Spin Injection in a Ballistic Two-Dimensional Electron Gas

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We explore electrically injected, spin polarized transport in a ballistic two-dimensional electron gas. We augment the Büttiker-Landauer picture with a simple, but realistic model for spin-selective contacts to describe multimode reservoir-to-reservoir transport of ballistic spin 1/2 particles. Clear and unambiguous signatures of spin transport are established in this regime, for the simplest measurement configuration that demonstrates them directly. These new effects originate from spin precession of ballistic carriers; they exhibit strong dependence upon device geometry and vanish in the diffusive limit. Our results have important implications for prospective “spin transistor” devices.

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The concept of a spin transistor, first proposed almost ten years ago [1], has attracted widespread interest [2] but its experimental realization remains elusive [3]. It is based upon electrical injection of spin polarized carriers from a ferromagnetic conductor into an electron gas within a semiconductor. Electrons propagating in the interfacial electric field confining them to the device channel experience an effective magnetic field that induces spin precession; this is called the Rashba effect [4]. The rate of spin precession should be tunable through an external gate voltage, which will add to the confinement potential [5]. With ferromagnetic source and drain contacts acting as spin polarizer and analyzer, an electron device analogous to an electro-optic modulator [6] is envisaged.

Experiments to date demonstrate that the spin transistor geometry [Fig. 1(a)], in general, leads to strong Hall phenomena that are unrelated to true spin transport [7]. They arise due to the requisite proximity of miniature magnets and the low density (and, hence, high Hall coefficient) electron gas. Since these Hall phenomena depend directly upon the magnetization state of these magnetic contacts, they often closely mimic the signals expected from spin transport experiments—especially those in which the relative magnetic orientation of the “spin polarizer” and “analyzer” contacts is varied. However, in early experiments on (diffusive) spin injection in metals [8], spin precession phenomena provided an alternate and crucial experimental proof of spin transport. In this Letter, we establish for the first time the analog of these precessional effects will constitute a definitive experimental demonstration of electrical spin injection in semiconductor systems.

The spin transresistance, \( R_S \) [Fig. 1(a)] provides the most direct demonstration of spin transport [9]. This nonlocal transport coefficient is free of obfuscating background signals unrelated to spin injection. In the diffusive limit, if current contact \( F_1 \) is replaced with one that is unpolarized, no voltage will appear between the analyzer contact \( F_2 \) and a suitably defined ground reference \( R \). In this case these voltage contacts, being well outside the net current path, remain at equipotential. With a polarized current contact \( F_1 \), injected magnetization can lead to steady-state spin accumulation that persists over the entire length of the channel if \( \delta_S > L \). This induces disequilibrium between the local spin-resolved electrochemical potentials of \( F_2 \) and \( R \), and yields a finite \( R_S \). Here \( \delta_S = \sqrt{\ell_0 \ell_S / 2} \) is the spin diffusion length, \( \ell_S = v_F \tau_S \) and \( \ell_0 = v_F \tau_0 \) are the spin and momentum mean free paths, \( \tau_S \) and \( \tau_0 \) are the effective spin and momentum relaxation times, and \( v_F \) is the Fermi velocity.

In the ballistic regime, however, it is not appropriate to speak of spin accumulation since a local chemical potential cannot be meaningfully defined within the channel. Accordingly, our description of the ballistic spin transresistance is based upon Büttiker’s picture for mesoscopic transport within a multiprobe conductor [10]. Here we augment this with a new model describing spin-selective contacts. Our procedure is as follows: (a) We first develop a simple description of spin-selective contacts, based upon careful consideration of the ferromagnetic/semiconductor (F/S) contacts in (our) real devices [11]. (b) We construct an 8-reservoir model, after Büttiker, to describe the spin injection experiment. (c) Boundary conditions are used to constrain the spin-resolved currents and chemical potentials. These lead to a simpler 4-reservoir problem for the spin transresistance, \( R_S \), in terms of reservoir-to-reservoir, spin-resolved transmission probabilities, \( T_{ij}^{\alpha \beta} \), of the 2DEG forming the device conduction channel. Here the indices \( i, j \) specify the reservoirs themselves, and \( \alpha, \beta \) their constituent spin bands. (d) The requisite \( T_{ij}^{\alpha \beta} \) are then calculated semiclassically, using a modified Monte Carlo numerical technique (described below). We follow the electrons’ ballistic trajectories and the phase of their spin wavefunctions as they pass through the device, while ignoring the phase of their spatial wavefunctions. For unpolarized ballistic systems, this semiclassical approach has proven remarkably consistent with experimental data at \( T \sim 4\text{K} \), where the electron phase coherence length is smaller than typical dimensions of nanoscale devices [12].
Measurement of $R_S$ involves four terminals [Fig. 1(a)], two that are spin-selective, $\mathbf{F}_1$ and $\mathbf{F}_2$, and two that are conventional, i.e. momentum- and spin-relaxing, $\mathbf{L}$ and $\mathbf{R}$. As depicted in Figures 1(d, e, f), the full problem separates into three sub-components. Fig. 1(d) represents the spin-up and spin-down currents $(I_{L\uparrow}, I_{L\downarrow})$ that flow between $\mathbf{F}_1$, $1\uparrow$, $1\downarrow$, and $\mathbf{L}$. A Sharvin resistance $R_{sh} = (h/2e^2)(k_F w)/\pi = (h/2e^2)N_{ch}$, arises between $1\uparrow$, $1\downarrow$ and the multichannel conductors connecting them to $\mathbf{L}$. Under conditions of current flow this yields the spin-resolved electrochemical potential differences $\mu_{\uparrow} - \mu_{L} = 2eR_{sh}I_{L\uparrow}$ and $\mu_{\downarrow} - \mu_{L} = 2eR_{sh}I_{L\downarrow}$. Here, the factors of 2 arise because transport is spin resolved; $k_F$, $w$, and $N_{ch}$ are the Fermi wave vector, channel width, and number of occupied modes within the 2DEG device channel, respectively. Similarly, at the rightmost side of Fig. 1(f), current flow between the reservoirs $2\uparrow$, $2\downarrow$ and $\mathbf{R}$ establishes the electrochemical potential differences $\mu_{2\uparrow} - \mu_{R} = 2eR_{sh}I_{R\uparrow}$ and $\mu_{2\downarrow} - \mu_{R} = 2eR_{sh}I_{R\downarrow}$. Also, $\mu_{2\uparrow} = \mu_{2\downarrow}$ since no current flows between these reservoirs. Note that all $I$’s here represent net currents (forward minus reverse contributions). In our model, the following sum rules hold: $I = I_{L\uparrow} + I_{L\downarrow}$, $I = I_{R\uparrow} + I_{R\downarrow}$, and $I_{1\uparrow} + I_{1\downarrow} = I_{2\uparrow} + I_{2\downarrow} = 0$. As the reservoirs in Fig. 1(f) are voltage contacts, net current is conserved separately for each spin band, $I_{R\uparrow} + I_{2\uparrow} = I_{R\downarrow} + I_{2\downarrow} = 0$. These expressions can be manipulated to yield

$$
\begin{align*}
\begin{pmatrix} 
\mu_{1\uparrow} \\
\mu_{1\downarrow} \\
\mu_{2\uparrow} \\
\mu_{2\downarrow}
\end{pmatrix} &= 
\begin{pmatrix} 
\mu_{L} + 2eR_{sh}(I - I_{1\uparrow}) \\
\mu_{L} + 2eR_{sh}I_{1\uparrow} \\
\mu_{R} - 2eR_{sh}I_{2\uparrow} \\
\mu_{R} + 2eR_{sh}I_{2\uparrow}
\end{pmatrix}.
\end{align*}
$$

Given these relations, calculation of $R_S$ reduces to a 4-terminal problem that solely involves the four spin-resolved reservoirs: $1\uparrow$, $1\downarrow$, $2\uparrow$, and $2\downarrow$ and the 2DEG device channel that connects them [Fig. 1(c)]. Modifying Büttiker’s formula to account for the spin-resolved channels, the 4-terminal linear response at zero temperature becomes

$$
I_{\alpha\beta} = \frac{e}{h} \left[ (N_{ch} - R_{\alpha\alpha})\mu_{\alpha} - T_{ij}^{\alpha\beta} \mu_{j\beta} \right] = \frac{e}{h} U_{ij}^{\alpha\beta} \mu_{j\beta}. 
$$

Transport within the ballistic multinode 2DEG conductor is fully represented by the transmission and reflection coefficients, $T_{ij}^{\alpha\beta}$ and $R_{ij}^{\alpha\alpha}$. These describe carriers incident from the lead $i$ with spin polarization $\alpha$, that are transmitted into lead $j$ with final spin state $\beta$ ; and carriers incident from $\alpha$ that are reflected back into same lead and spin channel, respectively. The coefficients $U_{ij}$ in Eq. 1 satisfy the sum rule $\sum_{\alpha} U_{ij}^{\alpha\beta} = \sum_{\beta} U_{ij}^{\alpha\beta} = 0$ ensuring the current sum rules of Eq. 2 and that all currents vanish when the $\mu_{ij}$ are equal.
and 0.19, appropriate for a typical metal and for InAs, re-

figurations and two channel lengths, \( L/w \) (for identical reasons). For parallel alignment of po-
nel, of length \( L \), and width \( w \), which we denote by the superscript \( \uparrow \uparrow \), these steps yield

\[
\begin{pmatrix}
I_1^\uparrow \\
I_1^\downarrow \\
I_2^\uparrow \\
I_2^\downarrow
\end{pmatrix}
= S
\begin{pmatrix}
\hat{\mu}_L + I \\
\hat{\mu}_L \\
\hat{\mu}_R \\
\hat{\mu}_R
\end{pmatrix}.
\]

For each path segment traversed by the electron be-
tween boundary reflections, the phase of its spin wave-
function evolves continuously via the local Larmor fre-
quency \( \omega_L = g^* e B / 2 m \). Here \( g^* \) is the effective elec-
tron \( g \)-factor, \( B \) the local magnetic field, and \( m \) the free
lectron mass. Total precession is accumulated for each
complete trajectory, which is the sum of these segments.

For each segment the electron’s spin precession is calcu-
lated analytically \[12\] and incorporated into the Monte
Carlo procedure.

In Fig. 2(a) we display \( R_S^{\uparrow \uparrow} \) as a function of perpen-
dicular magnetic field strength. The prominent and
striking new feature is that \( R_S \) is oscillatory, a ballis-
tic phenomenon not found in the diffusive regime. In
Fig. 2(b,c) we display \( R_S^{\uparrow \uparrow} \) calculated for three orienta-
tions of the external field—two that are in-plane and
the perpendicular case, displayed again for comparison.
In all three cases the \( \bf{F}_1 \), \( \bf{F}_2 \) magnetizations are parallel and \( \uparrow \uparrow \)-oriented.

When the external field is along \( \hat{y} \), the injected carriers
remain in spin eigenstates and do not precess. In this
situation \( R_S^{\uparrow \uparrow} \) is a positive constant [Fig. 3(b,c)]. However,
with an \( \hat{x} \)-oriented field precession is maximal, and \( R_S \)
oscillates. Since orbital effects are absent for an in-plane
field, the oscillations in this case arise purely from spin
precession and the oscillation period, \( \Delta B \), is determined
by the condition \( 2 \pi n = \omega_L t_{TR} \), i.e., \( \Delta B = h / (g^* \mu_B t_{TR}) \).
Here \( t_{TR} = S / v_F \) is a typical transit time from \( 1 \to 2 \), and
\( \mu_B \) is the electronic Bohr magneton. \( \Delta B \) is thus
inversely proportional to \( S \), a typical path length aver-
ged over the injection distribution function. The decay
of \( R_S \) occurs on a field scale where \( \omega_L \delta t_{TR} \sim \pi \); i.e. for
\( B = h \pi / (g^* \mu_B t_{TR}) \) beyond which precession amongst
the different contributing trajectories tends to get out of
step. Here \( \delta t_{TR} = (\hat{t}_{FR}^2 - \langle \hat{t}_{FR} \rangle^2)^{1/2} \) is the variance
in path lengths traversed while propagating from \( 1 \to 2 \).

Perpendicular field \( (B_{ext}|\hat{z}|) \) is special—it induces both
spin and orbital effects. (The characteristic field scale for
the latter is \( B_{ext} = p_F / e w \), at which the cyclotron ra-
dus, \( r_c = v_F / e w \), equals the channel width, \( w \).) The
frequency ratio, \( \omega_L / \omega_c = (g^* / 2)(m^*/m) \) describes
the relative importance of orbital and spin transport phe-
nomena. Here, \( p_F \) is the Fermi momentum, \( \omega_c = e B / m^* \)
the cyclotron frequency, and \( m^* \) the effective mass. For
InAs (\( m^* = 0.025 \), \( g^* = 15 \)) this ratio is \( \sim 0.19 \), for In-
GaAs \( \sim 0.1 \), whereas it is roughly 1.0 for most metals.
In the latter spin and orbital effects have similar periodicity
so disentangling them is difficult [Fig. 3(a)].

As mentioned, electrons confined within an InAs het-
terostructure are subject to an internal Rashba field, present
even for zero applied magnetic field. This can be modeled by a Hamiltonian \[1\], \( H_R = \alpha_R (\sigma \times \hat{z}) \). Comparing \( H_R \) to the Zeeman term we write the ef-
ective Rashba field as \( B_R = 2 \alpha_R \sigma k \times \hat{z} / (g^* \mu_B) \). Here
\( \alpha_R = \Delta_R / 2 k_F \) is the spin-orbit coupling parameter \( \Delta_R \)
is the Rashba splitting \[13\], and \( k \) and \( k_F \) are the elec-
tron and Fermi wave vectors, respectively. Using data

\[\text{FIG. 2. Ballistic spin transresistance in an external field}
\text{normalized to } B_0 = p_F / e w, \text{ at which the cyclotron radius}
\text{equals the channel width. (a) For a channel with } L/w = 15 \text{ in}
\text{perpendicular field, we plot two traces representing } \omega_L / \omega_c = 1
\text{ and } 0.19, \text{ appropriate for a typical metal and for InAs, re-
}\text{spectively. (b,c) Spin transresistance for three different con-
}\text{figurations and two channel lengths, } L/w = 3 \text{ and } 15. \text{ Here}
\text{ } \omega_L / \omega_c = 0.19 \text{ (InAs).}
\text{Simplification of Eq. 1 and 2 yields}
\begin{align*}
\begin{pmatrix}
I_1^\uparrow \\
I_1^\downarrow \\
I_2^\uparrow \\
I_2^\downarrow
\end{pmatrix}
&= S
\begin{pmatrix}
\hat{\mu}_L + I \\
\hat{\mu}_L \\
\hat{\mu}_R \\
\hat{\mu}_R
\end{pmatrix}.
\end{align*}
\]

\[\text{where } \hat{\mu}_L,R = \mu_{L,R} / 2 e R_{sh} \text{ and } S \equiv (1 + U)^{-1} U. \text{ The elements of } S \text{ satisfy the same sum rules that constrain } U \text{ (for identical reasons). For parallel alignment of pol-
arizer and analyzer, } \bf{F}_1 \text{ and } \bf{F}_2, \text{ which we denote by the superscript } \uparrow \uparrow, \text{ these steps yield}
\begin{align*}
R_S^{\uparrow \uparrow} &= -2 \frac{S_{31} S_{42} - S_{32} S_{41}}{S_{31} + S_{32} + S_{41} + S_{42}} R_{sh},
\end{align*}
\]

For antiparallel alignment, only the sign changes: \( R_S^{\downarrow \downarrow} = -R_S^{\uparrow \uparrow} \).

We obtain the requisite elements of \( S \) numerically,
extending the semiclassical billiard model \[10\] to al-
low tracking of an electron’s spin wavefunction along
ballistic trajectories linking the spin-resolved reservoirs
(1\uparrow,1\downarrow,2\uparrow,2\downarrow) at either end of the 2DEG device
cannel. We consider electrons confined within a hard-wall chan-
nel, of length \( L \) and width \( w \). The \( T_{ij}^{\alpha \beta} \) are calculated by
injecting and following a large number of electron trajec-
tories (typically \( > 10^4 \)) propagating at \( v_F \) \[10\].
The result of current division. For finite $\hat{\omega}_R$, precessional effects are maximal.

For large $\hat{\omega}_R$ the oscillations decay quickly initially, but exceedingly slowly thereafter. No spin precession occurs for $\hat{\omega}_R = 0$, hence, $t = 1$ yielding $R_S^{(\uparrow\downarrow)} = R_S^{(-\rightarrow)} = R_{sh}/3$, a simple result of current division. For finite $\hat{\omega}_R$, $R_S$ displays strong dependence upon the orientation of the magnetizations $\mathbf{M}$ (of $\mathbf{F}_1, \mathbf{F}_2$, assumed parallel), in relation to the device channel’s principal axis ($\hat{x}$). For $\mathbf{M}|\hat{x}$ (parallel to the channel), precessional effects are maximal. With increasing Rashba field, the variance in contributing path lengths causes the oscillations in $R_S$ to decay, as described previously for the case of finite external field. Here, however, the contributions from short paths (direct propagation between the DSPR’s involving few or no boundary reflections) continue to add coherently for large $\hat{\omega}_R$, resulting very slow decay. For $\mathbf{M}|\hat{y}$ most of the injected carriers experience a Rashba field nearly aligned with their spin. At intermediate Rashba field these yield small oscillations that center about a finite value of $R_S$. The other carriers make a contribution to $R_S$ at small $\hat{\omega}_R$ but this becomes incoherent and thus quickly decays for large $\hat{\omega}_R$.

The original idea of the spin transistor involved use of an external gate potential, acting in concert with the intrinsic confinement potential, to control of the spin precession rate. We note, however, that gate tuning of $\alpha_{so}$ for electrons has been experimentally demonstrated in relatively few narrow gap semiconductor heterostructures. Two such systems are InGaAs/InP and InGaAs/InAlAs. In the latter, tuning over about a 30% range has been reported. In Fig. 3 we show how this range of tunability translates into a direct modulation of $R_S$, for three device widths. Our calculations clearly illustrate that the “conventional” spin transistor configuration, $\mathbf{M}|\hat{x}$, (which is most easily fabricated) is not optimal—even for a very short channel ($L \sim \ell_S$). We find that tunability is maximized for $\mathbf{M}|\hat{x}$.

The spin transistor was originally envisaged as a one-dimensional device, with only a single populated transverse subband. Realizable devices in the near term will more likely be two-dimensional or, perhaps, quasi-one-dimensional channels. Their increased phase space for scattering can lead to quick suppression of $R_S$, especially in the presence of moderate scattering. Hence it appears that an extremely narrow channel is a basic requirement for a spin transistor.

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Among these are partial spin polarization and additional unpolarized bands at $E_F$ in $F_1$, $F_2$; magnetic disorder and spin scattering at the F/S interface; momentum and spin scattering in the 2DEG; and thermal smearing.

We assume that the $T_{\alpha\beta}^{\alpha\beta}$ are essentially energy-independent on the scale of the spin splitting. The latter is small compared to the Fermi energy in narrow gap, low density semiconductor heterostructures.

We find that introduction of diffuse boundary or bulk scattering to our calculations tends to wash out both the asymptotic effects and the high field oscillations. We also find that our results are insensitive to the precise formulation of the DSPR’s; more complex (8-reservoir) calculations that include junction scattering in these regions display similar physics.