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Observation of spin–orbit coupling induced Weyl points in a two-electron double quantum dot

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Recent years have brought an explosion of activities in the research of topological aspects of condensed-matter systems. Topological phases of matter are accompanied by protected surface states or exotic degenerate excitations such as Majorana modes or Haldane’s localized spinons. Topologically protected degeneracies can, however, also appear in the bulk. An intriguing example is provided by Weyl semimetals, where topologically protected electronic band degeneracies and exotic surface states emerge even in the absence of interactions. Here we demonstrate experimentally and theoretically that Weyl degeneracies appear naturally in an interacting quantum dot system, for specific values of the external magnetic field. These magnetic Weyl points are robust against spin–orbit coupling unavoidably present in most quantum dot devices. Our transport experiments through an InAs double-dot device placed in magnetic field reveal the presence of a pair of Weyl points, exhibiting a robust ground-state degeneracy and a corresponding protected Kondo effect.
Mathematical tools borrowed from topology find more and more applications in contemporary condensed-matter physics. In Weyl semimetals, e.g., the electronic band structure exhibits isolated degeneracy points, where two bands touch; this effect is also accompanied by exotic surface states. In three-dimensional systems, these degeneracy points can be protected by topology—and classified by a suitably chosen Chern number: continuous perturbations may displace these Weyl points in momentum space, but cannot break their degeneracy. Weyl point-related degeneracies of electronic states in molecules termed point-related degeneracies of electronic states in molecules termed in momentum space, but cannot break their degeneracy. Weyl points are protected by topology and cannot break their degeneracy.

The simplest example of a Weyl point arises when a spin-1/2 electron is placed in a homogeneous magnetic field (see Fig. 1a–c). In this example, the parameter space is spanned by the magnetic-field vector $\mathbf{B} = (B_x, B_y, B_z)$ and the two energy eigenstates are degenerate at $\mathbf{B} = 0$. We can associate a non-zero topological charge to this degeneracy point: the ground-state Chern number $C(S) = 1$ is evaluated on an arbitrary closed surface $S$ surrounding the degeneracy point (see Methods for details). This non-zero Chern number promotes this $\mathbf{B} = 0$ degeneracy point to a Weyl point and underlines the robustness of its (Kramers) degeneracy against perturbations.

Let us now turn to the case of two coupled interacting spins and investigate the fate of Weyl points in the presence of—possibly strong—spin–orbit interaction (SOI). In the most general case, this system is described by the Hamiltonian $H = H_Z + H_{\text{int}}$, where $H_Z = \sum_{\alpha=\uparrow,\downarrow} \mu_B B_\alpha S_\alpha^z \sigma_\alpha^z + \sum_{\alpha=\uparrow,\downarrow} \mu_B B_\alpha (\sigma_\alpha^x S_\alpha^z + \sigma_\alpha^z S_\alpha^x)$ describes the Zeeman coupling and $H_{\text{int}} = \sum_{\alpha=\uparrow,\downarrow} (\sum_{\beta=\alpha} J_{\alpha\beta} S_\alpha^+ S_\beta^- + \text{h.c.})$ is just the exchange interaction. The SOI appears here through the anisotropic and dot dependent $g$-tensors, $g_{\alpha/R}$, and the anisotropic exchange coupling tensor, $J$.

In the absence of SOI (Fig. 1d–f), the $g$-tensors as well as the exchange coupling are just scalars, $g_{\alpha/R} \rightarrow g_{\alpha/R}$ and $J \rightarrow J$. The energy spectrum (Fig. 1e) is therefore isotropic as a function of the magnetic field. For an antiferromagnetic coupling, the ground state becomes degenerate at a sphere of radius $B = J/(\mu_B g)$, where a singlet to triplet transition occurs (Fig. 1f).

According to naive expectations, a small SOI should mix the singlet and triplet states close to the sphere of degeneracies, and thereby remove the degeneracy immediately (Fig. 1g–i). Simple models, however, indicate that this may not be the case. For example, quasi-two- and quasi-one-dimensional quantum dots are often described in terms of a Rashba field, $B_{\text{Rash}} \| \mathbf{p}$, with $\mathbf{p}$ a substrate- or gate-induced (usually) vertical electric field, and $\mathbf{p}$ the momentum of the carriers. In such models, singlet–triplet mixing is found to be absent if the external field is aligned with the effective Rashba fields, whereas the degeneracy is lifted in other external field directions.

Unfortunately, the somewhat simplistic Rashba-field model is not quite appropriate for three-dimensional or disordered quantum dots, similar to ours: it fails to account, e.g., for the strongly distorted, dot-dependent $g$-tensors measured earlier in InAs nanowire double quantum dots (DQDs), as well as in our device, which are rather aligned with each other not correlated with the geometry of the sample.

In this work, we consider two interacting spins confined in a spin–orbit-coupled DQD. Theoretically, we study a generic Hamiltonian with spin-dependent tunneling and distorted $g$-tensors, and use topological considerations to provide the conditions for the existence of ground-state degeneracy points in the three-dimensional magnetic-field parameter space. Experimentally, we demonstrate the presence of such magnetic Weyl points by doing transport spectroscopy in an InAs DQD as we explore the magnetic-field parameter space. Furthermore, we show that these degeneracies lead to a two-electron Kondo effect. Our results establish generic, robust, topologically protected degeneracy points, insensitive to microscopic details, in spin–orbit-coupled interacting spin systems.

Results

Magnetic Weyl points are topologically protected. We show now that the presence of degeneracy points is not a consequence of an oversimplified description or the simple dot geometries, but is rooted in topology. To show this, we consider first the most generic spin–orbit-coupled two-spin Hamiltonian $H$ above. Consider a sphere in the magnetic parameter space, centered at the origin, with a radius approaching infinity, and calculate the corresponding ground-state Chern number, $C_{\text{infinite}}$. Each spin follows just the external field, yielding a Chern number, $C_{\text{infinite}} = \text{sign } \det (g_{\uparrow}) + \text{sign } \det (g_{\downarrow}) \equiv C_{\text{infinite}}^+ + C_{\text{infinite}}^-$. As long as both $g$-tensor determinants are positive, we simply obtain $C_{\text{infinite}} = 2$. As, by definition, $C_{\text{infinite}}$ counts the total topological charge carried by the degeneracy points in the entire magnetic field space, the finiteness of $C_{\text{infinite}}$ signals the existence of ground-state degeneracies with nonzero topological charge, typically located at single points, which we call magnetic Weyl points. Time reversal constrains the locations of these points (see Methods): a degeneracy point at $B_0$ must have a partner at $-B_0$, carrying the same topological charge.

By these topological considerations, we expect that such Weyl points at $B_0$ carry the total topological charge $C_{\text{infinite}} = 2$. Using random spin Hamiltonians, we have numerically verified that this scenario of two magnetic Weyl points is generic and is indeed realized in cca. 98% of randomly generated two-spin Hamiltonians. In the remaining cca. 2% of cases, the number of Weyl points is six, but the sum of their topological charges remains two.

Experimental observation of magnetic Weyl points. To demonstrate experimentally the existence of magnetic Weyl points in a spin–orbit-coupled interacting two-spin system, we carried out low-temperature electric transport measurements through a serial
InAs nanowire DQD\(^{16,18-21}\) in the temperature range 60–300 mK. The setup is sketched in Fig. 2a (for sample fabrication and characterization see Methods and Supplementary Note 1). Alternative experimental techniques to explore these magnetic Weyl points are Landau–Zener\(^{22-24}\) or electrically driven spin resonance spectroscopy\(^{15}\), as well as Pauli-blockade spectroscopy\(^{21}\), as applied to various two–electron double-dot devices.

In the experiments, we focused on the (1,1) charge configuration of the device (see Fig. 2b), where the DQD contains two spatially separated and exchange-coupled spins. In this region, we expect that the ground state of the system is a singlet and the first excited state separated by \(\Delta E \approx J_0\) is a triplet (see Supplementary Note 2). The finite exchange splitting \(J_0 \approx 0.055\) meV is demonstrated by the bias-dependent differential conductance data presented in Fig. 2c. At the center of the (1,1) configuration, i.e., along the vertical dashed line, the conductance is suppressed at small biases, but increases once the bias is sufficiently high to induce inelastic co-tunneling processes populating the triplet states. The differential conductance \(G = dI/dV_{\text{bias}}\) (white curve) has therefore two finite-bias peaks (white lines) placed symmetrically at the first excited state of the DQD, at \(eV_{\text{bias}} = \pm \Delta E \approx \pm J_0\). The asymmetry \(G(V_{\text{bias}}) \neq G(-V_{\text{bias}})\) can be attributed to asymmetric coupling to the leads.

We now switch on the magnetic field to tune the relative energies of the ground and excited states, and explore by the cotunneling spectroscopy outlined above, how the energy gap \(\Delta E = \Delta E(B)\) between the ground and first excited states varies with the field (Fig. 1h)\(^{25-27}\). Two examples are shown in Fig. 2d, e, where we present the conductance \(G(B, V_{\text{bias}})\) for magnetic fields \(B = B_y \sin \theta \cos \phi, B_x \sin \theta \sin \phi, B_z \cos \phi\) oriented along two different directions (see reference frame in Fig. 2a).

In Fig. 2d, e, the magnetic-field dependence of the gap \(\Delta E(B)\) is traced by the large-conductance features close to zero bias, also indicated by solid lines. The observed behavior is markedly different in the two cases: Fig. 2d displays a behavior in line with the naive singlet–triplet mixing argument and the gap remains open for all values of \(B\). In Fig. 2e, however, the gap closes at around \(B_0 \approx 70\) mT, where a zero-bias conductance peak develops (white continuous line), suggesting that this magnetic-field vector corresponds to a magnetic Weyl point.
The scenario of the two magnetic Weyl points at opposite magnetic fields ±B_{0,m} fits perfectly our experimental observations. To demonstrate this, we display the conductance G(θ, V_{bias}) in Fig. 2f for a fixed magnetic-field length B = 75 mT and ϕ = 90°, while varying the polar angle θ over a range of 360° (see also Supplementary Note 3 for additional data). Our data indicate ground-state degeneracies at two opposite isolated points, θ = 130° and θ = 310°, but a finite gap otherwise. The solid lines in Fig. 2d–f, indicating the gap, are not only guides to the eye: they were computed from a two-site Hubbard model (see Methods), with parameters adjusted to yield a good overall match to experimental observations. Figure 2g visualizes the Berry curvature fields (see Methods) and the associated topological charges at the two Weyl points, as computed numerically from this two-site Hubbard model.

We support further the scenario of the two magnetic Weyl points by showing a more complete scan of the zero-bias magnetococonductance ΔG(B) = G(B) − G(B = 0) in Fig. 2h. The four panels of Fig. 2h correspond to four azimuthal angles, ϕ ∈ {−45°, 0°, 45°, 90°} of the magnetic field, as depicted in the sketch on the left side of panel h. Each panel of Fig. 2h displays the zero-bias magnetococonductance ΔG(θ, B) as the function of the polar angle θ and strength B of the magnetic field. The most prominent local maximum in the bottom right panel of Fig. 2h indicates that the two-electron double dot has a magnetic Weyl point close to that region, ϕ = 90°, θ = 130°, and B = 70 mT (also seen in Fig. 2e).

Two-electron Kondo effect at the Weyl points. In our device, the observed ground-state degeneracy is accompanied by an increased zero-bias conductance in the vicinity of the magnetic Weyl points (see white curve on Fig. 2e). This increased conductance is due to a two-electron Kondo effect29–31, as clearly revealed by the temperature and voltage dependence of our transport data in Fig. 3, complying with the Kondo behavior seen in other experiments30–32. The differential conductance at the Weyl point exhibits, in particular, a pronounced zero-bias Kondo peak with a height increasing upon decreasing temperature (see Fig. 3a). This increased low-temperature conductance appears to be characteristic of the whole charge (1,1) domain, as demonstrated in Fig. 3b, presenting the temperature dependence of the zero-bias conductance along the diagonal dashed line in Fig. 2b. In contrast, in the regions corresponding to (2,0) and (0,2) charge configurations, the ground state is unique; there the conductance shows thermal activation and is suppressed with decreasing temperatures.

Discussion
So far, we have argued that for two interacting spins, the appearance of ground-state degeneracies at a pair of time-reversed magnetic Weyl points is generic and robust due to topological protection. The two Weyl points can change their positions in the magnetic-field space as the two-spin Hamiltonian is modified, but they cannot disappear, i.e., the corresponding degeneracy cannot split. Of course, this protection carries over to the observed Kondo effect.

In our experiments, we have tested the robustness of the Weyl points in two different ways. (1) We have also investigated other double-dot charge states and verified that the signatures of two Weyl points are indeed present there too, as predicted, although the quantum dot parameters as well as the location of the Weyl points changed considerably. (2) We have modified the microscopic parameters within the same charge state by varying the gate potentials. The Weyl points were displaced but have never disappeared (for further details, see Supplementary Notes 4 and 5).

Interestingly, in our simulations, we can find cases where four additional Weyl points emerge. In such cases, four out of the six Weyl points have a topological charge +1, whereas two of them have charge −1, adding up to a total topological charge +2, in agreement with our sum rule. This is shown in Fig. 4a, where the red (blue) spheres represent Weyl points with Chern number +1 (−1). As discussed below, this case—never realized in the simple Rashba-field model—may be relevant, e.g., in Si double dots33,34.

The numerical simulation in Fig. 4a illustrates further topological protection: Weyl points can move around upon continuous deformation of the Hamiltonian and degeneracies of opposite topological charge can annihilate each other, but the total topological charge remains unchanged and assures the presence of degeneracies. A particular example of an annihilation process is presented in Fig. 4a, where arrows indicate the motion of degeneracy points, while their colors refer to their charge (see Supplementary Note 6 for a description of the deformation protocol).

The creation and annihilation of magnetic Weyl points should be observable in double-dot devices. Although using completely random parameters the probability of finding six Weyl points is only cca. 2% as illustrated of Fig. 4b, engineering the parameters allows for increasing ratio of six Weyl points. An analysis of our InAs double-dot setup using the experimentally determined g-tensors and the ratio of spin-conserving and spin-rotating tunnel amplitudes (see Supplementary Note 6) yields that the probability of creating additional Weyl points is still low, cca. 10–20%. However, in other experimental systems with lower spin-rotating tunneling amplitude, the likelihood of finding six Weyl points can increase significantly (50–100%) (see Supplementary Note 6). Silicon devices33,34 or bent carbon nanotubes could be promising candidates in this regard.

Remarkably, the argument applied to two coupled spins can be generalized to interacting multi-spin systems, such as magnetic...
trimers, atomic clusters, or multi-dot arrangements. In fact, the total topological charge for $N$ non-interacting spins with size 1/2 and isotropic $g$-tensors is $C_n = N$. This suggests that for a generic $N$-spin system there are $N$ magnetic-field values where Weyl points of topological charge $+1$ appear and the ground state is degenerate. For even values of $N$, these degeneracies must appear as time-reversed pairs, whereas for an odd number of spins a Weyl point must appear at $B = 0$, as also implied by Kramers’ theorem. These arguments can be readily extended to systems of spin $S$ impurities as well, where the total topological charge adds up to $C_n = 2S N$.

These general considerations have direct experimental implications. To demonstrate this, we generated an ensemble of random Hamiltonians for a spin–orbit-coupled three-spin quantum dot system (see Supplementary Note 6) and analyzed the statistics of the number of magnetic Weyl points (Fig. 4c). We dominantly observe three Weyl points of charge $+1$ ($\sim 75\%$), but with a $25\%$ likelihood, additional pairs of degeneracy points appear. From the trend observed in Fig. 4b, we anticipate that the more complex interaction space, the easier it is (1) to find more Weyl points than the number of spins and (2) to measure the controlled creation and annihilation of Weyl points.

We thus established that magnetic Weyl degeneracies are generic in interacting quantum dot devices in the presence of SOI. The precise location and structure of these degeneracies may depend on microscopic details, but their presence is assured by topology. Their robustness has important physical implications such as the corresponding topologically protected Kondo effect observed.

### Methods

**Sample fabrication and measurement details.** An array of Cr/Au (with 5/25 nm thickness) bottom gates (see Fig. 2a) with a width of 40 nm and a period of 100 nm was prepared by e-beam lithography and e-beam evaporation on a Si/SiO$_2$ substrate. Exfoliated hexagonal boron-nitride (hBN) flakes with a thickness of 20 nm were positioned on top of the bottom gates by a transfer microscope to electrically isolate the bottom-gate electrodes from the nanowire. The 80 nm diameter InAs nanowire was placed on the hBN by a micromanipulator setup. The nanowire and the bottom gates were contacted by Ti/Au electrodes (10/80 nm), defined in a second e-beam lithography and e-beam evaporation step.

The sample was measured in Leiden Cryogenics CF-400 cryo-free dilution refrigerator, equipped with a two-dimensional vector magnet. To vary the magnetic field in three dimensions, the sample holder probe was rotated manually to four different orientations $\phi \in \{-45^\circ, 0^\circ, 45^\circ, 90^\circ\}$. After each rotation, the base temperature was different due to the different thermal contact between the probe and the cryostat. The differential conductance of the DQD was measured in a two-point geometry by lock-in technique at 237 Hz with 10 $\mu$V ac excitation with a home-built I/V converter. The conduction band was not fully depleted by the gates: charge-configuration labels in Fig. 2b therefore correspond to the number of electrons above closed shells in each quantum dot holding an unknown, large number of electron pairs.

### Berry curvature and Chern number

Consider the ground-state manifold $\psi_0(B)$ of a family of Hamiltonians $H(B)$ parameterized by the magnetic field $B$. Assuming that $\psi_0$ is differentiable in the vicinity of $B$, we define the Berry connection vector field $A = (A_x, A_y, A_z)$ as

$$A(B) = i \langle \psi_0(B) | \nabla_B | \psi_0(B) \rangle.$$  

(1)

The Berry curvature vector field $B = (B_x, B_y, B_z)$ is defined as the curl of the Berry connection,

$$B(B) = \nabla_B \times A(B).$$  

(2)

It is noteworthy that although the Berry connection $A$ is gauge dependent, the Berry curvature $B$ is not.

Consider now a closed surface $S$ in the magnetic-field space, such that the ground state is non-degenerate at any point of $S$. The (ground state) Chern number associated with this surface is then

$$C(S) = \frac{1}{4\pi} \oint_S d\Sigma \cdot B.$$  

(3)

For details, see Supplementary Note 7.

### Magnetic Weyl points form time-reversed pairs

If there is a magnetic Weyl point at $B_0$, then—by time reversal—there is also one at $-B_0$. This follows from the properties of time reversal, $r(\psi) = \psi^\dagger$, for any of the $\tau$-invariant operators, $x, y$ and $z$. (ii) $r(B)^2 = H(\tau - B)$. (iii) From (ii) it follows that if $H(B)|\psi\rangle = E|\psi\rangle$, then $H(\tau - B)|\psi\rangle = E^\dagger|\psi\rangle$. Thus, apart from an overall phase, $r(B)|\psi\rangle = |\psi(B)\rangle$. Thus a degeneracy at $B_0$ implies a degeneracy at $-B_0$, and at non-degenerate points $B(B) = -B(-B)$.

### Two-site Hubbard model of the double quantum dot

Theoretical results in Fig. 2d–f were produced by a spin–orbit-coupled two-site Hubbard model, with Hamiltonian $H = H_0 + H_0$. Here, the Hamiltonian in the absence of magnetic field is

$$H_0 = \frac{t}{2} \sum_{\langle \sigma \rangle} n_L (n_L - 1) + \frac{t_v}{2} \sum_{\langle \sigma \rangle} n_R (n_R - 1) + \varepsilon_L n_L + \varepsilon_R n_R + \sum_{\sigma \in \{ \uparrow, \downarrow \}} (\delta_{\sigma L} c_{\sigma L}^{\dagger} c_{\sigma R} + h.c.).$$  

(4)

with $U_{\sigma \pi}$ the strength of the Coulomb interaction on the left/right dot, $n_{\sigma \pi}$ the occupation numbers, $\varepsilon_{\sigma \pi}$ the gate-controlled on-site energies, and $t' = t_0 \delta_{\sigma L} - t_2 \sum_{\langle \sigma \rangle} \delta_{\sigma L} c_{\sigma L}^{\dagger} c_{\sigma R}$. In an external magnetic field, we also add the Zeeman terms $H_0 = \mu_B B \cdot (\sum_{\sigma \in \{ \uparrow, \downarrow \}} \sigma_{\sigma \pi})$. The spin-rotating interdot hopping as well as the nontrivial g tensors can be attributed to strong SOI in the InAs nanowires.

We have determined the values of the model parameters to provide a good overall agreement with the experimentally observed. For the methodology, see Supplementary Note 2. These parameters were then used to derive the theoretical results in Fig. 2g–l. The g tensors used were as follows:

$$\mathbf{g}_{\sigma \pi} = \begin{pmatrix} 2.136 & -1.089 & 0.443 \\ -1.089 & 11.696 & -5.315 \\ 0.443 & -5.315 & 6.617 \end{pmatrix}.$$  

(5)
\[ R_b = \begin{pmatrix} 8.739 & -1.703 & 5.835 \\ -1.703 & 8.637 & 1.532 \\ 5.835 & 1.532 & 14.713 \end{pmatrix} \]

Hoppings were set to \( t_0 = 0.0525 \text{meV}, t_e = -0.0151 \text{meV}, t = 0.0565 \text{meV}, t_L = -0.0697 \text{meV} \), and Coulomb energies to \( U_L = 1 \text{ meV}, U_b = 0.6 \text{ meV} \). The on-site energies corresponding to the center of the (1,1) hexagon of the charge stability diagram in Fig. 2b read \( e_L = -U_b/2 \) and \( e_b = -U_b/2 \).

**Data availability**
The data that support the plots within this paper and other findings of this study are available from the corresponding author upon reasonable request.

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**References**
1. Lutchyn, R. M., Sau, J. D. & Das Sarma, S. Majorana Fermions and a topological phase transition in semiconductor-superconductor heterostructures. *Phys. Rev. Lett.* 105, 077001 (2010).
2. Oreg, Y., Refael, G. & von Oppen, F. Holical liquids and majorana bound states in quantum wires. *Phys. Rev. Lett.* 105, 177002 (2010).
3. Haldane, F. D. M. Nonlinear field theory of large-spin Heisenberg antiferromagnets: semiclassically quantized solitons of the one-dimensional easy-axis Néel state. *Phys. Rev. Lett.* 50, 1153 (1983).
4. Turner, A. & Vishwanath, A. in *Topological Insulators* (eds Franz, M. & Molenkamp, L.) Ch. 11, 294–324 (Elsevier, Oxford, 2013).
5. Armitage, N. P., Mele, E. J. & Vishwanath, A. Weyl and Dirac semimetals in three-dimensional solids. *Rev. Mod. Phys.* 90, 015001 (2018).
6. Herring, C. Accidental degeneracy in the energy bands of crystals. *Phys. Rev.* 52, 365 (1937).
7. Murakami, S. Phase transition between the quantum spin Hall and insulator phases in 3D emergence of a topological gapless phase. *New J. Phys.* 9, 356 (2007).
8. Yankovic, G., Baebich, A., Asboth, A., Rougerie, S., Aassime, A., Jonckheere, A. & Abs vat, T. Magnetic control of a single electron spin in a semiconductor quantum dot. *Nature* 468, 490 (2010).
9. Armitage, N. P., Mele, E. J. & Vishwanath, A. Weyl and Dirac semimetals in three-dimensional solids. *Rev. Mod. Phys.* 90, 015001 (2018).
10. Weng, L. et al. Photonics in magnetized plasma. *Nat. Commun.* 7, 12435 (2016).
11. Schrock, M. D. et al. Imaging a topological transition in an artificial spin-1/2 system. *Phys. Rev. Lett.* 113, 056402 (2014).
12. Roushan, P. et al. Observation of topological transitions in interacting quantum circuits. *Nature* 515, 241 (2014).
13. Nowack, K. C., Koppens, F. H. L., Nazarov, Y. V. & Vandervorst, L. M. K. Coherent control of a single electron spin with electric fields. *Science* 318, 1430 (2007).
14. Nadj-Perge, S., Frolov, S. M., Bakkers, E. P. A. M. & Kouwenhoven, L. P. Spin-orbit qubit in a semiconductor nanowire. *Nature* 468, 1084 (2010).
15. Nadj-Perge, S. et al. Spectroscopy of spin-orbit quantum bits in indium antimonide nanowires. *Phys. Rev. Lett.* 108, 166801 (2012).
16. Schrock, M. D., Petersson, K. D., Jung, M. & Petta, J. R. Field tuning the g factor in InAs nanowire double quantum dots. *Phys. Rev. Lett.* 107, 176811 (2011).
17. Meier, A. & Waller, S. Single-electron transistor with a single InAs nanowire. *Phys. Rev. Lett.* 106, 216808 (2011).
18. Fauth, C., Fuhrer, A., Samuelson, L., Golovach, V. N. & Loss, D. Direct measurement of the spin-orbit interaction in a two-electron InAs nanowire quantum dot. *Phys. Rev. Lett.* 98, 266801 (2007).
19. Pfund, A., Sorubalko, I., Enslin, K. & Leturcq, R. Suppression of spin relaxation in an InAs nanowire double quantum dot. *Phys. Rev. Lett.* 99, 036801 (2007).
20. Fauth, C., Fuhrer, A., Björk, M. T. & Samuelson, L. Tunable double quantum dots in InAs nanowires defined by local gate electrodes. *Nano Lett.* 5, 1387 (2005).
21. Wang, J.-Y. et al. Anisotropic Pauli spin-blockade effect and spin-orbit interaction field in an InAs nanowire double quantum dot. *Nano Lett.* 18, 4741 (2018).
22. Reilly, D. J. et al. Suppression of spin and orbital effects by nuclear state preparation. *Science* 321, 817 (2008).
23. Petta, J. R., Lu, H. & Gossard, A. C. A coherent beam splitter for electric spin states. *Science* 327, 669 (2010).
24. Fogarty, M. A. et al. Integrated silicon qubit platform with single-spin addressability, exchange control and single-shot singlet-triplet readout. *Nat. Commun.* 9, 4370 (2018).
25. Jeong, H., Chang, A. M. & Melloch, M. R. The Kondo effect in an artificial quantum dot molecule. *Science* 293, 2221 (2001).

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**Author contributions**
Experiments were designed by S.C. Wires were developed by J.N. and hBN crystals by K.W. and T.T. Devices were developed and fabricated by G. Frank, I.L., B.F., G. Fulop, and S.C. Measurements were carried out and analyzed by Z.S., G. Frank, and S.C. Theoretical analysis was given by A.P., Z.S., G. Frank, and G.Z. prepared the manuscript.

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