Direct mapping of the formation of a persistent spin helix

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The spin-orbit interaction (SOI) in zincblende semiconductor quantum wells can be set to a symmetry point, in which spin decay is strongly suppressed for a helical spin mode. Signatures of such a persistent spin helix (PSH) have been probed using the transient spin-grating technique, but it has not yet been possible to observe the formation and the helical nature of a PSH. Here we directly map the diffusive evolution of a local spin excitation into a helical spin mode by a time-resolved and spatially resolved magneto-optical Kerr rotation technique. Depending on its in-plane direction, an external magnetic field interacts differently with the spin mode and either highlights its helical nature or destroys the SU(2) symmetry of the SOI and thus decreases the spin lifetime. All relevant SOI parameters are experimentally determined and confirmed with a numerical simulation of spin diffusion in the presence of SOI.

Conduction-band electrons in semiconductors experience SOI from intrinsic¹ and extrinsic sources, leading to spin dephasing, current-induced spin polarization and spin Hall effects². These physical mechanisms are of great fundamental and technological interest, recently also in the context of topological insulators³ and Majorana fermions⁴. Intrinsic SOI arises from an inversion asymmetry of the bulk crystal (Dresselhaus term) and of the grown layer structure (Rashba term). In a quantum well (QW), these two components can be tailored by means of the confinement potential⁵, and the Rashba SOI can be externally tuned by using gate electrodes⁶. In general, SOI leads to precession of electron spins. In the diffusive limit, in which the scattering length is much smaller than the spin–orbit (SO) length λ_{SO}, a random walk of the spins on the Bloch sphere will dephase a non-equilibrium spin polarization⁷.

Of special interest is the situation in a two-dimensional electron gas (2DEG) with balanced Rashba and Dresselhaus contributions⁸. There, the SOI attains SU(2) symmetry and the spin polarization of a helical mode is preserved. The reason for this conservation of the spin polarization is a unidirectional effective SO magnetic field B_{SO}, which depends linearly on the component of the electron momentum along a specific in-plane direction. This causes the precession angle of a moving electron to vary linearly with the distance travelled along that direction, irrespective of whether the electron path is ballistic or diffusive⁹,¹°. In such a situation, a local spin excitation is predicted to evolve into a helical spin mode termed a PSH (Fig. 1a). Transient spin-grating measurements⁶ showed that a spin excitation with a spatially modulated out-of-plane spin component decays with two characteristic lifetimes that correspond to two superposed spin modes of opposite helicity.

Here we directly measure the diffusive evolution of a local spin excitation into a PSH by time-resolved Kerr rotation microscopy (Fig. 1b). We employ a pump–probe approach, in which a circularly polarized pump pulse excites electrons into the conduction band of a (001)-grown GaAs/AlGaAs QW with their spins polarized along z || [001]. The out-of-plane spin polarization S_{z} is then measured by a probe pulse delayed by a time t, using the polar magneto-optical Kerr effect. The position of the incident pump beam is scanned to record the spatial spin distribution S_{x}(x, y) at time t. We define x along the [1T0] and y along the [110] direction. The SOI of the 2DEG is tuned close to the SU(2) symmetry point, |α| ≈ |β_1 − β_3|, by controlling the Rashba (α), the linear Dresselhaus (β_i) and the cubic Dresselhaus (β_3) SO coupling coefficient via asymmetric modulation doping on the two sides of the QW.

The experimental observation of the PSH is exemplified by three maps of S_{i}(x, y) recorded at different t (Fig. 1c). The first map at t = 10 ps still shows the local excitation of S_{z} > 0 centred at x = y = 0. Because of the initially rapid spin diffusion, the Gaussian shape of S_{z}(x, y) is already broader than the size of the focused pump-laser spot. Spins further diffuse in the (x, y) plane, but the second and the third map (recorded at t = 240 and 840 ps) in addition feature alternating stripes of S_{z}(x, y) > 0 and S_{z}(x, y) < 0 caused by spin precession about B_{SO}. To explain this unidirectional oscillation along the y-direction, B_{SO} must be more strongly correlated with k_{y} than with k_{x}, k_{x} and k_{y} being the components of the electron wave vector k. With our definition of α and β and from the symmetry of B_{SO} (see Supplementary Information), it follows that α and (β_3 − β_1) must have the same signs and that therefore the x-component of B_{SO} is much larger than the y-component. For opposite signs of α and (β_3 − β_1), the PSH would oscillate along the x-direction and the y-component of B_{SO} would be larger.

The formation of the PSH is best illustrated if spin dynamics is tracked in space and time. For that purpose, we position the pump pulse at x = 0 and scan it along the y-direction. A collection of such line scans S_{z}(y, t) recorded at various t is shown as a colour-scale plot in Fig. 2c. Starting with S || z, the excited spin distribution expands along ±y, and thereby S_{z} starts to oscillate with y. As we will discuss in the following, this oscillation is indeed the footprint of a helical spin mode.

A schematic of helical spin modes is shown in Fig. 2a. The spin polarization rotates by an angle which depends linearly on the position y, with the direction of rotation defining a helicity (ω^+ or ω−). Therefore, both S_{x} and S_{z} oscillate with y (Fig. 2a). The helicity of the emerging spin mode depends on the sign of the correlation ⟨B_{SO}, k_{y}⟩, that is, on whether B_{SO,x} is positive or negative for k_{y} > 0, and is given by the absolute sign of α + β_1 − β_3 (see Supplementary Information). The helicity cannot be directly determined from S_{z}(y), as for both modes S_{x} oscillates in the same way (Fig. 2a). This is especially true for experiments employing the

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transient spin-grating method, in which, as a further complication, spin waves with both helicities are excited simultaneously.

We uncover the helical nature of the measured spin mode by rotating in-plane spin components \( S_z \) out of plane with the help of an external magnetic field \( B \parallel x \) (Fig. 2b). The effect of \( B_y \) on \( S_y(x, t) \) is shown in Fig. 2d. \( S_y(x, t) \) still oscillates with \( y \)—indicating that the PSH is preserved—but the phase of the oscillation now shifts with \( t \) (dashed line). This shift can be understood from the spin precession about \( B_y \), as shown in Fig. 2b: for an \( \omega^+ \) mode and \( B_y < 0 \), the position \( y_0 \) of equal phase shifts towards \( -y \) with increasing \( t \). As we observe the same direction in the measurement (Fig. 2d), we conclude that the three spin components support an \( \omega^+ \) mode, which means that \( \alpha + \beta_1 - \beta_3 \) must be positive. With this, we have directly determined the sign of the cross-correlation \( B_{Bx,k_y} \).

Figure 2d also shows that for all positions \( y \), \( S_y(x, t) \) oscillates in \( t \) with the same frequency \( v \). Therefore, \( S_y(x, t) \) represents a collective precession of the \( \omega^+ \) mode about \( B_y \). According to \( v = |g| \mu_B B_y / h \), we determine the electron moment \( g \) to equal \(-0.17 \) (based on the QW thickness, we assume \( g < 0 \); \( \mu_B \) is the Bohr magneton, and \( h \) is Planck’s constant).

We now quantitatively analyse \( S_y(x, t) \) to extract information on the SOI and on spin diffusion. Fig. 3a shows \( S_y(x, t) \) for different \( t \) with \( B_y = 1 \) T. The experimental data (symbols) are fit to the product of a Gaussian, \( A \exp(-y^2/4w^2) \), and a cosine function, \( \cos(2\pi(y - y_0)/A_{30}) \) (solid lines). The Gaussian envelope describes the solution of the spin-diffusion equation, for which \( w^2 = D_s t \) (\( D_s \) is the spin-diffusion constant). Spin dephasing is included in the time dependence of the amplitude \( A(t) \).

A linear fit of \( w^2 \) to \( D_s t \) provides a direct measure of \( D_s \) (Fig. 3b), and we obtain \( D_s = (385 \pm 15) \, \text{cm}^2 \text{s}^{-1} \). Note that our experiment uses the spin as a label of the electrons and therefore tracks the spin and not the charge diffusion. The sensitivity of spin diffusion to electron–electron interactions\(^{14,15}\) explains the tenfold smaller value of \( D_s \) as compared to the charge-diffusion constant \( D_c = 4 \times 10^{-4} \, \text{cm}^2 \text{s}^{-1} \) (as calculated for the measured electron mobility \( \mu = 22 \, \text{m}^2 \text{V} \text{s}^{-1} \) and a sheet carrier density \( n_s = 5 \times 10^{13} \, \text{m}^{-2} \)).

Figure 3b shows the decay of \( A(t) \) with \( t \). The diffusive expansion of the excited spins into an area proportional to \( D_s t \) decreases \( S_y \) proportional to \( (D_s t)^{-1} \). Further spin decay is mainly induced by deviations from perfect SU(2) symmetry\(^{16}\) and can be described by an exponential decay proportional to \( \exp(-t/\tau_s) \), where \( \tau_s \) is the spin lifetime. \( A(t) \) is therefore fit to \( (t - t_0)^{-1} \exp(-t/\tau_s) \), yielding \( \tau_s = (1.1 \pm 0.1) \, \text{ns} \). This is about 30 times longer than the Dyakonov–Perel spin dephasing time calculated from the measured \( D_s \) and SOI strength, in agreement with a spin decay time \( \tau_s \approx 35 \, \text{ps} \) obtained for 20-μm-wide laser spots, where the spatial correlations of the PSH are averaged out (see Supplementary Information).

The measured \( \tau_s \) is limited by the two SU(2)-breaking contributions, namely, the cubic Dresselhaus SOI \( \beta_3 \) and the imbalanced SOI \( |\alpha| \neq |\beta_1 - \beta_3| \). It can be shown that (Supplementary Information)

\[
\tau_s^{-1} \approx 2D_s \frac{m^2}{\hbar^2} \left( 3\beta_3^2 + (\alpha - \beta_1 + \beta_3)^2 \right)
\]

where \( m = 6.1 \times 10^{-26} \, \text{kg} \) is the effective electron mass. From the measured \( \tau_s \) and \( D_s \), we find \( 3\beta_3^2 + (\alpha - \beta_1 + \beta_3)^2 = 1.7 \times 10^{-26} \, \text{eV}^2 \, \text{m}^2 \). This relation restricts \( \beta_3 \) to an upper limit of
0.77 × 10⁻¹³ eV m, which would be reached for \( \alpha = \beta_1 - \beta_2 \). From the measured \( \tau_c \) alone, it is not possible to differentiate cubic Dresselhaus contributions from imbalanced SOI. However, as we will discuss later, the latter can be separately determined from a small asymmetry in \( S_y(y) \) maps that appear if B is applied along y. This will allow us to quantitatively describe all SO coefficients.

The fitted PSH period \( \lambda_{SO} \) is shown in Fig. 3d. A decrease from \( \approx 10 \mu m \) at \( t = 200 ps \) to \( \approx 7.5 \mu m \) at \( 1.5 \) ns is well represented in several periods of \( S_y(y) \) and must be related to a continuous change of \( \lambda_{SO} \) with \( t \). From the relation \( \lambda_{SO} = \frac{\pi h^2 m^{-1}(\alpha + \beta_1 - \beta_2)}{2|\mu_B B_0|} \), we determine that \( |\alpha + \beta_1 - \beta_2| \) increases from 3.5 to 4.9 \times 10⁻¹³ eV m during this time. Starting with a broad positive peak in \( S_y(y) \) at \( t = 10 \) ps, the spin helix continuously adapts to the decreasing \( \lambda_{SO} \). The initially weaker SOI is most probably induced by the photoexcited charge carriers that recombine with time, and is possibly also affected by cooling of hot electrons. The SO coefficients are sensitive to the screening of the confinement potential and modifications of the Fermi energy: both an increase of \( \beta_2 \) with charge density and a reduction of \( \alpha \) and \( \beta_1 \) with screening could explain the observed decrease of \( \lambda_{SO} \) with \( t \). Supporting this interpretation, measurements at lower pump intensities yield an initially smaller \( \lambda_{SO} \) (Fig. 3d).

\( \lambda_{SO} \) is found to be independent of \( B_0 \) (Fig. 3d). Together with the insensitivity of \( A(t) \) to \( B_0 \), this demonstrates the decoupled influence of the Zeeman and the SO energy on the electron spins in the case where a magnetic field is applied along the unidirectional \( B_0 \).

To understand this decoupling better, it is instructive to plot the directional dependence of \( B_{SO} \) on \( k = k(\cos\theta, \sin\theta) \). For (001)-grown QWs, \( B_{SO} \) is in the \((x, y)\)-plane for all \( k \). It can be written as the sum of two terms, \( B_{SO}^{(1)} \) and \( B_{SO}^{(2)} \). The former is responsible for the PSH formation, whereas the latter leads to spin dephasing (see Supplementary Information). The x- and y-components of \( B_{SO}^{(1)} \) are proportional to \( k_y(\alpha + \beta_1 - \beta_2) \) and \( -k_x(\alpha - \beta_1 + \beta_2) \), respectively. In our case, \( |\alpha + \beta_1 - \beta_2| > |\alpha - \beta_1 + \beta_2| \), which means that \( B_{SO}^{(1)} \) breaks the SU(2) symmetry and leads to spin dephasing. In the following discussion, we omit the superscript and refer to \( B_{SO} \) when we write \( B_{SO} \).

As shown in Fig. 4b, \( B_0 < 0 \) supersedes with \( B_{SO} \) such that the total field, \( B_{tot} = |B + B_{SO}| \), is larger for \( k_y > 0 \) than for \( k_y < 0 \). Translating the momentum \( h k_y \) into a position \( y \) using \( y = h k_y t / m \), this is exactly what is seen in the measurement at \( B_0 = -1 T \) in Fig. 2d: the total precession frequency \( \nu \) of an electron with momentum \( h k_y \) is determined by the sum of the Zeeman splitting, \( g \mu_B B_z \), and the SO splitting, \( 2k_x(\alpha + \beta_1 - \beta_2) \). Electrons with \( k_y = 0 \) remain at \( y = 0 \) and precess with \( \nu = |g \mu_B B_0 / h| \). There is a path \( y_0(t) \) (dashed line in Fig. 2d) on which Zeeman and SO energies cancel each other and consequently the spins do not precess. This path is characterized by

\[
\frac{\partial y_0}{\partial t} = \frac{-h g \mu_B B_0}{2m(\alpha + \beta_1 - \beta_3)}
\]

In agreement with this equation, the fitted \( y_0 \) increases linearly with \( t \) and \( \partial y_0 / \partial t \) changes sign with \( B_0 \) (Fig. 3c). Inserting the measured \( \partial y_0 / \partial t \) and \( g \) into equation (2), we obtain \( \alpha + \beta_1 - \beta_3 = 4.8 \times 10^{-13} \) eV m, consistent with the value obtained from \( \lambda_{SO} \), but here the sign is directly determined by the sign of \( \partial y_0 / \partial t \).

We now investigate how close the SOI in our QW is tuned to the SU(2) symmetry point. For that purpose we apply B along y. Figure 4c shows that in this case \( B_{tot} \) is no longer unidirectional.
Figure 3 | Spin diffusion and SOI characterization. a. Normalized line scans $S_y(t)$ for various $t$. The experimental data (symbols) are fits to $A\exp(-y^2/4w^2) \times \cos(2\pi(y-y_0)/\lambda_{SO})$ (solid lines). The parabola follows $\pm 2 / \sqrt{D_t}$, b. Symbols: amplitude $A$ and squared width $w^2$ of the Gaussian envelope. Solid lines: fits to $(t-t_0)^{-1} \exp(-t/\tau_c)$ and $D_t$, respectively. Error bars: uncertainty of $w^2$ in a single fit. c. Shift $y_0$ of the spatial oscillation for $-1$, $0$ and $1$ T. d. Symbols: PSH period $\lambda_{SO}$ for $0$ and $1$ T. Solid line: smoothed data at $1$ T. ($\times$) and ($\pm$): fits to data at threefold and ninefold lower pump intensities.

Figure 4 | Dependence of the total magnetic field $B_{tot}$ on $k$. a. The $x$-component of $B_{10,x}^{(1)}$ is shown (arrows) as a function of $k$ for $\alpha > 0$, $\beta_1 - \beta_3 > 0$ and $g < 0$. b. If an external magnetic field $B$ is applied along $x$, $|B_{tot}|$ is different for $k_x > 0$ and $k_x < 0$. The arrows represent the size and direction of $B_{10,x}^{(1)} + B_x$. c. For $B \parallel y$, $|B_{tot}|$ is the same for $\pm k_y$. d. Away from the SU(2) symmetry point ($\alpha \neq \beta_1 - \beta_3$), the $y$-component of $B_{10}^{(1)}$ is non-zero, and the superposition with $B_y$ leads to a different $B_{tot}$ for $k_x > 0$ than for $k_x < 0$. This allows us to determine $\alpha - \beta_1 + \beta_3$.

Even though a map of $S_y(x, t)$ at $B_y = -1$ T (Fig. 5a) still indicates that a helical mode evolves for $t < 400$ ps, $S_y$ quickly decays at longer $t$. In addition, an oscillation with $t$ is seen at $(x, y) = (0, 0)$. This is related to the precession of $S_y$ about $B_y$. On the other hand, the superposition of $B_y$ and $B_{SO,y}$ (Fig. 4d) leads to an asymmetry for opposite signs of $k_x$, which can be observed in the maps of $S_y(x, t)$ (Fig. 5b). We find a small tilt of the positions $x_0(t)$ of constant spin precession phase. In analogy to equation (2), the tilt is given by

$$\frac{\partial x_0}{\partial t} = \frac{\hbar g \mu_B B_y}{2m(\alpha - \beta_1 + \beta_3)}$$

This $\partial x_0/\partial t$ is a measure of the detuning from balanced SOI. We obtain $\partial x_0/\partial t \approx -280 \mu$m ps$^{-1}$ for $B_y = 1$ T, and about twice that value for $B_y = -2$ T (dashed line in Fig. 5b). From this, it follows that $\alpha - \beta_1 + \beta_3 \approx -0.3 \times 10^{-13}$ eV m, where the sign is directly determined by the sign of $\partial x_0/\partial t$. 

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We can now derive the size of all SO coefficients in our sample. From the known sum and difference of $\alpha$ and $\beta_1 - \beta_3$, respectively, we find $\alpha = (1.6 \pm 2.3) \times 10^{-13}$ eV m and $\beta_1 - \beta_3 = (1.9 \pm 2.6) \times 10^{-13}$ eV m. The cubic Dresselhaus SOI $\beta_3$ is estimated from equation (1) based on the observed $\tau_s$, yielding $\beta_3 \approx 0.7 \times 10^{-13}$ eV m. We then obtain $\beta_1 \approx (2.6 \pm 3.3) \times 10^{-13}$ eV m, which is in excellent agreement with results from a similar QW structure.\(^6\) Using the theoretical expression $\beta_1 = -\gamma \langle k_z^2 \rangle$ (ref. 1), with $\langle k_z^2 \rangle = 3.7 \times 10^{16}$ m$^{-2}$ as obtained by solving the one-dimensional Poisson and Schrödinger equations, we determine the Dresselhaus coupling parameter $\gamma \approx -9 \text{ eV Å}^3$. This is in agreement with a previous work where $\gamma$ was found to be in the range of $-4$ to $-8 \text{ eV Å}^3$ (ref. 8). Inserting $\gamma = -9 \text{ eV Å}^3$ into $\beta_3 = -(1/2)\gamma n_s$, with $n_s = 5 \times 10^{13}$ m$^{-2}$, we find $\beta_3 = 0.7 \times 10^{-13}$ eV m, exactly the same value as obtained from $\tau_s$. This indicates that the PSH decay is well described by equation (1). The Rashba coefficient $\alpha$ can be related to the electric field $E_{\text{QW}}$ in the QW by $\alpha = r_{\text{QW}} E_{\text{QW}}$, defining a proportionality

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**Figure 5** | Interplay of the PSH with an external magnetic field. a, Experimentally measured and numerically simulated maps of $S_z(y, t)$ for $B_y = -1 \text{ T}$. The formation of a helical spin mode is challenged by simultaneous spin precession about $B_y$. b, Maps of $S_z(x, t)$ recorded at $B_y = -1$ and $-2 \text{ T}$ show the spin precession about $B_y$ as well as an asymmetry of the precession phase with $\pm x$ (dashed lines), which is attributed to a slight detuning from the SU(2) symmetry point.

**Figure 6** | Detuning from the PSH regime. a, Left: maps of $S_z(y, t)$ for $B_y = -0.375 \text{ T}$. The PSH is still visible, but the spin lifetime has decreased to a few 100 ps. Right: maps of $S_z(y, t)$ for $B_y = -2 \text{ T}$. Instead of the formation of a PSH, a spin precession about $B_y$ is seen. b, Time traces $S_z(t)$ exhibit no spin precession for $|B_y| \leq 0.375$ mT, and $\tau_s$ decreases rapidly with $|B_y|$. c, Precession frequency $\nu$ and spin lifetime $\tau_s$ versus $B_y$. Open symbols: experimental data. Filled symbols: numerical simulation. Arrows indicate positions of maps shown in a.
constant $r_{\text{QW}}$. From the calculated conduction-band profile, we estimate $E_{\text{QW}} = 4 - 6 \times 10^{-3} \text{V m}^{-1}$. Using $\alpha \approx 2.3 \times 10^{-13} \text{eV m}$, we obtain $r_{\text{QW}} = 4 - 6 \text{Å}^2$, in good agreement with theoretical prediction\(^1\) and experiment\(^6\).

To cross-check the consistency of our experimentally determined SO coefficients, we simulate spin diffusion in the presence of SOI using a Monte Carlo approach that combines semiclassical spin dynamics and diffusion (see numerical Methods). The simulation uses the experimentally determined values for $\alpha$, $\beta_1$, $\beta_3$, and $D_s$, and reproduces the spin dynamics in the combined field of $B$ and $B_{\text{SO}}$ remarkably well (see Fig. 5a,b). It also confirms the analytical relation of equation (2) (dashed lines in Fig. 5b).

Finally we want to discuss the transition from a PSH-dominated regime to a regime in which the external magnetic field dominates. Figure 6a shows two maps of $S_z(x, y)$ recorded at $B_t = -0.375$ and $-2$. At $B_t = -0.375$, the formation of a helical mode is readily identifiable, but $\tau$, is significantly lower than for $0$ T (see Fig. 2c). At $-2$, no signature of a helical mode is observed and spins precess about $B_t$ with $v = [g \mu_B B_t / h]$. To determine $v$ and $\tau$, versus $B_t$, the time traces $S_z(t)$ at $(x, y) = (0, 0)$ (Fig. 6b) are fits to $(t - t_0)^{-1/2} \exp(-t/t_0) \cos(2\pi v t)$. Starting from the PSH at $0$ T, $\tau$ rapidly decreases with increasing $|B_t|$ (Fig. 6c). Interestingly, spin precession about $B$ is suppressed in the regime of $|B_t| \leq 0.5$ T. Even though at $|B_t| = 0.5$ T the Zeeman spin splitting (5 eV) is more than one order of magnitude smaller than the SO spin splitting (up to 170 eV at the Fermi energy), the symmetry breaking induced by $B$, dramatically decreases $\tau$, from 1 ns to below 100 ps (Fig. 6c).

In an intermediate regime of $0.5$ T $< |B_t| < 1.5$ T, the helical spin mode coexists with spin precession about $B$, and $\tau$, recovers from its minimum value. For $|B_t| > 1.5$ T, $\tau$, monotonically decreases. In the regime where the Zeeman energy approaches the SOI energy, we expect that $\tau$, approaches the value given by the Dyakonov–Perel expression for homogeneous polarization. Also shown in Fig. 6c are fits to numerical simulations of $S_z(t)$ at $(x, y) = (0, 0)$, which exhibit the same features as the measured data.

**Methods**

Experimental methods. The 2DEG investigated is confined to a 12-nm-wide (001)-oriented GaAs/AlGaAs QW placed 95 nm below the surface. Asymmetric Si modulation doping provides a sufficiently strong Rashba SOI to nearly balance the Dresselhaus SOI contribution. We employ a pump–probe approach in which a circularly polarized pulse of a mode-locked laser at $\lambda = 785$ nm (full-width at half-maximum 10 nm) excites spin-polarized electrons into the conduction band of the QW. The spin-polarization component $S_z$ along the $z \parallel [001]$ direction is measured by a probe pulse at $\lambda = 799 – 801$ nm using the polar magneto-optical Kerr effect\(^7\)-\(^10\). Both pulses are directed through a single high numerical aperture lens placed inside the cryostat. The spatial position of the pump beam on the sample surface is scanned by controlling the angle of the incident pump beam with a galvo mirror. In spatial maps of the spin distribution $S_z(x, y)$, the coordinates $x$ and $y$ indicate the position of the probe beam relative to the scanning pump beam. A nonlinearity in the angle control was corrected to obtain calibrated values for $x$ and $y$. The pump pulse is spectrally removed with a low-pass filter before detection. The focal distance is optimized to achieve a maximal spatial resolution of 2 μm in diameter. The probe pulse is synchronized to the pump pulse and delayed with a mechanical stage by a time $t$ to monitor the time evolution of the spin polarization $S_z(t)$. The pulse lengths of the pump and the probe beam are typically 60 ps and 3 ps, respectively. Typical power intensities of pump and probe were 250 and 50 W, respectively. Figure 3b includes data with the same probe power, $W$, respectively, with pulses arriving at a repetition rate of 200 ps. The pump pulse is spectrally removed with a low-pass filter before measurement. Wavelengths $(k_x, k_y)$ are uniformly distributed on the Fermi disc, and the spins $S$ are oriented along the $z$-direction. In this regime we steps, updates of $(x, y)$ and $S$ are calculated, treating the spin precession about the sum of $B_{\text{SO}} + B$ semiclassically. Scattering is accounted for by isotropically redistributing the charge carriers on the Fermi disc with scattering probability $v = 2D/v_F$, where $v_F$ is the Fermi velocity. For the simulation, we used a single set of SO coefficients, $\alpha = 1.7$, $\beta_3 = 2.7$, $\beta_1 = 0.7 \times 10^{-13} \text{eV m}$, and $D_S = 385 \text{cm}^2 \text{s}^{-1}$ to reproduce the measured $S_z$, as shown in Figs 5, 6 and Supplementary Fig. 2. The variation of the SO coefficients with $t$, as observed in the measurement, has not been included in the simulations.

Received 25 April 2012; accepted 3 July 2012; published online 12 August 2012

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**Acknowledgements**

We would like to acknowledge financial support from the Swiss National Science Foundation through NCCR Nano and NCCR QSIT, as well as valuable discussions with A. Alenbach, K. Ensslin and Y. Chen.

**Author contributions**

M.P.W. and G.S. designed the experiment, interpreted the data and wrote the manuscript. M.P.W. performed the time-resolved experiment. C.R. and W.W. grew the samples. G.S. performed numerical simulations.

**Additional information**

Supplementary information is available in the online version of the paper. Reprints and permissions information is available online at www.nature.com/reprints. Correspondence and requests for materials should be addressed to G.S.

**Competing financial interests**

The authors declare no competing financial interests.