Doppler velocimetry of spin propagation in a two-dimensional electron gas

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Controlling the flow of electrons by manipulating their spin is a key to the development of spin-based electronics. Recent demonstrations of electrical-gate control in spin-transistor configurations have shown great promise, but operation at room temperature remains elusive. Further progress requires a deeper understanding of the propagation of spin polarization, particularly in the high-mobility semiconductors used for devices. Here we report the application of Doppler velocimetry to resolve the motion of spin-polarized electrons in GaAs quantum wells driven by a drifting Fermi sea. We find that the spin mobility tracks the high electron mobility precisely as a function of temperature. However, we also observe that the coherent precession of spins driven by spin–orbit interaction, which is essential for the operation of a broad class of spin logic devices, breaks down at temperatures above 150 K, for reasons that are not yet understood theoretically.

The transistor, the iconic invention of twentieth century science, is a semiconductor device in which the flow of electrons is modulated by voltages applied via electrodes, known appropriately as gates. In a conventional transistor the gate electrode controls the number of mobile electrons in the current carrying pathway, or ‘channel’. In pursuit of transistors with faster response and lower rates of energy dissipation, there has been intense investigation aimed at modulating current through the manipulation of spin by applied electric fields, a coupling that occurs because of the spin–orbit (SO) interaction. Recently, gate-controlled modulation of current by means of SO coupling has been demonstrated in prototype device structures that operate below room temperature.

Further progress towards spintronic logic requires a deeper understanding of the basic physical principles on which such devices are based. Essentially the question is this: how far, and how fast, can spin polarization propagate in a current-carrying electron gas? This question was first addressed in pioneering work that used magneto-optic imaging to follow the drift of spin polarization packets in real space. These experiments were enabled by the enhanced spin lifetimes (in excess of 10 ns) that arise because of the metal–insulator transition of a doped semiconductor at the expense of electron mobility, . However, the high– electron gas needed for fast devices is in a very different dynamical regime, where spin lifetimes are ~10–100 ps, during which time spin may propagate only 10–100 nm (depending on the temperature, , and applied field, ). To resolve spin propagation on picosecond timescales and nanometre length scales we have developed a technique to measure velocity through the Doppler shift of light scattered from propagating waves of spin density. Our method extends transient grating spectroscopy (TGS), which has traditionally been used to measure rates of diffusion, to the measurement of drift velocity.

**Transient grating spectroscopy and Doppler velocimetry**

Excitation of a semiconductor with a single beam of photons with energy above the bandgap injects an equal population of electrons and holes, whose spatial distribution follows the intensity of the laser spot. Although recombination rates can be readily determined from the lifetime of this excited state, information about the motion of electrons and holes is inferred only indirectly. TGS adds the element of spatial resolution to time-resolved optical measurements, enabling direct probing of transport. In TGS, two non-collinear beams of light pulses interfere at the sample surface, creating a pattern of intensity and photon helicity that depends on the relative angle and polarization state of the two beams. When the two beams are polarized parallel to each other, interference creates a standing wave of laser intensity, generating the sinusoidal pattern of photo-induced electron–hole (e–h) density shown in Fig. 1a. On the other hand, orthogonal polarization generates a standing wave of photon helicity, while maintaining spatially uniform intensity (on the scale of the laser spot) and, therefore, e–h density. Optical selection rules in GaAs cause photon helicity to be imprinted in the out-of-plane (z) component of the angular momentum of the photo-induced e–h gas. Because hole spins depolarize on a subpicosecond timescale, the excited state for t > 1 ps comprises an electron spin density wave (SDW) accompanied by a charge compensating gas of unpolarized holes, as illustrated in Fig. 1b. The SDW induces variation in the local index of refraction and therefore acts as a transient optical grating, whose subsequent dynamics can be monitored by the diffraction of a time-delayed probe pulse.

In the experiments reported here we photo-inject the SDW into a two-dimensional electron gas (2DEG) subject to an in-plane field that is parallel to the grating wave vector and measure the resulting propagation of spin polarization. If the polarization wave undergoes normal drift and diffusion, the spin density will evolve according to

\[ S(x, t) = S_0 \exp(-t/\tau(q)) \cos(qx - x_0(t)), \]

where \( q \) is the wave vector and \( x_0(t) \) is the displacement. Measurement of the amplitude of a diffracted pulse yields the wave vector dependent lifetime \( \tau(q) \), from which the spin memory time and diffusion coefficient can be determined. Information about \( x_0(t) \) is contained in the phase, rather than the amplitude, of the diffracted light. For example, light diffracted from an SDW drifting at constant velocity, such that \( x_0(t) = v_d t \), will contain the optical phase factor \( \phi(t) = qv_d t \). The linear advance of phase with time is equivalent to a Doppler shift, \( \Delta \omega = v_d q \), as illustrated in Fig. 2. To measure the phase of the
that induces spin precession at a rate that depends on the electron’s momentum, \( \mathbf{p} \). In a symmetric QW the Dresselhaus SO coupling\(^1\) dominates, which is characterized by the precession rate vector field, \( \mathbf{g}(\mathbf{r}) = 2\hbar^{-1} \mathbf{\beta}_1 (p_x \hat{x} + p_y \hat{y}) \), where \( \mathbf{\beta}_1 \) is the linear Dresselhaus coupling strength and \( \hat{x} \) and \( \hat{y} \) are the [110] and [110] crystal axes, respectively. The connection between the precession vector and momentum induces a strong correlation between the diffusion of electrons in real space and of the orientation of their spins on the Bloch sphere. Theoretical analysis of this correlation yields a pair of normal modes at each \( q \) that are helical waves of spin density with opposite sense of rotation\(^2\)–\(^7\). The lifetime of the helical whose sense of rotation matches that of the electron’s precession is strongly enhanced by SO coupling whereas the lifetime of the helix with opposite rotation is reduced. Both lifetimes, \( \tau_{\pm}(q) \), are observed in the TSG experiment because the photogenerated initial state, a wave of pure \( \pm \), is a superposition of the two helices of opposite pitch. The solid lines through the data in Fig. 3b are fits to the spin helix theory\(^13\)–\(^15\) with \( \mathbf{\beta}_1 = 3.4 \times 10^{-3} \text{ eV Å} \).

From analysis of the measured \( \tau_{\pm}(q) \) we also obtain the spin diffusion coefficient, \( D_s(T) \), plotted in Fig. 3c. For comparison we plot the electron diffusion coefficient, \( D_e(T) \), obtained by applying the Einstein relation to the electron transport mobility \( \mu_e \). \( D_s(T) \) is smaller than \( D_e(T) \) as a result of spin Coulomb drag (SCD)\(^18\), which is a frictional force between oppositely oriented spins that is generated by electron–electron (e–e) collisions. As spin diffusion requires a counterflow of opposite spin populations, it is damped by SCD, whereas charge transport is protected from e–e collisions by momentum conservation. The reduction of the spin diffusion coefficient relative to the electron diffusion coefficient seen here is considerably larger than in previously reported measurements\(^19\) — \( D_s \) is only about 5\% of \( D_e \) measured above the Fermi temperature of 80 K. The SCD effect is more pronounced in the cleaner sample studied here because, whereas its low-\( T \) resistivity is approximately eight times smaller than the previously studied QW with the same electron density, the intrinsic spin-drag transresistivity \( \rho_{11}(T) \) is unchanged. This is evident when we invert the measured \( D_s/D_e \) to extract the transresistivity (See Supplementary Section SI). The \( \rho_{11}(T) \) thus obtained (plotted in Fig. 3c inset) is quantitatively consistent with earlier reports and in excellent agreement with the random phase approximation (RPA) theory of SCD in two dimensions in the range of \( T \) below the Fermi temperature\(^20\),\(^21\).
Spin drift

We turn now to Doppler shift measurements of spin helix drift under the influence of an $E$ field applied parallel to the SDW wave vector. Figure 4a shows the phase, $\phi(q,t)$, of light diffracted from a transient spin grating as a function of $t$ for several values of $q$, measured at 30 K. For wave vectors larger than $q_0$, the phase increases linearly with time, indicating near-uniform drift in the same direction as the Fermi sea of electrons. However, $\phi(q,t)$ is clearly more complex for $q < q_0$. Although $\phi$ starts out positive, it quickly crosses zero and becomes negative for $t \geq 50 \text{ ps}$, indicating counter-propagation with respect to the Fermi sea. It is natural to associate the anomalous behaviour of the phase with the presence of the two helical modes discussed previously, and to describe the overall $\phi(q,t)$ as the weighted average of their individual rates of phase advance, $\phi_\pm(q)$,

$$\phi(q,t) = \sum_{i=\pm} \frac{A_i \exp[-t/\tau_i(q)] \phi_i(q)}{\sum_{i=\pm} A_i \exp[-t/\tau_i(q)]}$$

(1)

The lines through the data in Fig. 4a are fits obtained with this expression, using the values of $\tau_i(q)$ obtained previously. The high quality of the fits suggests that the complicated behaviour of $\phi(q,t)$ reflects contributions from the two helices of opposite pitch, each propagating with its own uniform phase velocity. Shown in Fig. 4b are values $\phi_\pm(q)$ obtained using equation (1); the solid lines are fits to a theory of spin helix propagation described qualitatively below.

Whereas for a Fermi sea at rest the average of $\Omega(p)$ over occupied states is clearly zero, an electron gas drifting with velocity $v_D \mathbf{k}$ will experience a nonzero $\langle \Omega \rangle = 2\beta \mathbf{h} \cdot m^* v_D \mathbf{k}$, where $m^*$ is the effective mass of electrons in the conduction band. Consider first, spin helices injected into a drifting Fermi sea in the absence of SO coupling. In this case the angle, $\theta$, of the local spin polarization with respect to $\mathbf{k}$ would be static in a frame moving with $v_D \mathbf{k}$, $\theta(x',t) = \pm q x'$, where $x' = x - v_D t$. However, the nonzero SO coupling will cause the spins to precess as they drift, such that $\theta(x',t) = \pm \alpha (x' + t)$. When viewed in the stationary frame, $\theta(x,t) = \pm (q - v_D q_0) t$, where $q_0 = 2\beta \mathbf{h} \cdot m^* \mathbf{k}$, which corresponds to the two rates of phase advance, $\phi_\pm(q,t) = \phi_\pm(q \pm q_0)$. Thus, for example, the long-lived helix will seem to be stationary when $q = q_0$ and counter-propagate for $q < q_0$. The lines through the data points in Fig. 4b are the predictions of quantitative theories of helix drift22,23, which differ from the qualitative picture outlined above only at very low values of $q$ not accessible in our experiments.

The propagation of spin in the Dresselhaus field can also be visualized in the spatial rather than wave-vector domain. As we have shown, TGS measures $S_1(q,t)$ over a broad range of $q$. The Fourier transform of $S_1(q,t)$ yields the real-space spin propagator, $S_1(x,t)$, which describes the time evolution of spin polarization following pulsed injection of a narrow stripe of $z$-oriented spin density along the y axis. Fourier transformation of the theoretical fits to the amplitude and phase of $S_1(q,t)$ shown in Figs 3b and 4b yield the propagator illustrated in Fig. 4c, which has the form of an envelope function that moves with uniform velocity $v_D$ while modulating a stationary SDW (see refs 22,23 and Supplementary Section SI). The polarization wave that emerges as the envelope propagates is closely related to the stripe-like patterns imaged in steady-state measurements on low-$\mu_0$ semiconductors24.

A surprising feature of our results is that the spin propagation dynamics described above change drastically as $T$ is increased towards room temperature. As shown in Fig. 4d,e, at 150 K the two helices with phase dispersion $\phi(q,t) = v_D(q \pm q_0) t$ have been replaced by a single mode with $\phi(q,t) = v_D(q)t$. The latter corresponds to drift with velocity $v_D$ without spin precession. Fourier transformation of fits to $S_1(q,t)$ at 150 K yields a spin packet that propagates at $v_D$ and does not modulate a polarization wave, as shown in Fig. 4f (see Supplementary Section SI). Thus a spin-transistor based on control of SO-induced precession will not operate in this $T$ regime. We note that the clear crossover in spin dynamics that has taken place can be seen only in spatial or wave-vector-resolved measurements, as the lifetime of the uniform spin polarization, $S_1(q = 0, t)$ increases monotonically from 30 to 150 K (ref. 25).

The absence of spin precession at 150 K cannot be attributed to a change in the SO coupling strength, $\beta$, which is an intrinsic property of the GaAs band structure. Instead, our results suggest that the effective precession vector $\langle \Omega \rangle$ does not survive increased thermal averaging. One possible reason for this is the cubic (in $p$) Dresselhaus coupling, which causes the net precession angle between scattering events to depend on the electron’s velocity, and has been shown to degrade the spin-spatial correlations described previously17,26.

Temperature dependence of spin mobility

We have seen that, when viewed in the spatial domain, an injected spin packet moves with $v_D = \partial \phi / \partial q$, regardless of whether the propagation is accompanied by coherent spin precession. Thus at each $T$ we can determine a spin packet velocity from the dispersion of $\phi(q)$, obtain a spin packet mobility, $\mu_s = v_D/E$, and compare it with the electron mobility, $\mu_e$, as determined from d.c. transport. In
the course of such measurements we discovered that $\mu_s$ depends strongly on the intensity, $I$, of the laser pulse that generates the spin grating. Figure 5a is a plot of $\mu_s$ as a function of $I$ for various $T$. As $I$ is reduced from its maximum value $I_0 = 0.25 \mu \text{J cm}^{-2}$, $\mu_s$ initially increases and then approaches an asymptotic value $\mu_{s0}$ in the limit that $I \to 0$. The curves through the data points are fits to the relation, $\mu_s(I) = \mu_{s0}(T)/(1 + 2.86I/I_0)$. The red circles in Fig. 5b represent $\mu_{s0}(T)$ as determined from fits to the intensity dependence. The plot shows that $\mu_{s0}(T) = \mu_s(T)$ over the entire $T$ range of the experiment. Furthermore, this equality holds even as the nature of spin propagation crosses over from the precession regime, where $\phi_s(q) = v_d(q \pm q_0)$, to the incoherent regime, where $\phi_s(q) = v_d q$.

We argue below that the dependence of $\mu_s$ on $I$ indicates that the direct force of the electric field on the spin polarization is zero, and that spin waves (or packets) are propelled solely by momentum transfer from the surrounding Fermi sea. The basis of this claim is that the same dependence of mobility on pump intensity is observed when the Fermi sea drives another neutral excitation of the 2DEG, namely packets of e-h density\textsuperscript{27,28}. Individual packets of electron and hole density cannot separate in weak applied fields—the constraint of local charge neutrality forces the two charge species to drift together at a speed $v_d$, where $v_d$ is the ambipolar mobility. Because the driving force of the Fermi sea scales with the equilibrium carrier concentration, $n_0$, and the packet’s inertia varies as the photo-induced carrier concentration, $\Delta n$, the packet drift velocity depends on the ratio $\Delta n/n_0$, which is proportional to $I$. Solving the appropriate force balance equations yields $\mu_s \propto I^{-1}$ in the limit that $\Delta n \gg n_0$ and $\mu_s \to \mu_{s0}$ in the limit $\Delta n \ll n_0$ (refs 27,28), that is, the same trends that we observe in the spin packet mobility. In Fig. 5b we compare $\mu_{s0}(T)$ determined from TGS measurements on the same QW using parallel polarization (see Fig. 1a) with the electron and spin mobility. The slower rate of ambipolar propagation reflects the fact that, in contrast with electron spin propagation, the low-mobility holes must be dragged along with the drifting electrons.

Overall, the phenomena described above demonstrate that spin density propagates in a Fermi sea as a distinct, neutral degree of freedom that couples to electron motion through various interactions (for example SO and Coulomb). Figure 1c illustrates the distinct nature of the spin polarization—given sufficient time,
the more rapidly drifting SDW will leave the unpolarized packet of e−h density in its wake. In the limit that the number of photoexcited carriers is much less than the equilibrium number, we find that spin density propagates at the drift velocity of the Fermi sea. Although this result has been predicted theoretically in models that exclude SO-induced precession\(^{29,30}\), it is striking that this equality is preserved even as spin dynamics cross over from incoherent relaxation to coherent precession with decreasing T. Hopefully, the ability to measure spin mobility through the Doppler effect will encourage a deeper understanding of the underlying physics of spin propagation in SO-coupled metallic systems, providing the basis for extending the temperature range of spin-based logic.

**Methods**

**Sample preparation.** The experiments were performed on a 9-nm-wide n-doped GaAs/AlGaAs quantum well grown by molecular beam epitaxy on a semi-insulating GaAs (001) substrate (VB0355). A 200 nm Al\(_{0.33}\)Ga\(_{0.67}\)As etch stop layer was grown on the substrate, followed by a second 10-nm etch stop of GaAs. The lower barrier was 210 nm Al\(_{0.33}\)Ga\(_{0.67}\)As with Si \(\delta\)-doping 95 nm below the 9 nm GaAs quantum well. The upper barrier of 190 nm Al\(_{0.33}\)Ga\(_{0.67}\)As includes a \(\delta\)-doped layer 75 nm above the quantum well. The top layer is a 10 nm GaAs cap. The slight asymmetry in the \(\delta\)-doped layers compensates for the upward drift of the Si atoms and results in nearly symmetric doping when the growth is complete. The 2DEG channel was defined by mesa etching, and ohmic contact was made by annealing NiGeAu into the sample. After patterning, the sample was epoxied onto a non-birefringent sapphire disk and mechanically thinned to 50 \(\mu\)m. A citric acid etch removed the remaining GaAs substrate. The final thickness was reached after a second selective etch with hydrofluoric acid to remove the Al\(_{0.33}\)Ga\(_{0.67}\)As layer. After processing, the carrier density and mobility of the 2DEG in the dark were 1.9 \(\times\) 10\(^5\) cm\(^{-2}\) and 5.5 \(\times\) 10\(^5\) cm\(^2\)/V s\(^{-1}\) at 5 K, respectively, as measured by standard van der Pauw techniques at Sandia National Laboratories. The sample resistance was also monitored during the course of the optical measurements at Lawrence Berkeley National Laboratory. In the temperature range of the experiments reported here, very little difference in the sample resistance was observed between light and dark cool-down conditions (see Supplementary Section SIII).

**Error analysis.** The error bars in Fig. 5 stem from the uncertainty in fitting the Doppler velocimetry data such as that in Fig. 4a and d. The uncertainty in these fits (barely visible error bars in Fig. 4b and e) was used to weight subsequent fitting. The uncertainty in the carrier density and mobility of the 2DEG in the dark were 1.9 \(\times\) 10\(^5\) cm\(^2\)/V s\(^{-1}\) and 5.5 \(\times\) 10\(^5\) cm\(^2\)/V s\(^{-1}\) at 5 K, respectively, as measured by standard van der Pauw techniques at Sandia National Laboratories. The sample resistance was also monitored during the course of the optical measurements at Lawrence Berkeley National Laboratory. In the temperature range of the experiments reported here, very little difference in the sample resistance was observed between light and dark cool-down conditions (see Supplementary Section SIII).

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**Author contributions**

L.Y., J.D.K., J.O. and M.P.L. devised the experiment and wrote the manuscript. L.Y. and J.D.K. performed optical measurements and M.P.L. and L.Y. carried out transport measurements. L.Y. and J.O. performed analysis and theoretical modelling of the data. J.L.R. and D.R.T. carried out growth and fabrication of the quantum–well device.

**Additional information**

The authors declare no competing financial interests. Supplementary information accompanies this paper on www.nature.com/naturephysics. Reprints and permissions information is available online at http://www.nature.com/reprints. Correspondence and requests for materials should be addressed to J.O.