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*.[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,[2] Zeeman's splitting due to magnetic fields, and optical excitations.[3] Very recently some alternative proposals have been put forward. Mucciolo *et al.*[4] 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; Brataas *et al.*[5] proposed a spin battery relying on a ferromagnet with precessing magnetization.

Adiabatic charge pumping[6–8] consists of 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.[6] Numerous works (e.g., Refs. 9–14 and references therein) addressed different aspects of adiabatic pumping as, for example, the counting statistics of the pumped current, the generalization to multiterminal 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. 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 paper 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.[15–17] 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 $H_{\rm so} = (\hbar k_{\rm so}/m)(\sigma_x p_y - \sigma_y p_x)$, where m is the effective mass. It is important to note that the strength of the SO coupling, denoted as $k_{\rm so}$, can be tuned by changing the asymmetry of the quantum well via externally applied voltages, as shown in several experimental studies.[18–20] 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 semiinfinite 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 $V_{\rm bar}$. The Hamiltonian of the wire can be written as $H = H_{\rm 1D} + H_{\rm tras} + H_{\rm mix}$, with

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

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

$$H_{\rm mix} = \frac{\hbar k_{\rm so}}{m} \sigma_x p_y , \tag{1c}$$

where $H_{\rm 1D}$ describes the longitudinal motion along the wire, $H_{\rm tras}$ is the transverse part of the Hamiltonian (it contains the transverse confining potential $V_{\rm conf}$), and $H_{\rm mix}$ the part of the SO coupling that is responsible for subband mixing.[21,22] If

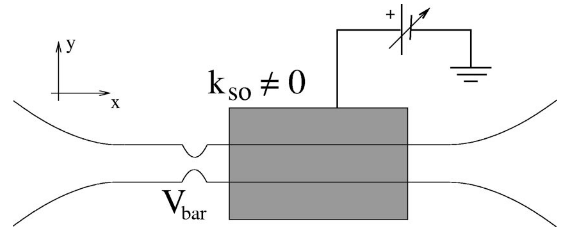

FIG. 1. Schematic setup for the adiabatic spin pump. It consists of a quantum wire with Rashba spin-orbit coupling $k_{\rm so}$ modulated by a gate (gray region) and controlled through a variable time-dependent voltage generator. A potential barrier $V_{\rm bar}$ (represented by the constriction) is present at the interface between the left lead and the wire.

the spin precession length $l_{\rm so}=\pi/k_{\rm so}$ is much larger than the typical width of the wire then $H_{\rm mix}$ 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: $\epsilon_{n,\sigma}=(\hbar^2/2m)(k_x-\sigma k_{\rm so})^2-\Delta_{\rm so}+E_n$, where E_n is the transverse energy, $\sigma=\pm$ is the quantum number of σ_y, and $\Delta_{\rm so}=\hbar^2k_{\rm so}^2/2m$. 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 $\Delta_{\rm so}$ is much smaller than the Fermi energy.[23] The analytical results are presented for a δ function potential $V_{\rm bar}(x)=V\delta(x)$. All these assumptions do not affect the basic principles of our proposal. Indeed we will show that when these hypotheses 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 $\overline{V}=2mV/\hbar^2$), and the spin-orbit coupling $k_{\rm so}$. In the following we discuss both the average current and the noise spectrum. The average spin-σ particle current is given by[6]

$$I_\sigma=\frac{\omega}{2\pi^2}\int d\overline{V}dk_{\rm so}D_\sigma(\overline{V},k_{\rm so}), \tag{2}$$

where the integral is over the surface spanned during a cycle in parameter space, ω is the frequency of the pumping fields, and

$$D_\sigma(\overline{V},k_{\rm so})=\mathrm{Im}\left\{\frac{\partial r_\sigma'^*}{\partial\overline{V}}\frac{\partial r_\sigma'}{\partial k_{\rm so}}+\frac{\partial t_\sigma^*}{\partial\overline{V}}\frac{\partial t_\sigma}{\partial k_{\rm so}}\right\} \tag{3}$$

with r_σ' and t_σ being the reflection and transmission coefficient for an electron with spin σ, respectively. From the spin-σ currents, Eq. (2), we can define a charge\spin current as $I_{\rm charge\backslash spin}=I_+\pm I_-$ (note that $I_{\rm charge}$ 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

$$S_\sigma=\frac{\omega}{2\pi}\int d\tau\int_0^{(2\pi/\omega)}d\tau'\langle\Delta\hat{I}_\sigma(\tau)\Delta\hat{I}_\sigma(\tau')\rangle, \tag{4}$$

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 $S_{\rm spin}=S_{\rm charge}=S_+ +S_-$. Several authors have studied noise in quantum pumps.[9,13,10] We make use of the formulation of Moskalets *et al.*

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

$$D_\sigma(\overline{V},k_{\rm so})=\sigma\frac{4k_{\rm F}^2L\overline{V}}{(4k_{\rm F}^2+\overline{V}^2)^2}, \tag{5}$$

where $k_{\rm F}$ is the Fermi wave vector in the leads.[24] From Eq. (5) we immediately obtain that the pumped charge current is zero and the pumped spin current is $I_{\rm spin}=2I_+$. For a sinusoidal pumping cycle, $\overline{V}=V_0+\Delta V\sin(\omega\tau)$ and $k_{\rm so}=k_{\rm so,0}+\Delta k_{\rm so}\sin(\omega\tau-\phi)$, with $\Delta V\ll V_0$ (weak-pumping limit) we can determine the explicit form of the average current

$$I_{\rm spin}=\frac{\omega}{2\pi}\sin(\phi)\Delta V\Delta k_{\rm so}\frac{8k_{\rm F}^2LV_0}{(4k_{\rm F}^2+V_0^2)^2} \tag{6}$$

and noise

$$S_\sigma=\frac{|\omega|}{\pi}\frac{2k_{\rm F}^2}{(4k_{\rm F}^2+V_0^2)^2}(\Delta V^2+\Delta k_{\rm so}^2L^2V_0^2). \tag{7}$$

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 ϕ. The spectrum of Eq. (7) shows that the fluctuations introduced by the modulation of $\overline{V}$ and $k_{\rm so}$ are uncorrelated. We can define a signal-to-noise ratio as $|I_{\rm spin}|/S_{\rm spin}$, which in the present case reads

$$\frac{|I_{\rm spin}|}{S_{\rm spin}}=\frac{2}{\hbar}\frac{|\sin(\phi)V_0^2\Delta V\Delta k_{\rm so}L|}{\Delta V^2+\Delta k_{\rm so}^2L^2V_0^2}. \tag{8}$$

The signal-to-noise ratio, Eq. (8), reaches its maximum at fixed ϕ for $\Delta V=\Delta k_{\rm so}LV_0$.

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 θ with the y axis, the spin state of the electrons in the ferromagnetic lead is $|F\rangle=\cos(\theta/2)|+\rangle+i\sin(\theta/2)|-\rangle$ ($|\pm\rangle$ are the eigenstates of σ_y). Furthermore, only to keep formulas compact, we assume that the Fermi velocity in the ferromagnetic lead is same as in the rest of the system. In this case the pumped charge and spin are given by

$$I_{\rm charge}^{\rm F}=\frac{I_{\rm spin}}{2}\cos\theta, \tag{9}$$

$$I_{\rm spin}^{\rm F}=\frac{I_{\rm spin}}{2}. \tag{10}$$

$I_{\rm spin}$ is the pumped spin current with normal leads. There are several remarks that should be made: (1) the spin current is independent of the magnetization direction; (2) the charge current is no more zero and it reaches its maximum when the magnetization is aligned with the spin-quantization axis in the wire; (3) the charge current can be reversed changing θ into $\theta+\pi$. The dependence of the pumped charge on the magnetization direction can be exploited to verify that the pumping mechanism is taking place.

As a second possibility, we consider a magnetic field in the y direction, which introduces only a Zeeman term in the Hamiltonian $H_B=(\hbar/2)\Omega_{\mathrm{B}}\sigma_y$. The effect of the Zeeman field is to modify the Fermi velocities for the two spin species. We can take this effect into account simply by replacing k_{F} in Eq. (5) with $k_{\mathrm{F},\sigma}=k_{\mathrm{F}}-\sigma\Delta k_{\mathrm{F}}$, where k_{F} is the Fermi wave-vector in the leads in the absence of magnetic field. Assuming that $|\Delta k_{\mathrm{F}}|\ll k_{\mathrm{F}}$ we can write $\Delta k_{\mathrm{F}}=(\Omega_{\mathrm{B}}/2)\times(m/\hbar k_{\mathrm{F}})$. Finding for D_σ in the presence of the magnetic field (to lowest order in $\Delta k_{\mathrm{F}}/k_{\mathrm{F}}$)

$$D^{\mathrm{B}}_\sigma(\overline{V},k_{\mathrm{so}})=D_\sigma(\overline{V},k_{\mathrm{so}})-4L\overline{V}\frac{4k_{\mathrm{F}}^2-\overline{V}^2}{(4k_{\mathrm{F}}^2+\overline{V}^2)^3}\,\frac{m\Omega_{\mathrm{B}}}{\hbar}, \quad (11)$$

where D_σ is the expression in the absence of magnetic field given in Eq. (5). The lowest-order contribution in $\Delta k_{\mathrm{F}}/k_{\mathrm{F}}$ to the pumped spin current is zero, while the pumped charge current is

$$I^{\mathrm{B}}_{\mathrm{charge}}=-\frac{\omega}{2\pi^2}\int d\overline{V}dk_{\mathrm{so}}8L\overline{V}\frac{4k_{\mathrm{F}}^2-\overline{V}^2}{(4k_{\mathrm{F}}^2+\overline{V}^2)^3}\,\frac{m\Omega_{\mathrm{B}}}{\hbar}. \quad (12)$$

The direction of charge flow can be reversed by changing the sign of Ω_B, i.e., of the magnetic field. The detection of this effect would constitute an indirect evidence of spin pumping.

Until now we have studied an idealized model, which allowed us to understand the physical phenomena that adiabatic spin pumping relies on. We now consider a more realistic model, which includes several modes, subband mixing induced by Eq. (1c), and the effect of the time modulation of Δ_{so}. We numerically calculate the scattering matrix within the tight-binding model, using a recursive Green's function technique.[25] The tight-binding version of the Rashba SO coupling can be written as[21]

$$H_{\mathrm{so}}=-i\gamma_{\mathrm{so}}\sum_{\sigma,\sigma'}\sum_{i,j}[c^\dagger_{i+1,j,\sigma'}(\sigma_y)_{\sigma,\sigma'}\;c_{i,j,\sigma}-c^\dagger_{i,j+1,\sigma'}(\sigma_x)_{\sigma,\sigma'}\;c_{i,j,\sigma}]+\mathrm{H.c.}, \quad (13)$$

where $c^\dagger_{i,j,\sigma}$ is the creation operator of an electron in site (i,j) with spin σ and γ_{so} is the Rashba nearest-neighbor coupling. Note that γ_{so} is related to the parameter k_{so} through the relation: $\gamma_{\mathrm{so}}=(ak_{\mathrm{so}})\gamma$, where γ is the tight-binding hopping potential and a 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 $\gamma_{\mathrm{so}}=0$, through a hopping potential Γ_{L} on the left-hand side and Γ_{R} on the right-hand side (in the following we set $\Gamma_{\mathrm{R}}=\gamma$). The Fermi energy is chosen so that three bands are occupied (from now on we express γ_{so}, Γ_{L}, and Γ_{R} in units of γ). Adiabatic pumping is obtained by performing a square cycle in the parameter space $(\gamma_{\mathrm{so}},\Gamma_{\mathrm{L}})$, with Γ_{L} varying in the range $\Gamma_0\pm\delta\Gamma/2$ and γ_{so} varying between zero and $\gamma_{\mathrm{so}}^{\mathrm{max}}$. 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 $\gamma_{\mathrm{so}}^{\mathrm{max}}$ and for fixed Γ_0. As expected, both I_{spin} and I_{charge} are increasing functions of $\delta\Gamma$. It is remarkable that for $\gamma_{\mathrm{so}}=0.042$ and $\gamma_{\mathrm{so}}=0.125$, corresponding to typical values for Rashba splitting in semiconductor,[26] I_{spin} is about two orders of magnitude larger than I_{charge} almost over the whole $\delta\Gamma$ range. Note that even for $\gamma_{\mathrm{so}}=0.25$, values which exceeds the maximum reported Rashba coupling strength,[26] I_{spin} is still much larger than I_{charge}. The pumped charge I_{charge} remains much smaller compared to I_{spin} as long as we are in the weak Rashba coupling regime,[21] in which the intersubband mixing due to Rashba coupling is negligible (in our case $\gamma_{\mathrm{so}}<0.38$). Since in our simulations $\gamma_{\mathrm{so}}\simeq0.165$ 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 $|I_{\mathrm{spin}}|$, with respect to the clean case, but keeping $|I_{\mathrm{spin}}|\gg|I_{\mathrm{charge}}|$. In the quasiballistic regime[27] adiabatic spin pumping still takes place with no qualitative difference.

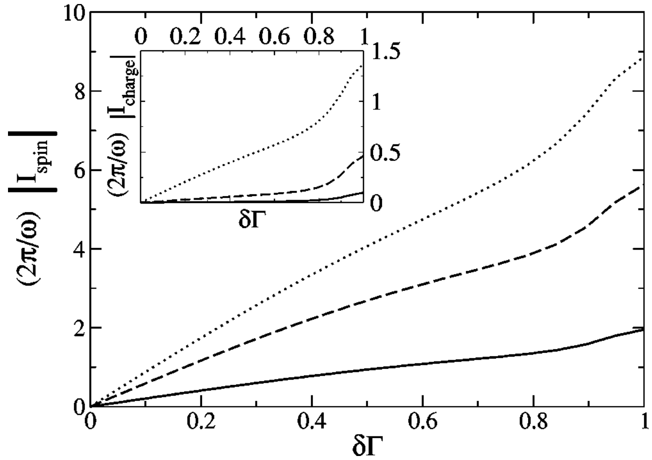

FIG. 2. Average spin and charge (in the inset) transmitted within a cycle as a function of $\delta\Gamma$ for $\gamma_{\mathrm{so}}^{\mathrm{max}}=0.042$ (full line), $\gamma_{\mathrm{so}}^{\mathrm{max}}=0.125$ (dashed line), and $\gamma_{\mathrm{so}}^{\mathrm{max}}=0.25$ (dotted line) for fixed $\Gamma_0=0.5$. For this choice of the parameters I_{spin} and I_{charge} have the same sign. The Fermi energy is set at 41/12 (in units of γ).

We finally reanalyze the spin pumping from a different perspective. To this end we start noticing that the Hamiltonian $H_{1\mathrm{D}}$ [see Eq. (1a)] in the basis of eigenstates of σ_y can be recast in the following form $H_{1\mathrm{D}}=1/2m(p_x-e\vec{A}_{\sigma,x}\cdot\hat{x})^2$, where the spin-dependent vector potential is given by $\vec{A}_\sigma=(\hbar k_{\mathrm{so}}/e)\sigma\hat{x}$. As $\vec{\nabla}\times\vec{A}_\sigma=0$, this vector potential does not describe a magnetic field. But if k_{so} varies with time τ it describes a spin-dependent electric field

$$\vec{E}_\sigma(\tau)=-\partial_\tau\vec{A}_\sigma(\tau)=-\sigma\frac{\hbar}{e}\partial_\tau k_{\mathrm{so}}(\tau)\,\hat{x}.$$

This electric field leads to a spin-dependent potential drop along the wire $V_\sigma = -\int_0^L \vec{E}_\sigma \cdot d\vec{x} = \sigma\hbar/e\,\partial_\tau k_{\rm so} L$. Thus, we can consider the wire without SO coupling but with spin-σ electrons subject to the time-dependent potential V_σ. An additional time-dependent barrier $V_{\rm bar}$ [as it was the case for the pumping cycle that led to Eq. (5)] leads to rectification of the oscillating potential V_σ. Provided that the voltage V_σ 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-σ current reads[14]

$$I_\sigma = \frac{\omega}{2\pi}\frac{e}{h}\int_0^{2\pi/\omega} |t(\tau)|^2 V_\sigma(\tau)\, d\tau, \qquad (14)$$

where t is the transmission coefficient for $k_{\rm so}=0$. 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 $k_{\rm so}$ depends on time $\Delta_{\rm so}$ also does. The only effect of the time variation of $\Delta_{\rm so}$ 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_σ, read $t_\sigma = 2k_{\rm F} e^{i\sigma k_{\rm so}L}/(2k_{\rm F} + i\overline{V})$, $r'_\sigma = -i\overline{V}/(2k_{\rm F}+i\overline{V})$.

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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 $1/10E_{\rm F}$.

[28] Further predictions can be made using Eq. (14). If in a time T, the SO coupling strength goes from zero to a value $\Delta k_{\rm so}$ 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 $(2e/h)|t|^2\int_0^T V_+(\tau)d\tau$. If one of the leads is replaced by a ferromagnet or a transverse in-plane magnetic field is present, this particular variation of $k_{\rm so}$ will result in a pulse of charge current of duration T.