

Ballistic Spin Valve in Graphene Realized via Electron Optics

Daniel Burrow¹, Oktay Deveci¹, Rares Dragomir¹, Thomas Thomson², and Ivan J. Vera-Marun^{1,*}

¹*Department of Physics and Astronomy, University of Manchester, Manchester M13 9PL, United Kingdom*

²*Nano-Engineering and Spintronic Technologies (NEST), Department of Computer Science, University of Manchester, Manchester M13 9PL, United Kingdom*



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Graphene is an ideal platform for supporting ballistic, spin-coherent transport—crucial ingredients for spin-based quantum technologies. Here, we demonstrate a spin-dependent electron-optic mechanism that enables spin-transistor-like control of spin transport in a graphene spin valve, in the absence of spin-orbit coupling in the transport channel. Using transverse magnetic focusing in a high-mobility graphene device, we couple spin and orbital degrees of freedom through charge transfer doping and proximity exchange at ferromagnetic contacts. This coupling allows for gate-tunable modulation of both the amplitude and polarity of the spin signal, reminiscent of the functionality of a Datta-Das spin field-effect transistor. Measurements confirm that spin-coherent ballistic transport is observable up to room temperature. Our results demonstrate an operational principle for spintronic devices based on spin-dependent electron optics in low spin-orbit coupling materials such as graphene.

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I. INTRODUCTION

Advancements in the electronic quality of graphene heterostructures have enabled 2D devices that show high charge carrier mobility and micrometer-scale ballistic transport up to room temperature [1–4]. Increased material quality has also had a substantial impact on the field of graphene spintronics [5], where enhanced spin transport has been observed in devices encapsulated in hexagonal-boron nitride (*h*-BN) [6–9], due to a reduction in spatial disorder relative to graphene in direct contact with a SiO₂ substrate [10,11]. The momentum scattering of spin carriers in the diffusive transport regime contributes to spin relaxation in graphene [12–16], motivating the exploration of spin transport in the ballistic regime. Theoretical work has highlighted how the diffusive formalism for spin dynamics [17] underestimates spin transport parameters in the ballistic regime [18], while a limited number of experimental studies in semiconductors indicate enhanced spin signals for high-mobility devices [19,20]. Hence, the crossover from diffusive to ballistic spin transport is an enabling step toward the realization of practical spintronic devices.

Long-range ballistic transport in graphene has been exploited to implement electron optics experiments, such as transverse magnetic focusing (TMF) [21,22], Veselago lensing [23,24], and spin-dependent beam splitting [25,26]. In such experiments, electrons injected by a point contact are bent by external out-of-plane magnetic fields and reflected or refracted from potential barriers [27–30]. Electron optics have been leveraged to study changes in the graphene band structure in the presence of proximity induced spin-orbit coupling (SOC) [26] and the formation of moiré minibands in twisted graphene bilayers [31]. Furthermore, such techniques have lead to the development of novel two-dimensional quantum devices, such as microcavities [32] and interferometers [33].

Here, we employ transverse magnetic focusing to unequivocally demonstrate ballistic, spin-polarized transport in a high-quality graphene spin valve. We show nonlocal ballistic transport over micrometer scales at low temperature, with quasiballistic transport present up to room temperature. The observed TMF signal is sensitive to the magnetic configuration of the contacts, indicating that the ballistic signal is spin dependent. An enhanced spin signal is observed under optimized electron optics focusing conditions, relative to a standard spin-valve measurement in the absence of an out-of-plane magnetic field. We also demonstrate electrostatic control over the amplitude and polarity of the focused spin signal by varying a back gate voltage. These results demonstrate a versatile device architecture for studying fundamental aspects of electronic spin dynamics in the ballistic transport regime and represent a practical route to developing a spintronic

*Contact author: ivan.veramarun@manchester.ac.uk

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transistor, along the lines of the Datta-Das spin-FET proposal [34], where electron optics can be incorporated as a means of modulating spin signal, in the absence of SOC in the transport channel.

II. RESULTS

Figure 1(a) shows a schematic of the spin-valve device used in this study, set up in the nonlocal measurement configuration, with nonlocal signal defined as $R_{\text{NL}} = V_{\text{NL}}/I$. The device is fabricated using methods described in prior work [9]. Nanowire electrodes of cobalt (Co, 30 nm), capped with gold (Au, 50 nm), define magnetic one-dimensional (1D) contacts at the exposed edges of the encapsulated graphene channel [9,35,36]. The nanoscale widths of the electrodes (200–350 nm) mean that they behave as pointlike contacts [37]. All measurements in this paper are performed at $T = 25$ K, unless stated otherwise.

Figure 1(b) shows a TMF map, where nonlocal resistance R_{NL} is plotted as a function of out-of-plane magnetic field $B_{\perp}$ and carrier density in the graphene channel, n . Carrier density is controlled by applying a back gate voltage, V_{BG} , to the p -doped silicon substrate, separated from the h -BN–graphene– h -BN heterostructure by 290 nm of SiO_2 (see Sec. V). A series of periodic peaks are present in the top right ($B_{\perp} > 0, n > 0$) and bottom left ($B_{\perp} < 0, n < 0$) quadrants of the map, seen as red stripes in the region between $-0.5\text{ T} < B_{\perp} < 0.5\text{ T}$, which occur when ballistic charge carriers are focused on the detector contact d . Under the Lorentz force exerted by $B_{\perp}$, carriers injected at contact i follow cyclotron orbits with radius $r_c = \hbar k_F / eB_{\perp}$, which can be tuned to the focusing

condition $L = j \times 2r_c$, where k_F is the Fermi wave number, j is an integer, and L is the separation between contacts [Fig. 1(c)] [21,22]. Provided that no scattering occurs along their trajectories, the charge carriers are collected at contact d , causing a spike in nonlocal voltage. The value of $B_{\perp}$ required to focus charge carriers on contact d is given by [22]

$$B_{\perp}^j = \pm j \frac{2\hbar k_F}{eL}. \quad (1)$$

Higher-order modes ($j > 1$) occur due to specular reflections from the edges of the 2D channel [Fig. 1(c)]. The emergence of such peaks is direct evidence of ballistic transport over a distance of at least $\ell_{\text{mfp}} = j\pi r_c$. In the present case, the contacts are spaced by $L = 1.9\text{ }\mu\text{m}$, implying that the mean free path of the device is $\ell_{\text{mfp}} \gtrsim 3\text{ }\mu\text{m}$. The expected positions of the first three focusing modes, calculated via Eq. (1), are indicated in Fig. 1(b) by black dashed lines, which align closely to the red stripes representing the measured peaks.

Figure 2(a) shows TMF signals at a series of carrier densities for either electron transport [$n > 0$, (ii) and (iv)] or hole transport [$n < 0$, (i) and (iii)]. Before each measurement, the magnetization configuration of the ferromagnetic injector and detector electrodes is set to be either parallel [$\uparrow\uparrow$, (iii) and (iv)] or antiparallel [$\uparrow\downarrow$, (i) and (ii)] by sweeping an in-plane magnetic field $B_{\parallel}$ along their length (coinciding with the easy axis of magnetization). Distinct TMF spectra are observed for opposite charge carrier polarities, with the hole spectra exhibiting magnetic focusing peaks of larger magnitude in each case.

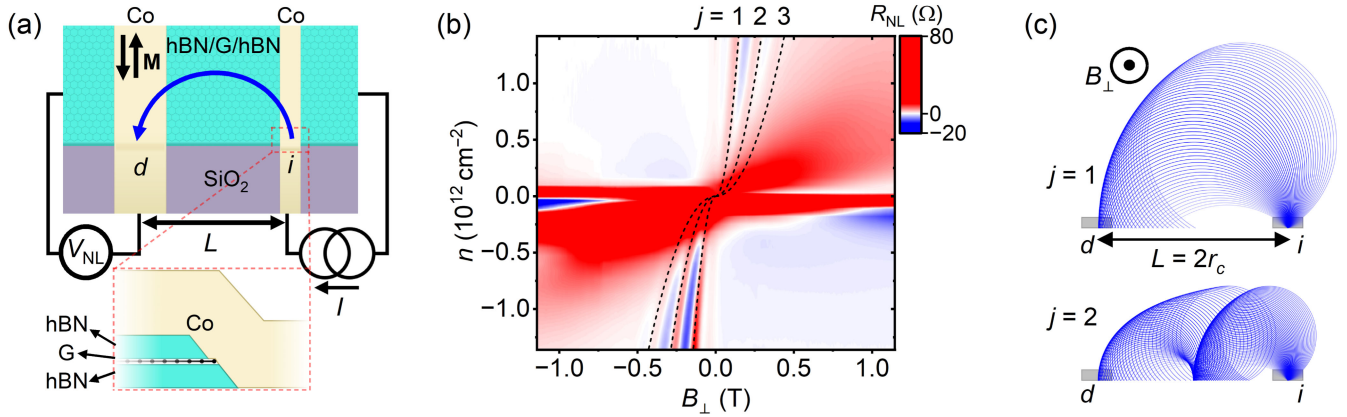


FIG. 1. Measurement concept and TMF map. (a) Schematic of an h -BN encapsulated graphene device on a Si/SiO₂ substrate, set up in the nonlocal measurement configuration. Co nanowires with magnetization $\mathbf{M}$ form magnetic 1D contacts at the exposed graphene edge, which inject (i) and detect (d) spin-polarized current; the inset depicts a profile of the heterostructure and 1D contact. A single-electron cyclotron orbit from i to d is illustrated (blue arrow). (b) Nonlocal resistance R_{NL} measured as a function of carrier density n and perpendicular magnetic field strength $B_{\perp}$ at low temperature (25 K). Three characteristic peaks ($j = 1, 2, 3$) are resolved in R_{NL} ; dashed black lines indicate the focusing conditions given by Eq. (1). (c) Illustration of ballistic charge carrier trajectories (blue lines), for the first ($j = 1$) and second ($j = 2$) focusing conditions. A perpendicular magnetic field $B_{\perp}$ tunes the cyclotron radius of charge carriers injected at i , focusing the majority of trajectories onto a caustic at d . Peaks in R_{NL} occur when the cyclotron radius is an integer multiple of half the contact separation ($L = j \times 2r_c$).

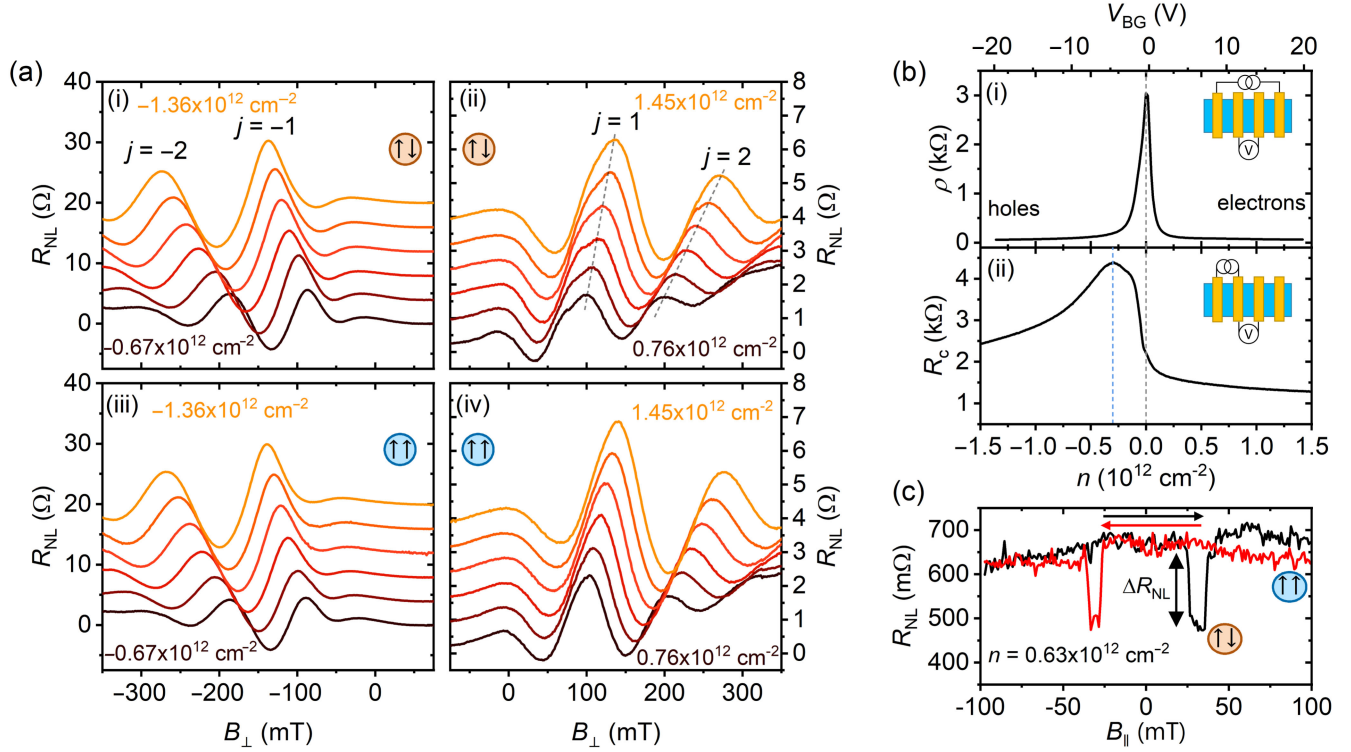


FIG. 2. Ballistic charge and spin transport at low temperature (25 K). (a) TMF signal for a range of carrier densities in the hole [(i) and (iii)] and electron [(ii) and (iv)] transport regimes, compared in the parallel [(iii) and (iv)] and antiparallel [(i) and (ii)] magnetic configurations. For presentation, each curve in (i) and (iii) is consecutively offset by 0.75 V, while in (ii) and (iv) the offset is 4 V—in each case, the lowest carrier density (black) represents the true baseline. Dashed lines in (ii) are a guide to the evolving positions of the $j = 1$ and $j = 2$ focusing peaks. (b) Local charge transport measurements: (i) graphene resistivity, ρ , measured in the four-terminal configuration (inset), and (ii) contact resistance, measured in the three-terminal configuration (inset). Both are plotted as a function of carrier density in the channel (back gate voltage), n (V_{BG}). Dashed lines in (i) and (ii) represent the voltages at which the graphene Dirac point ($V_D = -0.42$ V) and contact resistance maximum ($V_c = -4.62$ V) occur, respectively. (c) Spin-valve measurement— R_{NL} is measured as a function of in-plane magnetic field $B_{||}$. The parallel state ($\uparrow\uparrow$), antiparallel state ($\uparrow\downarrow$), and the magnitude of the spin signal, ΔR_{NL} , are indicated.

Remarkably, there is also a change in peak value between magnetic configurations, for the same doping polarity. For example, at $n = 0.76 \times 10^{12} \text{ cm}^{-2}$ in (iv) ($\uparrow\uparrow$), a peak value of $R_{NL} = 2.3 \Omega$ is measured at $j = 1$, while in the case of (ii) ($\uparrow\downarrow$) the peak value is $R_{NL} = 1.7 \Omega$ at $j = 1$ (peak values measured with respect to $R_{NL} = 0$). The difference between these values gives a spin signal of $\Delta R_{NL} \approx 600 \text{ m}\Omega$, representing an approximately 25% change in nonlocal resistance arising solely from a change in the magnetic configuration of the contacts. We focus on $j = 1$, as this is the only trajectory without edge scattering that can flip spin polarization. This change in focusing peak magnitude is direct evidence of spin-dependent ballistic transport.

Another distinctive feature is an apparent splitting of the first-order focusing peak observed in the electron transport regime. All TMF curves in Fig. 2(a)(ii), taken in the $\uparrow\downarrow$ configuration, show a clear double peak structure in the $j = 1$ peak. This structure is present, yet weaker for the curves shown in Fig. 2(a)(iv), taken at the same carrier

densities in the $\uparrow\uparrow$ configuration. This dependence of the structure of the focusing peak on the magnetic configuration further confirms the presence of spin-dependent focusing and implies that the double peak is the sum of two components, arising from distinct focusing spectra for each of the two spin channels. Neither magnetic configuration in the hole transport regime shows evidence of a double peak structure, indicating that the spin-dependent spectra are more similar in this case. Such contrast between TMF line shapes for electron and hole transport correlates with the distinct magnitudes of the peaks for each carrier polarity, with sharper focusing peaks in the hole regime.

Figure 2(b)(i) shows graphene resistivity ρ as a function of carrier density. Channel quality is quantified by the mobility μ , extracted by taking the gradient of conductivity $\sigma = 1/\rho$ with respect to carrier density. At low temperature (25 K), we extract a value of $\mu = 270 \times 10^3 \text{ cm}^2 \text{ V}^{-1} \text{ s}^{-1}$, a record high for graphene spintronic devices [9]. From diffusive transport data, we estimate the mean free path

to be $\ell_{\text{mfp}} \approx 1.25 \mu\text{m}$ (Supplemental Material Sec. SI [38]). The high mobility value and the observation of TMF over cyclotron orbits on the order of $\pi r_c \approx 3 \mu\text{m}$ both imply that ℓ_{mfp} is greater than the diffusive estimate, which is most likely limited by edge scattering due to the narrow channel width, $W \approx 1.5 \mu\text{m}$.

Contact resistance R_c for contact i is plotted as a function of carrier density in the channel in Fig. 2(b)(ii). The maximum of R_c is shifted toward negative n , indicating n -doped graphene in the vicinity of the contact, arising from charge transfer from the Co nanowires [9,39]. The offset between the neutrality point in the graphene channel and the maximum of R_c is $\delta n_c \approx 3 \times 10^{11} \text{ cm}^{-2}$. The significant modulation with varying n indicates that the value of R_c is determined by both the interfacial properties of the nanowire-graphene junction and the doped region of graphene adjacent to it [9]. This, along with the electron-hole asymmetry of R_c , evidences the formation of a tunable

p - n junction around the 1D contacts, expected to extend a short distance into the channel ($\leq 100 \text{ nm}$) [37,40–42].

A nonlocal spin-valve measurement, taken from the same region, is shown in Fig. 2(c). In this case, R_{NL} is measured as a function of $B_{\parallel}$, which causes sudden changes in R_{NL} when the magnetization of either electrode reverses relative to the other. The spin signal is defined as $\Delta R_{\text{NL}} = R_{\text{NL}}(\uparrow\uparrow) - R_{\text{NL}}(\uparrow\downarrow)$, which, for the spin valve presented here, gives $\Delta R_{\text{NL}} = 190 \pm 16 \text{ m}\Omega$.

III. DISCUSSION

A. Tunable spin-dependent electron optics

Figure 3(a) shows focusing signals taken in the parallel ($\uparrow\uparrow$) and antiparallel ($\uparrow\downarrow$) magnetic configurations for a single carrier density in the (i) p -doped and (ii) n -doped transport regimes. The difference between these signals isolates the spin-dependent component ΔR_{NL} (green line,

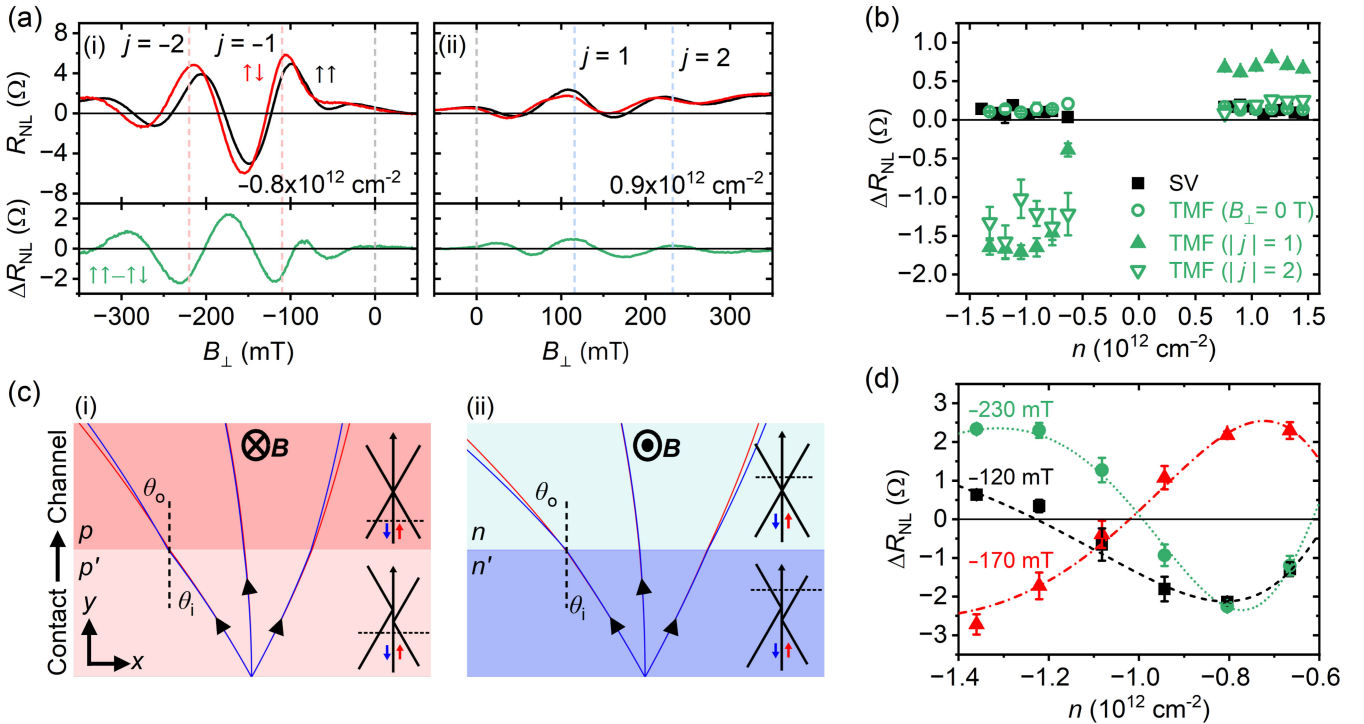


FIG. 3. Analysis of spin-dependent focusing spectra. (a) Parallel (black, $\uparrow\uparrow$) and antiparallel (red, $\uparrow\downarrow$) TMF signals for a single carrier density in the cases of (i) p doping and (ii) n doping. The spin-dependent signal is calculated as $\uparrow\downarrow - \uparrow\uparrow$ (green, bottom). Vertical dashed lines indicate the nominal focusing conditions at which ΔR_{NL} is evaluated for comparison with spin-valve data. (b) Comparison of ΔR_{NL} extracted from spin-valve measurements (black squares, SV), and from the focusing signals at $B_{\perp} = 0$ (green open circles), at the first-order mode (green closed triangles, $j = 1$) and at the second-order mode (green open triangles, $j = 2$). (c) Idealized schematics of carrier refraction across the contact-channel boundary in the case of (i) a p -doped and (ii) an n -doped graphene channel, where carriers are injected from a single point and move radially outward across the contact-channel boundary; θ_i (θ_o) is the incoming (outgoing) angle. The contact region is consistently more n doped than the channel due to proximity to the Co electrode. Spin-up (red) and spin-down (blue) trajectories have different refraction conditions due to an exchange split density of states in the contact region, as indicated by the schematic Dirac cones (stronger doping indicated by darker background). (d) Modulation of focused spin signal, ΔR_{NL} , via tuning of the carrier density in the channel, n , at three fixed perpendicular magnetic field strengths: -120 (black squares), -170 (red triangles), and -230 mT (green circles). Dashed (-120 mT), dotted (-230 mT), and dash-dotted (-170 mT) lines are sinusoid fits to the data, described in the text.

lower panel). We extract the value of this focused spin signal at three perpendicular magnetic field strengths: (i) at $B_{\perp} = 0$, (ii) at the first-order focusing condition ($|j| = 1$), and (iii) at the second order ($|j| = 2$). ΔR_{NL} evaluated at these points in the TMF experiment is plotted as a function of carrier density in Fig. 3(b) and compared with values extracted from standard nonlocal spin-valve measurements, where the magnetic field is oriented in plane. The values obtained from the magnetic focusing spectra at $B_{\perp} = 0$ are consistent with those extracted from the spin-valve measurements, as expected due to the absence of TMF in both cases.

As the magnitude of $B_{\perp}$ is increased, the spin signal ΔR_{NL} exhibits an oscillatory response as a function of the magnetic field, implying it is modulated by TMF [Figs. 3(a)(i) and 3(a)(ii)]. This is demonstrated by the extracted value of ΔR_{NL} at the nominal first- and second-order focusing conditions [Fig. 3(b)]. The spin signal evaluated at $|j| = 1$ shows two distinct differences compared to the standard spin-valve measurement. First, there is a significant enhancement in magnitude for both doping polarities. For n doping, there is an enhancement of $\geq 5\times$ throughout the carrier density range explored, reaching a maximum of $\Delta R_{\text{NL}} = 0.79 \pm 0.02\Omega$ at $n \approx 1.2 \times 10^{12} \text{ cm}^{-2}$. For p doping, there is a greater enhancement of up to $-20\times$, reaching $\Delta R_{\text{NL}} = -1.71 \pm 0.10\Omega$ at $n \approx -1.0 \times 10^{12} \text{ cm}^{-2}$. Second, the presence of TMF leads to an inversion of the spin signal at the $j = -1$ focusing mode (p doping), whereas the spin signal remains of a positive polarity at $j = 1$ (n doping). Thus, TMF introduces an electron-hole asymmetry in the spin signal that is absent in the standard spin-valve measurement and at $B_{\perp} = 0$. A similar behavior is observed for the spin signal at $|j| = 2$, albeit with a lower enhancement in magnitude, in particular, for n doping [Fig. 3(b)].

The enhanced spin signal is explained, in part, by the suppression of momentum scattering during the TMF experiment, which reduces contributions to spin relaxation via the Elliot-Yafet mechanism [15]. Although nonlocal transport is also ballistic during a standard spin-valve experiment ($\ell_{\text{mfp}} \geq 3 \mu\text{m}$), this can be considered as a chaotic billiard model where electrons undergo random scattering from the device edges [43]. Furthermore, the dynamics of magnetic focusing occur on the picosecond timescale, faster than that of spin precession due to external fields (Supplemental Material Sec. SII [38]) and the typical timescale of spin relaxation due to intrinsic SOC fields (the D'yakonov-Perel mechanism) [13,44]. Suppression of these common spin relaxation mechanisms partially accounts for an enhanced spin signal but cannot explain the differences in focusing peak positions observed for the parallel and antiparallel magnetic configurations or for the resultant reversal of spin polarity.

Spin-dependent magnetic focusing is commonly understood to arise due to a spin-split band structure owing to SOC

within the spin transport channel [25,45,46], causing a spatial separation of the focal points for spin-up and spin-down polarized carriers. However, in the absence of proximity effects within the graphene channel, there is negligible intrinsic SOC strength [47], and any such spin-dependent spatial separation of the focal points would not be resolved in the present experiment. Instead, we propose that the changes in the magnitude, polarity, and structure of the focusing peaks are the result of spin-dependent electron optics arising due to the characteristics of the magnetic point contacts. These contacts introduce two key contributions to the graphene region in the immediate vicinity of the Co nanowire interface: (i) charge transfer doping and (ii) exchange splitting of the graphene band structure, $E_{\text{ex}} = E_{\text{Fi}\uparrow} - E_{\text{Fi}\downarrow}$, due to the magnetic proximity effect [42,48].

Charge transfer from cobalt to graphene results in a small n -doped region around the graphene-metal interface [9,40,49]. In analogy to light passing a boundary between two media with different refractive indices, ballistic electrons moving across the boundary between regions of different carrier density experience refraction—a consequence of the charge carriers having different Fermi energy on either side of the potential barrier [50]. Conservation of momentum across the barrier gives us an electrostatic equivalent of Snell's law [29]:

$$E_{\text{Fi}} \sin(\theta_i) = E_{\text{Fo}} \sin(\theta_o), \quad (2)$$

where E_{Fi} (E_{Fo}) and θ_i (θ_o) are the incoming (outgoing) Fermi energy and angle, respectively. Thus, the barrier can act as a convex ($|E_{\text{Fo}}| > |E_{\text{Fi}}|$) or concave ($|E_{\text{Fo}}| < |E_{\text{Fi}}|$) lens, collimating or dispersing the electron trajectories in these respective cases. This effect explains the greater magnitude and sharper focusing peaks for hole transport compared to electron transport [Fig. 2(a)], with the implication that the beam of ballistic electrons is more tightly collimated in the case of p doping in the channel, bringing the distribution of spin-resolved trajectories closer together. The situation is reversed for n doping in the channel, where the paths taken by spin-up and spin-down electrons are forced further apart, giving rise to a broader peak. This interpretation is in agreement with our observation of n doping at the Co nanowire-graphene junctions, which results in $|E_{\text{Fo}}| > |E_{\text{Fi}}|$ (convex lens) for p doping in the channel and $|E_{\text{Fo}}| < |E_{\text{Fi}}|$ (concave lens) for the case of n doping in the channel (see Supplemental Fig. S5 [38]). Our hypothesis is further supported by a simple electron-optics model based on Eq. (2) [see Fig. 3(c) and Supplemental Material Sec. SIV [38]].

The potential barrier surrounding the contact is spin dependent as a result of the exchange split graphene band structure adjacent to the Co nanowire interface. This splitting leads to inequivalent refraction conditions for electrons with different spin polarization, i.e., $E_{\text{Fi}\uparrow} \neq E_{\text{Fi}\downarrow}$ [Fig. 3(c)]. We estimate the exchange splitting to be

$E_{\text{ex}} \approx 10$ meV (Supplemental Material Sec. SIII [38]), a value consistent with those extracted for graphene-cobalt junctions [42]. These spin-dependent potential barriers, which are sensitive to the value of exchange splitting and charge transfer, alter the focusing conditions in a spin-dependent manner. The latter can either add to or subtract from the spin polarization achieved in the absence of magnetic focusing, providing a significant enhancement of the spin signal [Fig. 3(b)].

The inversions in ΔR_{NL} are reminiscent of precession curves, but the nature of these modulations cannot be explained by this effect due to a mismatch in the relevant timescales; the time of flight during the focusing experiment is calculated as $\tau_f = 0.51 \pm 0.01$ ps, while the precession period is at least an order of magnitude greater throughout the magnetic field range explored (Supplemental Material Sec. SII [38]). The inversions can be explained only by the unique ballistic transport regime of the TMF experiment, including modulations of the spin-dependent focusing conditions as the magnetic field is tuned, as well as the potential role of diffusive scattering at the edge of the graphene channel.

Further to modulating focusing conditions by varying the perpendicular magnetic field, it is desirable to have electrostatic control over the focused spin signal. Figure 3(d) displays ΔR_{NL} extracted as a function of carrier density in the channel, at three fixed values of $B_{\perp}$. We evaluate ΔR_{NL} at fixed magnetic field values that span the $j = -1$ and $j = -2$ peaks as carrier density is varied: $B_{\perp} = -120$ mT and $B_{\perp} = -230$ mT, respectively [Fig. 2(a)]. We also pick an intermediate value, $B_{\perp} = -170$ mT, which

encompasses a transition between the $j = -2$ and $j = -1$ peaks. In all cases, the spin signal is continuously tunable within $R_{\text{NL}} \approx \pm 2.5\Omega$. For the signal evaluated at $j = -2$ ($B_{\perp} = -230$ mT), a full reversal of the spin signal polarity is observed over a carrier density range of $\Delta n \approx 0.5 \times 10^{12} \text{ cm}^{-2}$ (corresponding to $\Delta V_{\text{BG}} \approx 8$ V). A similar degree of modulation is observed for the transition between two modes ($B_{\perp} = -170$ mT), but the oscillations are approximately 180° out of phase. The oscillatory nature of these signals can be described by the function $R_{\text{NL}} = A \sin(2\pi j)$, where j is given by Eq. (1) and A is the amplitude. We fit this expression to the data shown in Fig. 3(d), with the resulting line shapes correctly describing the experimental data. The period obtained from this model is consistent across the three datasets and yields an average effective length $L = 2.1 \pm 0.1 \mu\text{m}$, in close agreement with the nominal contact separation ($L = 1.9 \mu\text{m}$). Hence, spin-dependent electron optics enables electrostatic control over both the amplitude and polarity of the focused spin signal, while the perpendicular magnetic field strength controls the phase of oscillations.

In a practical implementation of a voltage controlled spin transistor, carrier density could be modulated by a top gate, separated from the graphene by approximately 10 nm of *h*-BN, rather than the bottom gate approach (290 nm of SiO_2) used here. For a device of the same quality, using a top gate, a simple capacitor calculation (see Sec. V) implies that the voltage required to induce the same variation in carrier density would be $\Delta V \approx 240$ mV, evidencing the potential for integrating such a device in low-power hybrid spintronic-CMOS microelectronics [51].

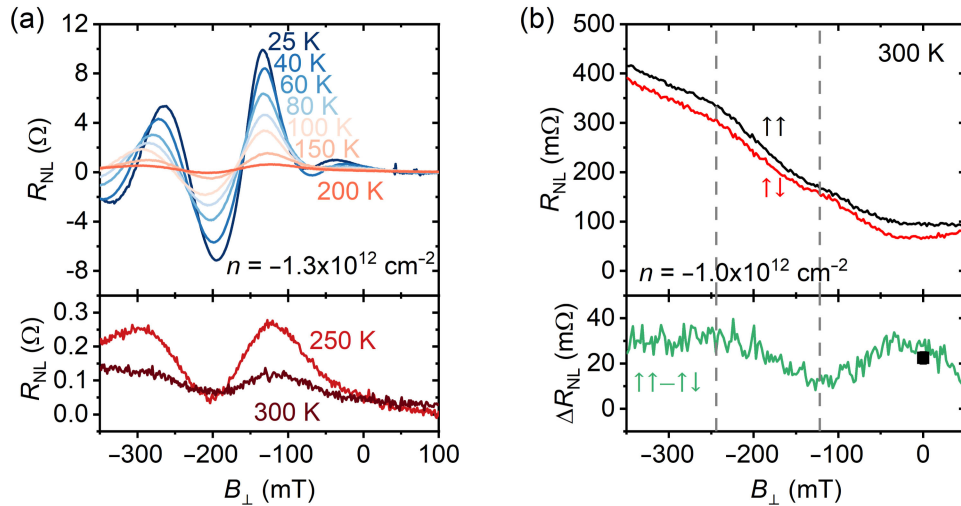


FIG. 4. (a) Temperature dependence of TMF. Individual TMF traces taken at a fixed carrier density ($n = -1.32 \times 10^{12} \text{ cm}^{-2}$) for a range of temperatures: 25–200 K (top) and 250–300 K (bottom). TMF signals are split into two panels for presentation purposes—different R_{NL} scales are necessary due to the substantial difference in peak amplitudes. (b) Spin-dependent TMF at room temperature—the parallel (black, $\uparrow\uparrow$) and antiparallel (red, $\uparrow\downarrow$) signals are measured, from which the spin signal is calculated (green, bottom panel). The black square in the bottom panel represents the magnitude of spin signal extracted from a room-temperature spin valve, at the same carrier density.

B. Temperature dependence

Figure 4(a) shows the evolution of the focusing signal as temperature is increased from 25 K up to room temperature. The TMF peaks are present throughout the measured temperature range but reduce in magnitude with increasing temperature. Phonon scattering is more significant at room temperature, reducing the electronic mean free path, which is estimated to be $\ell_{\text{mfp}} = 669 \pm 2$ nm. This value is approximately $3\times$ less than the contact separation, $L = 1.9$ μm , implying that injected charge carriers undergo an average of two scattering events between the injector and detector. Hence, transport at room temperature can be considered only quasiballistic.

We repeat the spin-dependent focusing experiment at room temperature, the result of which is seen in Fig. 4(b). At $B_{\perp} = 0$ mT, there is a clear distinction between focusing signals measured in the parallel and antiparallel configurations, consistent with the spin-valve measurement. At the first-order focusing condition, visible at $B_{\perp} = -115$ mT, there is a significant modulation of the spin signal, which becomes more negative. Although the TMF spectra peaks and the corresponding spin signal modulation are broadened compared to the low-temperature results, consistent with the quasiballistic transport regime and thermal broadening, the modulation in ΔR_{NL} resembles that observed at low temperature, indicating that spin-dependent magnetic focusing phenomena are still accessible at room temperature.

IV. CONCLUSIONS

In summary, we use TMF to demonstrate fully ballistic spin transport in a high-quality graphene spin valve. At low temperature (25 K), the magnitudes and positions of observed TMF peaks are sensitive to the magnetic configuration of the contacts, implying spin-polarized ballistic transport. Spin-dependent electron optics occurring at potential barriers surrounding the contacts give rise to an enhancement in spin signal, relative to the standard spin-valve response in the absence of focusing. At modest magnetic field strengths ($B_{\perp} \approx 100$ mT), the enhancement reaches up to approximately $5\times$ for electron doping in the channel, while hole-doped peaks show an inverted spin signal up to approximately $20\times$ greater in magnitude than that extracted from standard nonlocal spin valve measurements. Furthermore, we evidence electrostatic control over the magnitude and polarity of the ballistic spin signal via changing a back gate voltage at fixed perpendicular magnetic field strength. Finally, we show the survival of a quasiballistic signal at room temperature. These results constitute the first experimental demonstration of fully ballistic spin transport in a graphene spin valve and testify to the utility of TMF as a technique for enhancing and manipulating spin signals, with potential applications to spintronic and spin-based quantum devices.

V. METHODS

A. Electrical measurements

Our measurements are performed in high vacuum in a continuous flow cryostat, through which we pass liquid helium to reach cryogenic temperatures (25 K or above). Stable temperature is maintained by a proportional-integral-derivative feedback loop via an Oxford Instruments Mercury iTC, connected to a heater-thermometer pair in the cryostat. For electrical measurements, we use a SR830 lock-in amplifier, set to low frequency ($f = 5\text{--}20$ Hz), as a voltage source to drive an excitation current to the device ($I_{\text{AC}} = 1\text{--}5$ μA), via a series limiting resistance ($R = 1$ M Ω). We measure the resulting ac voltage across different sections of the device via the lock-in input channels; all signals returning to the lock-in are subject to a $10\times$ preamplifier, which increases input impedance from 10 to 100 M Ω .

B. Carrier density control via back gate voltage

Carrier density n in the graphene channel is varied by application of a dc voltage to the p -doped silicon substrate. We estimate carrier density by treating the graphene channel and p -doped Si back gate as the parallel plates of a capacitor, with the SiO_2 layer and bottom h -BN flake acting as two separate dielectric layers, connected in series. This gives us a way to convert between the back gate voltage V_{BG} and carrier density via the equation

$$n(V_{\text{BG}}) = C \frac{(V_{\text{BG}} - V_{\text{D}})}{e}, \quad (3)$$

where V_{D} is the gate voltage at which the Dirac point occurs ($V_{\text{D}} = -0.42$ V for this device) and e is the electron charge. C is the gate capacitance per unit area, calculated as

$$C_{\text{g}} = \frac{\epsilon_0 \epsilon_{\text{h}} \epsilon_{\text{Si}}}{(\epsilon_{\text{h}} d_{\text{Si}} + \epsilon_{\text{Si}} d_{\text{h}})}, \quad (4)$$

where $\epsilon_0 = 8.85 \times 10^{-12}$ Fm $^{-1}$ is the permittivity of free space and $\epsilon_{\text{h}} = 3.8$ and $\epsilon_{\text{Si}} = 3.9$ are the respective relative permittivities of the h -BN and SiO_2 layers (ϵ_{h} for bulk h -BN [8]), while $d_{\text{h}} = 20$ nm and $d_{\text{Si}} = 290$ nm are their respective thicknesses (d_{h} evaluated by AFM).

C. Electrostatic manipulation with top gate

Using Eqs. (3) and (4), along with the parameter values laid out in the above methods, we can estimate the change in voltage required to induce a change in carrier density of $\Delta n = 0.5 \times 10^{12}$ cm $^{-2}$, in the case of a top gate, separated from the graphene channel by 10 nm of h -BN. For such a gate architecture, the h -BN provides the only dielectric layer, so Eq. (4) reduces to $C_{\text{g}} = \epsilon_0 \epsilon_{\text{h}} / d_{\text{h}}$, giving a value of $C_{\text{g}} \approx 3.36$ mF for $d_{\text{h}} = 10$ nm. Rearranging Eq. (3) to give $\Delta V = \Delta n e / C_{\text{g}}$ and inserting the above values results in a change in gate voltage of $\Delta V \approx 0.24$ V.

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There are no competing interests to declare.

DATA AVAILABILITY

The data that support the findings of this article are openly available [52].

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