Pumping spin with electrical fields

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Spin currents can be obtained through adiabatic pumping by means of electrical gating only. This is possible by making use of the tunability of the Rashba spin-orbit coupling in semiconductor heterostructures. We demonstrate the principles of this effect by considering a single-channel wire with a constriction. We also consider realistic structures, consisting of several open channels where subband spin-mixing and disorder are present, and we confirm our predictions. Two different ways to detect the spin-pumping effect, either using ferromagnetic leads or applying a magnetic field, are discussed.

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The investigation of spin-dependent transport and its application to steer electrical currents is at the foundation of Spintronics.\textsuperscript{[1]} Both of fundamental interest and of practical importance, the success in operating spin-based devices relies on the ability to produce and control spin currents. At present, techniques to obtain a spin current include injection from ferromagnets, Zeeman’s splitting due to magnetic fields, and optical excitations\textsuperscript{[2]}. Very recently some alternative proposals have been put forward. Mucciolo\textsuperscript{[3]} suggested to obtain spin currents based on the use of pumping of electrons through a chaotic dot in the presence of an in-plane magnetic field; Bratatis et al.\textsuperscript{[5]} proposed a spin battery relying on a ferromagnet with precessing magnetization.

Adiabatic charge pumping\textsuperscript{[6, 7, 8]} consists in the transport of charge obtained, at zero bias voltage, through the periodic modulation of some parameters (e.g. gate voltages) in the scattering region. If the time variation of the scattering matrix occurs on a long time scale compared to the transport time then the charge transferred per period does not depend on the detailed time evolution of the scattering matrix but only on geometrical properties of the pumping cycle\textsuperscript{[6]}. Numerous works (e.g. Refs.\textsuperscript{[6, 10, 11, 12, 13, 14]} and references therein) addressed different aspects of adiabatic pumping as, for example, the counting statistics of the pumped current, the generalization to multi-terminal geometries and the question of the phase coherence.

Adiabatic pumping of spin seems to be quite attractive as well, although little attention has been payed to it so far (see however Ref.\textsuperscript{[1]})\textsuperscript{[4]}. A combined implementation of adiabatic charge pumping with a spin filter will ensure that if charge transport occurs also spin is transferred. In this Letter we discuss the possibility of spin pumping without using ferromagnetic materials or external magnetic fields. This is indeed possible by making use of the tunability of the Rashba effect\textsuperscript{[15, 16, 17]}. A spin current is then produced by electrical gating only. Adiabatic pumping plays a crucial role in the present mechanism since there is no spin-polarized current if the same device is dc biased and no time-dependent transport is involved.

Electrons confined in a two-dimensional electron gas, realized in a semiconductor heterostructure with some asymmetry in growth direction $z$, are subject to the Rashba spin-orbit (SO) coupling whose Hamiltonian reads $H_{\text{so}} = \frac{\hbar k_{\text{so}}}{m} (\sigma_x p_y - \sigma_y p_x)$, where $m$ is the effective mass. It is important to notice that the strength of the SO coupling, denoted as $k_{\text{so}}$, can be tuned by changing the asymmetry of the quantum well via externally applied voltages, as shown in several experimental studies\textsuperscript{[18, 19, 20]}. The system we have in mind to produce a spin current is schematically depicted in Fig. 1. It consists of a quantum wire (parallel to the $x$-axis) of length $L$ with Rashba spin-orbit coupling, connected to two semi-infinite leads, where spin-orbit coupling is absent. At the interface between the wire and the left lead a constriction will give rise to a potential barrier denoted by $V_{\text{bar}}$. The Hamiltonian of the wire can be written as $H = H_{1D} + H_{\text{tras}} + H_{\text{mix}}$, with

$$H_{1D} = \frac{1}{2m} p_x^2 - \frac{\hbar k_{\text{so}}}{m} \sigma_y p_x,$$  

$$H_{\text{tras}} = \frac{1}{2m} p_y^2 + V_{\text{conf}}(y),$$

$$H_{\text{mix}} = \frac{\hbar k_{\text{so}}}{m} \sigma_x p_y,$$

where $H_{1D}$ describes the longitudinal motion along the wire, $H_{\text{tras}}$ is the transverse part of the Hamiltonian (it contains the transverse confining potential $V_{\text{conf}}$), and $H_{\text{mix}}$ the part of the SO coupling that is responsible for subband mixing\textsuperscript{[21, 22]}. If the spin precession length $l_{\text{so}} = \pi/k_{\text{so}}$ is much larger than the typical width of the wire then $H_{\text{mix}}$ can be neglected and a common spin quantization axis can be found (perpendicular to the wire and the heterostructure growth direction). In this limit the quasi-one-dimensional subband dispersion relations read: $\epsilon_{n,\sigma} = \frac{\hbar^2}{2m} (k_x - \sigma k_{\text{so}})^2 - \Delta_{\text{so}} + E_n$, where $E_n$ is the transverse energy, $\sigma = \pm$ is the quantum number of $\sigma_y$, and $\Delta_{\text{so}} = \hbar^2 k_{\text{so}}^2/2m$. In order to compute the
pumped charge and spin the scattering matrix should be determined. For the sake of simplicity and clarity we present the general idea considering that only a single subband is occupied both in the wire and leads. To avoid cluttering the notation and allow for simple analytical expressions, we further assume that no Fermi-velocity mismatch is present between the leads and the wire and that Δso is much smaller than the Fermi energy \[23\]. The analytical results are presented for a delta function potential \[\bar{V}_{\text{bar}} (x) = V \delta (x)\]. All these assumptions do not affect the basic principles of our proposal. Indeed we will show that when these hypothesis are relaxed, only small quantitative changes occur. Since in this idealized model there is no spin-mixing mechanism, we can treat the two spin species separately. The pumping cycle is obtained by varying in time the height of the barrier (we define \[\bar{V} = 2mV \bar{V}/h^2\], and the spin-orbit coupling \[k_{so}\]. In the following we discuss both the average current and the noise spectrum. The average spin-σ particle current is given by \[2\]

\[
I_{\sigma} = \frac{\omega}{2\pi^2} \int d\bar{V} \, dk_{so} \, D_{\sigma}(\bar{V}, k_{so}),
\]

where the integral is over the surface spanned during a cycle in parameter space, \(\omega\) is the frequency of the pumping fields, and

\[
D_{\sigma}(\bar{V}, k_{so}) = \text{Im} \left\{ \frac{\partial r_{\sigma}^*}{\partial V} \frac{\partial t_{\sigma}^*}{\partial k_{so}} + \frac{\partial t_{\sigma}}{\partial V} \frac{\partial r_{\sigma}}{\partial k_{so}} \right\}.
\]

(3)

with \(r_{\sigma}\) and \(t_{\sigma}\) being, respectively, the reflection and transmission coefficient for an electron with spin \(\sigma\). From the spin-σ currents Eq. (2), we can define a charge/spin current as \(I_{\text{charge/spin}} = I_+ \pm I_-\) (note that \(I_{\text{charge}}\) is expressed in units of electron charge).

While the average pumped currents depend only on the geometrical properties of the pumping cycle, the current noise depends on the full time-dependence of the pumping parameters. Since we are interested in current fluctuations around the average current we calculate the zero-frequency component of the noise spectrum

\[
S_{\sigma} = \frac{\omega}{2\pi} \int d\tau \int_{0}^{\infty} d\tau' \langle \Delta \hat{I}_{\sigma}(\tau) \Delta \hat{I}_{\sigma}(\tau') \rangle,
\]

(4)

where \(\Delta \hat{I}_{\sigma} = \hat{I}_{\sigma} - \langle I_{\sigma} \rangle\). In the case of weak pumping the knowledge of the average number of transmitted particles and of the zero-frequency noise characterizes the full counting statistics \[10\]. As there are no correlations between electrons with different spin indexes, the noise of the charge current and of the spin current is simply \(S_{\text{spin}} = S_{\text{charge}} = S_+ + S_-\). Several authors have studied noise in quantum pumps \[2\, 11\, 13\]. We make use of the formulation of Moskalets et al.

Once the scattering matrix is determined, Eq. (3) yields

\[
D_{\sigma}(\bar{V}, k_{so}) = \frac{\sigma}{(4k_F^2 + V^2)^2} \frac{4\bar{V}^2 L V}{(4k_F^2 + V^2)^2},
\]

(5)

where \(k_F\) is the Fermi wave-vector for the leads \[24\]. From Eq. (4) we immediately obtain that the pumped charge current is zero and the pumped spin current is \(I_{\text{spin}} = 2I_+\). For a sinusoidal pumping cycle: \(\bar{V} = V_0 + \Delta V \sin(\omega \tau)\) and \(k_{so} = k_{so, 0} + \Delta k_{so} \sin(\omega \tau - \phi)\), with \(\Delta V \ll V_0\) (weak-pumping limit) we can determine the explicit form of the average current

\[
I_{\text{spin}} = \frac{\omega}{2\pi} \sin(\phi) \Delta V \Delta k_{so} \frac{8k_F^2 LV_0}{(4k_F^2 + V_0^2)^2}
\]

(6)

and noise

\[
S_{\sigma} = \frac{\abs{\omega}}{\pi} \frac{2k_F^2 L^2 V_0^2}{(4k_F^2 + V_0^2)^2} \frac{\Delta V^2 + \Delta k_{so}^2 L^2 V_0^2}{\Delta V^2 + \Delta k_{so}^2 L^2 V_0^2}.
\]

(7)

For the particular pumping cycle chosen, and for vanishing temperature, the zero-frequency component of the spin-σ current noise does not depend on the spin-index and on the phase \(\phi\). The spectrum of Eq. (7) shows that the fluctuations introduced by the modulation of \(\bar{V}\) and \(k_{so}\) are uncorrelated. We can define a signal-to-noise ratio as \(\abs{I_{\text{spin}}}/S_{\text{spin}}\) that in the present case reads

\[
\frac{\abs{I_{\text{spin}}}}{S_{\text{spin}}} = \frac{2}{\hbar} \frac{\abs{\sin(\phi) V_0^2 \Delta V \Delta k_{so} L}}{\Delta V^2 + \Delta k_{so}^2 L^2 V_0^2}
\]

(8)

The signal-to-noise ratio Eq. (8) reaches its maximum at fixed \(\phi\) for \(\Delta V = \Delta k_{so} L V_0\).

In the simplest arrangement the spin-pumping effect can be detected if one of the two leads has been replaced by a half-metallic ferromagnet (i.e. only majority spins are present). If its magnetization lies in the plane of the wire and makes an angle \(\theta\) with the y-axis the spin state of the electrons in the ferromagnetic lead is \(|F\rangle = \cos \frac{\theta}{2} |+\rangle + i \sin \frac{\theta}{2} |-\rangle\) (\(|\pm\rangle\) are the eigenstates of
magnetic field (to lowest order in $\Delta k$ of magnetic field). Assuming that scattering matrix within the tight-binding model, using the Hamiltonian $H$ more realistic model, which includes several modes, subband mixing induced by Eq. (14), and the effect of the time modulation of $\Delta s_0$. We numerically calculate the scattering matrix within the tight-binding model, using a recursive Green’s function technique [27]. The tight-binding version of the Rasha SO coupling can be written

\[ H_{so} = -i\gamma_{so} \sum_{\sigma} \sum_{i,j} \left[ c_{i+1,j+1,\sigma}^\dagger c_{i,j,\sigma} - c_{i+1,j+1,\sigma'} c_{i,j,\sigma'} \right] + \text{h.c.} \]

where $c_{i,j,\sigma}^\dagger$ is the creation operator of an electron in site $(i,j)$ with spin $\sigma$ and $\gamma_{so}$ is the Rashba nearest-neighbour coupling. Note that $\gamma_{so}$ is related to the parameter $k_{so}$ through the relation: $\gamma_{so} = (\alpha_{so})/\gamma$, where $\gamma$ is the tight-binding hopping potential and $\alpha$ is the lattice constant.

In our simulations the wire is modeled as a 2D lattice with $W = 3$ sites in the transverse $y$-direction and $N = 50$ sites in the longitudinal $x$-direction. The wire is then attached to the two leads, in which $\gamma_{so} = 0$, through a hopping potential $\Gamma_0$ on the left-hand-side and $\Gamma_\sigma$ on the right-hand-side (in the following we set $\Gamma_\sigma = \gamma$). The Fermi energy is chosen so that three bands are occupied (from now on we express $\gamma_{so}$, $\Gamma_\sigma$, and $\Gamma_\sigma$ in units of $\gamma$).

Adiabatic pumping is obtained by performing a square cycle in the parameter space $(\gamma_{so}, \Gamma_\sigma)$, with $\Gamma_\sigma$ varying in the range $\Gamma_0 \pm \delta \Gamma/2$ and $\gamma_{so}$ varying between zero and $\gamma_{so}^{max}$. In Fig. 2 the average number of spins and charges transmitted in a cycle are plotted as functions of $\delta \Gamma$ for different values of $\gamma_{so}^{max}$ and for fixed $\Gamma_0$. As expected, both $\Gamma_\sigma$ and $\Gamma_{I_{charge}}$ are increasing functions of $\delta \Gamma$. It is remarkable that for $\gamma_{so} = 0.042$ and $\gamma_{so} = 0.125$, corresponding to typical values for Rashba splitting in semiconductor [24], $\Gamma_{I_{spin}}$ is about two orders of magnitude larger than $I_{charge}$ almost over the whole $\delta \Gamma$ range. Note that even for $\gamma_{so} = 0.25$, values which exceeds the maximum reported Rashba coupling strength [26], $\Gamma_{I_{spin}}$ is still much larger than $I_{charge}$. The pumped charge $I_{charge}$ remains much smaller compared to $\Gamma_{I_{spin}}$ as long as we are in the weak Rashba coupling regime [21], in which the inter-subband mixing due to Rashba coupling is negligible (in our case $\gamma_{so} < 0.38$). Since in our simulations $\gamma_{so} \approx 0.165$ corresponds to the maximum reported value for $k_{so}$, there is no need to go beyond the weak Rashba regime, at least for narrow wires. The inclusion of an additional constant on-site energy in the leads (modeling a difference in the Fermi velocity between the leads and the wire) does not introduce any new time dependence in the scattering matrix, and hence it does not hinder the principle on which spin pumping is based on.

We also considered the presence of disorder by adding to the tight-binding on-site energies in the Rashba region a random potential. We find that averaging over disorder realizations yields a suppression of the average $|I_{spin}|$, with respect to the clean case, but keeping $|I_{spin}| \gg |I_{charge}|$. In the quasi-ballistic regime [27], adiabatic spin pumping still takes place with no qualitative difference.

We finally reanalyze the spin pumping from a different perspective. To this end we start noticing that the Hamiltonian $H_{1D}$ [see Eq. (13)] in the basis of eigenstates of $\sigma_y$ can be recast in the following form $H_{1D} =$
This electric field leads to a spin-dependent potential produced through adiabatic pumping with no use of temperature. Furthermore, the gauge transformation to spin-dependent electric fields shows that the pumping effect [28].

In conclusion we have shown that spin currents can be produced through adiabatic pumping with no use of ferromagnets or magnetic fields. Only electrical gating and the tunability of the Rashba coupling are exploited. This effect is robust also when several propagating modes are present, even in the case of non-negligible subband mixing and disorder. In addition, two possible different ways to detect spin pumping have been discussed and the zero-frequency noise spectrum has been calculated.

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\[ \frac{1}{2m}(p_{x} - e \mathbf{A}_{x} \cdot \hat{x})^{2}, \]

where the spin-dependent vector potential is given by \( \mathbf{A}_{\sigma} = \frac{\hbar}{e} \mathbf{r}_{\sigma} \). As \( \nabla \times \mathbf{A}_{\sigma} = 0 \) this vector potential does not describe a magnetic field. But if \( k_{so} \) varies with time \( \tau \) it describes a spin-dependent electric field

\[ \mathbf{E}_{\sigma}(\tau) = -\partial_{\tau} \mathbf{A}_{\sigma}(\tau) = -\frac{\hbar}{e} \mathbf{r}_{\sigma} \cdot k_{\text{so}}(\tau) \hat{x}. \]

This electric field leads to a spin-dependent potential drop along the wire \( V_{\sigma} = \int_{0}^{\delta_{\Gamma}} \mathbf{E}_{\sigma} \cdot d\hat{x} = \frac{\hbar}{e} \partial_{\tau} k_{\text{so}} L \). Thus, we can consider the wire without SO coupling but with spin-\( \sigma \) electrons subject to the time-dependent potential \( V_{\sigma} \). An additional time-dependent barrier \( V_{\text{bar}} \) (as it was the case for the pumping cycle that led to Eq. (6)) leads to rectification of the oscillating potential \( V_{\sigma} \). Provided that the voltage \( V_{\sigma} \) is small enough so that linear transport theory applies, and that it changes on a time scale much larger than the time an electron needs to go through the scattering region the average spin-current reads [14]

\[ I_{\sigma} = \frac{\omega}{2\pi} \frac{e}{\hbar} \int_{0}^{2\pi} |t(\tau)|^{2} V_{\sigma}(\tau) d\tau, \tag{14} \]

where \( t \) is the transmission coefficient for \( k_{\text{so}} = 0 \). From Eq. (14) all the results obtained so far can be found. In this framework spin pumping appears as a rectification effect [28]. Furthermore, the gauge transformation to spin-dependent electric fields shows that the pumping mechanism survives when no phase coherence is present (although the description that relies on the scattering matrix ceases to be valid). The spin current will be limited by the spin-relaxation rate (which depends on temperature).

In conclusion we have shown that spin currents can be produced through adiabatic pumping with no use of ferromagnets or magnetic fields. Only electrical gating and the tunability of the Rashba coupling are exploited. This effect is robust also when several propagating modes are present, even in the case of non-negligible subband mixing and disorder. In addition, two possible different ways to detect spin pumping have been discussed and the zero-frequency noise spectrum has been calculated.

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[23] If \( k_{\text{so}} \) depends on time also \( \Delta_{\text{so}} \) does. The only effect of the time variation of \( \Delta_{\text{so}} \) is to induce some small charge pumping, but it does not affect spin-pumping.

[24] The reflection and transmission coefficients for the present model, which are needed to compute \( D_{\sigma} \), read

\[
t_{\sigma} = 2k_F e^{i\sigma k_{\text{so}} L}/(2k_F + i\bar{V}), \quad r_{\sigma}' = -i\bar{V}/(2k_F + i\bar{V}).
\]

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[27] In the simulations (not reported here) we take the on-site energy randomly distributed in the range \([-U/2, U/2]\) with \( U \) as large as \( 1/10 E_F \).

[28] Further predictions can be made using Eq. (14). If in a time \( T \), the SO coupling strength goes from zero to a value \( \Delta k_{\text{so}} \) and that for \( \tau > T \) it stays fixed at this value, during the interval \( T \) a net spin current flows (but no charge current), and the total spin transferred is simply given by

\[
\frac{2e}{h} \int_{-T}^{T} V_4(\tau) d\tau.
\]

If one of the leads is replaced by a ferromagnet or a transverse in-plane magnetic field is present, this particular variation of \( k_{\text{so}} \) will result in a pulse of charge current of duration \( T \).