-8-7) as a function of global exchange coupling, reconstructing the energies of the first spin manifold (Fig. 4A). As for the shorter spin chains, we fitted Eq. 1 to the spectrum in Fig. 4B, once again finding good agreement after only small modifications to the fit parameters extracted from pairwise characterization, demonstrating robust control of seven simultaneously activated exchange interactions.

At interaction strengths $J/\Delta E_z \approx 0, 1, 2$ (vertical dashed lines, Fig. 4B), we extended spectroscopic reconstruction to the second manifold (Fig. 4C) and implemented statistical analysis of the energy levels to reveal signatures of quantum chaos. Higher excitation manifolds can, in principle, be reconstructed as well, through sequential SWAP protocols or driving, at the cost of added complexity in the pulse sequence.

From the energy levels $\{E_i\}$, we can compute the energy level uniformity $r_i = \frac{\min\{s_i, s_{i+1}\}}{\max\{s_i, s_{i+1}\}}$, where $s_i = E_i - E_{i-1}$ (5) (Fig. 4D). The level uniformity can capture energy level repulsion, characteristic of quantum chaotic systems, as the probability distribution $P(r_i \rightarrow 0) \rightarrow 0(2I)$. Such a correlation of energy level spacings is absent in localized systems, leading to randomly distributed energy levels. Figure 4E displays $P(r_i)$ as we increased the interaction strength. Indeed, we observed a reduction of counts for $P(r_i \rightarrow 0)$, signaling increased correlations in the energy levels. The gray dashed (black dashed

dotted) lines show the ideal distributions for localized (chaotic) systems, obtainable in large ensembles, as further detailed in the supplementary text.

Additionally, from $\{E_i\}$, we computed the spectral form factor (SFF), defined as $\text{SFF}(t) \equiv \frac{1}{N^2} \sum_{i,j} e^{i(E_i - E_j)t/\hbar}$ (45), where N is the system size.

This quantity can capture time correlations in the energy levels beyond nearest neighbor, making it a potentially more sensitive probe of quantum chaotic systems. Similar to r , the SFF displayed universal signatures that are different between localized and chaotic systems, described in more detail in the supplementary text. At relatively short times, the SFF develops a dip with a minimum value $\approx \frac{1}{N}$ ($\approx \frac{1}{N^{3/2}}$) for localized (chaotic) systems. Subsequently, energy level correlations lead to a linear in t ramp, up to the late-time limit $\frac{1}{N}$. Because the late time limit corresponds to the dip in the localized case, no linear ramp would develop, as correlations are absent (46). Notably, the SFF requires sufficient statistics to display all the universal features, which are otherwise masked by noise. Because we only observed small changes of the g -factors with gate voltage, an average over disorder was unfeasible in this device (27, 47). However, in (27), we showed that the early time behavior was robust even without averaging. Indeed, the SFF displayed in Fig. 4F systematically developed a deeper early

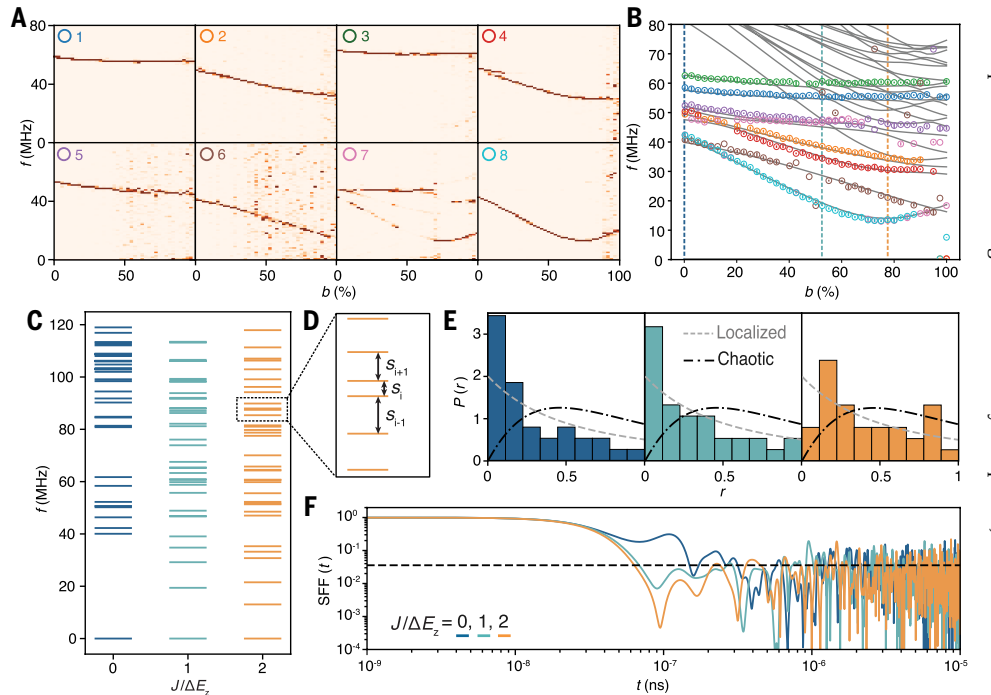


Fig. 4. Inferring global system parameters. (A) FFTs of oscillations for a chain of eight spins. The panels are ordered to reflect the position of the measured spin in the array. The barrier gates were all swept together from low exchange $J \ll \Delta E_z$ ($b = 0\%$) to high exchange $J > \Delta E_z$ ($b = 100\%$). Except for spin 7, we could always clearly identify the main frequency. (B) Extracted frequencies from (A), with overlaid fitted model in gray. Error bars represent the full width at half maximum of the fitted FFT peaks. At the colored dashed lines, we reconstructed higher spin manifolds. (C) Energy levels of the first and second spin manifold extracted for three different values of $\frac{J}{\Delta E_z}$ highlighted in (B). We observed clearly separated manifolds for $\frac{J}{\Delta E_z} \approx 0$. For $\frac{J}{\Delta E_z} \approx 2$, the energy levels tended to repel each other, signaling correlations and, consequently, a tendency for chaotic behavior. (D) Visualization of the definition of s_i . (E) Energy level uniformity distribution for the three exchange configurations calculated from the energy levels extracted in (C). The gray dashed (black dashed dotted) line represents the expected theoretical distribution for localized (chaotic) systems. The data displays a trend to more chaotic behavior as interactions increase, signaled by $P(r \rightarrow 0) \rightarrow 0$ for $\frac{J}{\Delta E_z} \approx 2$. (F) Spectral form factor calculated for the energy levels extracted in (C). We observed a systematic trend to reach a deeper early-time minimum smaller than the late-time plateau (black dashed line at a value of $1/N$) as we increased the interaction strength. For longer times, oscillatory behavior masked the universal features of the SFF, as we did not average over different disorder realizations.

time dip around $t \approx 10^{-7}$. The systematic trends of both $P(r_i)$ and the SFF as a function of interaction strength suggest a transition from localization to quantum chaotic behavior in the eight-spin chain.

Conclusions

We have demonstrated many-body interferometry of spin chains with up to eight spins. To perform the interferometry, we leveraged a spin-orbit-induced avoided crossing, which allowed us to initialize and disentangle superposition states using only baseband pulses. Furthermore, to access additional initial states, we implemented local SWAPs. By initializing different states and adiabatically turning on exchange interactions, we probed the energy levels in two-, four-, and eight-spin chains, supporting our findings with extensive simulations.

After mapping out the energies of higher spin manifolds, we computed the energy level uniformity and the spectral form factor as a function of interaction strength, finding trends consistent with a crossover from localized to quantum chaotic behavior. Further work will be required to unveil more universal features in the spectral form factor as well as the energy level uniformity.

We note that, as system size increases, the energy spectra become exponentially more dense and more avoided crossings appear, and tailored transformation strategies need to be implemented to maintain adiabaticity for various initial states, an active area of research in the context of quantum annealing (48). The present method combined with our earlier studies of magnon propagation (20) will allow us to engineer multispin Hamiltonians and probe their dynamics to reveal signatures of many-body localization.

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ACKNOWLEDGMENTS

We thank T.-K. Hsiao and members of the Vandersypen, Veldhorst, Scappucci, Rimbach-Russ, and Bosco groups for stimulating discussions. We acknowledge technical support from O. Benningshof, J. D. Mensingh, T. Orton, R. Schouten, R. Vermeulen, R. Birnholtz, E. Van der Wiel, B. Otto, D. Brinkman, and S. de Snoo. Claude opus 4.1 and Chat GPT 4o were used to refine parts of the manuscript (27). The views, conclusions, and recommendations contained in this document are those of the authors and are not necessarily endorsed nor should be interpreted as representing the official policies, either expressed or implied, of the Army Research Office or the US Government or of The Netherlands Ministry of Defence. The US Government and The Netherlands Ministry of Defence are authorized to reproduce and distribute reprints for government purposes notwithstanding any copyright notation herein.

Funding: This work was financed by the Advanced Grant from the European Research Council (ERC) under the European Union’s Horizon 2020 research (882848) and the ERC Starting Grant QUIST (850641) as well as the NWO Spinoza prize awarded to L.M.K.V. by the Netherlands Organisation for Scientific Research (NWO). This work was partly supported by the Army Research Office under award number W911NF-23-1-0110. This research was sponsored in part by The Netherlands Ministry of Defence under award no. QuBits R23/009.

Author contributions: D.J., S.R., and L.M.K.V. conceived the experiment. D.J. and S.R. performed the experiment and analyzed the data with contributions from P.C.F. and X.Z. M.P.N. and C.V.M. developed the theory and produced the simulations. S.D.O. fabricated the device. G.S. developed and provided the Ge/SiGe heterostructures. D.J., S.R., and L.M.K.V. wrote the manuscript with input from all the authors. **Competing interests:** M.V., G.S., and L.M.K.V. are founding advisers of Groove Quantum BV and declare equity interests. The remaining authors declare that they have no competing interests. **Data, code, and materials availability:** The data, analysis code, and simulation code are available in a Zenodo repository (49). Details on materials processing and device fabrications are outlined in (27). **License information:** Copyright © 2026 the authors, some rights reserved; exclusive licensee American Association for the Advancement of Science. No claim to original US government works. <https://www.sciencemag.org/about/science-licenses-journal-article-reuse>. This research was funded in whole or in part by European Union Horizon 2020 (882848), a cOAlition S organization, and by the European Research Council [ERC Starting Grant QUIST (850641)]. The author will make the Author Accepted Manuscript (AAM) version available under a CC BY public copyright license.

SUPPLEMENTARY MATERIALS

science.org/doi/10.1126/science.aed4177

Materials and Methods; Supplementary Text; Figs. S1 to S16; References (50–66)

Submitted 30 October 2025; accepted 6 February 2026

10.1126/science.aed4177



Many-body interferometry with semiconductor spins

D. Jirovec, S. Reale, P. Cova Fariña, C. Ventura-Meinersen, M. P. Nguyen, X. Zhang, S. D. Oosterhout, G. Scapucci, M. Veldhorst, M. Rimbach-Russ, S. Bosco, and L. M. K. Vandersypen

Science **392** (6794), . DOI: 10.1126/science.aed4177

Editor's summary

The use of quantum systems allows the simulation of complex many-body interacting systems otherwise intractable to classical computational methods. Whereas methods based on cold atoms, trapped ions, and superconducting circuits are now relatively large, solid-state-based systems offer the flexibility of individual qubit control and complementary metal-oxide semiconductor compatibility, but they have been limited in size. Jirovec *et al.* developed a scalable route to probing and controlling collective quantum states in semiconductor spin qubit arrays. By fabricating a 2×4 array of gate-defined quantum dots, they extracted spectroscopic information from the interacting array as the coupling between qubits was varied. These results demonstrate a key step toward realizing solid-state quantum simulators. — Ian S. Osborne

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