Mechanically assisted spin-dependent transport of electrons

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Spin-dependent tunneling of electrons through magnetic nanostructures containing a mechanically movable quantum dot is considered. It is shown that the mechanically assisted current can be made strongly sensitive to an external magnetic field, leading to a giant magnetotransmittance effect for weak external fields of order 1–10 Oe.

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I. INTRODUCTION

Metal-organic nanocomposite materials are interesting from the point of view of the "bottom-up" approach to building future electronic devices. The ability of the organic parts of the composite materials to identify and latch onto other organic molecules is the basis for the possible self-assembly of nanoscale devices, while the metallic components provide mechanical robustness and improve the electrical conductance.

Such composite materials are heteroelastic in the sense that the mechanical rigidity of the organic and metallic components are very different. This allows for a special type of deformation, where hard metallic components embedded in a soft organic matrix can be rearranged in space at a low deformation-energy cost associated with stretching and compressing the soft matrix. Strong Coulomb forces, due to accumulation of electronic charge in embedded nanoscale metallic particles, can be a source of such mechanical deformations. This leads to a scenario where the transport of electric charge, possibly due to tunneling of electrons between metal particles, becomes a complex nanoelectromechanical phenomenon, involving an interplay of electronic and mechanical degrees of freedom.[1] Such an interplay can lead to new physics, as was recently demonstrated theoretically for the simplest possible structure—a nanoelectromechanical single-electron transistor. The electromechanical instability predicted to occur in this device at large enough bias voltage was shown to provide a new mechanism of charge transport.[2] This mechanism can be viewed as a "shuttling" of single electrons by a metallic island—a Coulomb dot—suspended between two metal electrodes. The predicted instabilily leads to a periodic motion of the island between the electrodes shuttling charge from one to the other.

The shuttle instability appears to be a rather general phenomenon. It has, e.g., been shown to occur even for extremely small suspended metallic particles (or molecules) for which the coherent quantum dynamics of the tunneling electrons[3] or even the quantum dynamics of the mechanical vibration[4–7] becomes essential. Nanomechanical transport of electronic charge can, however, occur without any such instability, e.g., in an externally driven device containing a cantilever vibrating at frequencies of order 100 MHz. A small metallic island attached to the tip of the vibrating cantilever may shuttle electrons between metallic leads as has recently been demonstrated.[8] Further experiments with magnetic and superconducting externally driven shuttles, as suggested in Ref. 9, seem to be a natural extension of this work. Fullerene-based nanomechanical structures[10] are also of considerable interest.

The possibility to place transition-metal atoms or ions inside organic molecules introduces an additional "magnetic" degree of freedom that allows the electronic spins to be coupled to mechanical and charge degrees of freedom.[12] By manipulating the interaction between the spin and external magnetic fields and/or the internal interaction in magnetic materials, spin-controlled nanoelectromechanics may be achieved. An inverse phenomenon—nanomechanical manipulation of nanomagnets—was suggested earlier in Ref. 11. A magnetic field, by inducing the spin of electrons to rotate (precess) at a certain frequency, provides a clock for studying the shuttle dynamics and a basis for a dc spectroscopy of the corresponding nanomechanical vibrations.

A particularly interesting situation arises when electrons are shuttled between electrodes that are half-metals. A half-metal is a material that not only has a net magnetization as do ferromagnets, but all the electrons are in the same spin state; the material is fully spin-polarized. Examples of such materials can be found among the perovskite maganese oxides, a class of materials that shows an intrinsic, so-called colossal magnetoresistance effect at high magnetic fields (of order 10–100 kOe).[13]

A large magnetoresistance effect at lower magnetic fields has been observed in layered tunnel structures where two thin perovskite manganese oxide films are separated by a tunnel barrier.[13–17] Here the spin polarization of electronic states crucially affects the tunneling between the magnetic electrodes. This is because electrons that can be extracted from the source electrode have their spins aligned in a definite direction, while electrons that can be injected into the drain electrode must also have their spins aligned, possibly in a different direction. Clearly the tunneling probability and, hence, the resistance must be strongly dependent on the relative orientation of the magnetization of the two electrodes. An external magnetic field aligns the magnetization direction of the two films at different field strengths, so that the relative magnetization can be changed between high- and low-resistance configurations. A change in the resistance of

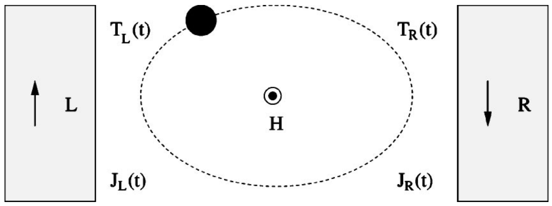

FIG. 1. Schematic view of the nanomechanical GMR device: a movable dot with a single-electron level couples to the leads due to tunneling of electrons, described by the tunneling probability amplitudes $T_{L,R}(t)$, and to the exchange interaction whose strength is denoted by $J_{L,R}(t)$. An external magnetic field H is oriented perpendicular to the direction of the magnetization in the leads (arrows).

trilayer devices by factors of order 2–5 have in this way been induced by magnetic fields of order 200 Oe.[14–16] The required field strength is determined by the coercivities of the magnetic layers. This makes it difficult to use a tunneling device of the described type for sensing very low magnetic-fields. In this paper we propose a different functional principle—spin-dependent shuttling of electrons—for low-magnetic-field sensing purposes. We will show that this principle can lead to a giant magnetoresistance effect in external fields as low as 1–10 Oe.

The idea that we propose to pursue is to use the external magnetic field to manipulate the *spin of shuttled electrons* rather than the magnetization of the leads. The possibility to "trap" electrons on a nanomechanical shuttle (decoupled from the magnetic leads) during quite a long time on the scale of the time it takes an electron to tunnel on/off the shuttle makes it possible for even a weak external field to rotate the electron's spin to a significant degree. Such a rotation allows the spin of an electron, loaded onto the shuttle from the spin-polarized source electrode, to be reoriented in order to allow the electron finally to tunnel from the shuttle to the spin-polarized drain lead. As we will show below, the magnetic field-induced spin-rotation of shuttled electrons is a very sensitive nanomechanical mechanism for a giant magnetoresistance (GMR) effect.

II. FORMULATION OF THE PROBLEM: GENERAL EXPRESSION FOR THE CURRENT

A schematic view of the nanomechanical GMR device to be considered is presented in Fig. 1. Two fully spin-polarized magnets with fully spin-polarized electrons serve as source and drain electrodes in a tunneling device. In this paper we will consider the situation when the electrodes have exactly opposite polarization. A mechanically movable quantum dot [described by a time-dependent displacement $x(t)$], where a single energy level is available for electrons, performs forced harmonic oscillations with period $T=2\pi/\omega$ between the leads. The external magnetic field is perpendicular to the orientation of the magnetization in both leads.

The Hamiltonian that governs the dynamical evolution of the system is

$$\hat{\mathcal{H}} = \hat{\mathcal{H}}_{leads} + \hat{\mathcal{H}}_{dot} + \hat{\mathcal{H}}_{int}; \tag{1}$$

$$\hat{\mathcal{H}}_{leads} = \sum_{\alpha,\kappa} \varepsilon_\alpha a^\dagger_{\alpha,\kappa} a_{\alpha,\kappa};$$

$$\hat{\mathcal{H}}_{dot} = \sum_\sigma \varepsilon_0 a^\dagger_\sigma a_\sigma - (g\mu H/2)\sum_\sigma a^\dagger_\sigma a_{-\sigma} - J_L(t)(a^\dagger_\uparrow a_\uparrow - a^\dagger_\downarrow a_\downarrow) - J_R(t)(a^\dagger_\downarrow a_\downarrow - a^\dagger_\uparrow a_\uparrow);$$

$$\hat{\mathcal{H}}_{int} = T_L(t)\sum_\alpha a^\dagger_{\alpha,L} a_\uparrow + T_R(t)\sum_\alpha a^\dagger_{\alpha,R} a_\downarrow + \text{H.c.},$$

where $a^\dagger_{\alpha,\kappa}, (a_{\alpha,\kappa}), \kappa=L,R$, are the creation (annihilation) operators of electrons with energy ε_α in the left (right) lead (we have suppressed the spin indices for the electronic states in the leads based on the assumption of full spin polarization), $a^\dagger_\sigma(a_\sigma), \sigma=\uparrow,\downarrow$, are the creation (annihilation) operators on the dot, ε_0 is the energy of the on-dot level, $J_{L(R)}(t) \equiv J_{L(R)}(x(t))$ are the exchange interactions between the on-grain electron and the left (right) lead, $T_{L(R)}(t) \equiv T_{L(R)}(x(t))$ are the tunnel coupling amplitudes, g is the gyromagnetic ratio, and μ is the Bohr magneton.

The single-electron density matrix describing electronic transport between the leads may be expressed as

$$\hat{\rho} = \sum_{\alpha,\kappa} w_{\alpha,\kappa} |\Psi^{\alpha,\kappa}\rangle\langle\Psi^{\alpha,\kappa}|. \tag{2}$$

Here $|\Psi^{\alpha,\kappa}\rangle$ are single-electron states that obey the time-dependent Shrödinger ($\hbar=1$) equation with the Hamiltonian (1). The initial condition has the form

$$|\Psi^{\alpha,\kappa}(t\to-\infty)\rangle = |\alpha,\kappa\rangle \exp(-i\varepsilon_\alpha t),$$

where $|\alpha,\kappa\rangle$ is a single-electron state on the lead κ with energy ε_α.

We will assume that internal relaxation in the leads is fast enough to lead to equilibrium distributions of the electrons. Referring to Eq. (2) this means that $w_{\alpha,L(R)}=f(\varepsilon_\alpha \mp V/2)$, where $f(\varepsilon)$ is the Fermi distribution function and V is the applied voltage.

The problem at hand is greatly simplified if one considers the large bias-voltage limit

$$V - \varepsilon_0 \gg \nu T^2_{\max}, \tag{3}$$

where ν is the density of states in the leads and $T_{\max} = \max T_{L,R}(t)$. The restriction (3) does not allow us to consider a narrow ($\sim \nu T^2_{\max}$) transition region, where the voltage-shifted Fermi level in one of the leads crosses the resonant level on the dot. However, it does cover the case most important in reality, i.e., when the fully transmissive junction is strongly affected by electronic spin polarization. Therefore, in our further considerations we will take $w_{\alpha,L}=1$, $w_{\alpha,R}=0$, and $\varepsilon_0=0$.

We now calculate the average current I through the system from the standard relation

$$I = \frac{1}{T}\int_0^T dt\, \text{Tr}\{\hat{\rho}\hat{j}\}, \tag{4}$$

$$\hat{j} = e\frac{\partial \hat{N}_R}{\partial t} = ieT_R(t)\sum_\alpha (a_\downarrow^\dagger a_{\alpha,R} - a_{\alpha,R}^\dagger a_\downarrow),$$

where $\hat{N}_R = \Sigma_\alpha a_{\alpha,R}^\dagger a_{\alpha,R}$ is the electron number operator for the right lead.

In general, the state $|\Psi^{\alpha,L}\rangle$ can be expressed as

$$|\Psi^{\alpha,L}(t)\rangle = \sum_\sigma c_\sigma^\alpha(t)|\sigma\rangle + \sum_{\beta,\kappa} c_\kappa^{\alpha,\beta}(t)|\beta,\kappa\rangle. \tag{5}$$

Thus the problem is reduced to determining the coefficients $c_\kappa^{\alpha,\beta}(t)$ and $c_\sigma^\alpha(t)$. At this point it is convenient to introduce the bivectors

$$\mathbf{c}^\alpha = \begin{pmatrix} c_\uparrow^\alpha \\ c_\downarrow^\alpha \end{pmatrix}, \quad \mathbf{e}_1 = \begin{pmatrix} 1 \\ 0 \end{pmatrix}, \quad \mathbf{e}_2 = \begin{pmatrix} 0 \\ 1 \end{pmatrix},$$

so that the coefficients $c_\kappa^{\alpha,\beta}$ can be expressed as (see Appendix A)

$$c_R^{\alpha,\beta} = -i\int_{-\infty}^t dt' e^{i\varepsilon_\beta(t-t')}T_R(t')(\mathbf{e}_2, \mathbf{c}^\alpha(t')),$$

$$c_L^{\alpha,\beta} = e^{-i\varepsilon_\beta t}\delta_{\alpha\beta} - i\int_{-\infty}^t dt' e^{i\varepsilon_\beta(t-t')}T_L(t')(\mathbf{e}_1, \mathbf{c}^\alpha(t')). \tag{6}$$

Here $(\mathbf{a}, \mathbf{b})$ is the inner product of two bivectors. As shown in Appendix A, by using the wide-band approximation (i.e., by taking the electron density of states in the leads ν to be constant) the equation for the bivectors $\mathbf{c}^\alpha$ takes the form

$$i\frac{\partial \mathbf{c}^\alpha}{\partial t} = \hat{R}(t)\mathbf{c}^\alpha + \mathbf{f}^\alpha(t). \tag{7}$$

Here $\mathbf{f}^\alpha(t) = T_L(t)e^{-i\varepsilon_\alpha t}\mathbf{e}_1$ and the matrix $\hat{R}(t)$ is

$$\hat{R}(t) = \begin{pmatrix} -J(t) - i\Gamma_L(t)/2 & -g\mu H/2 \\ -g\mu H/2 & J(t) - i\Gamma_R(t)/2 \end{pmatrix}, \tag{8}$$

where $J(t) = J_L(t) - J_R(t)$ and $\Gamma_\kappa(t) = 2\pi\nu T_\kappa^2(t)$ is the level width.

The formal solution of Eq. (7) can be written in the form

$$\mathbf{c}^\alpha(t) = -i\int_{-\infty}^t dt' \hat{L}(t,t')\mathbf{f}^\alpha(t'), \tag{9}$$

where the "evolution" operator $\hat{L}(t,t')$, with $\hat{L}(t,t) = \hat{I}$, is defined as the solution of the equation

$$i\frac{\partial \hat{L}(t,t')}{\partial t} = \hat{R}(t)\hat{L}(t,t'), \tag{10}$$

and obeys the multiplication and periodicity rules

$$\hat{L}(t,t') = \hat{L}(t,t'')\hat{L}(t'',t'), \quad \hat{L}(t+T,t'+T) = \hat{L}(t,t'). \tag{11}$$

Using Eq. (9) together with Eq. (4), one can write the average current on the form

$$I = \frac{e}{T}\int_0^T dt\Gamma_R(t)\int_{-\infty}^t dt'\Gamma_L(t')|L_{21}(t,t')|^2, \tag{12}$$

where $L_{21}(t,t')$ is a matrix element of the operator $\hat{L}(t,t')$; $L_{21}(t,t') = (\mathbf{e}_2, \hat{L}(t,t')\mathbf{e}_1)$.

Since the probability amplitude for tunneling is exponentially sensitive to the position of the dot, the maximum of the tunnel exchange interaction between an electron on the dot and an electron in one lead occurs when the tunneling coupling to the other lead is negligible. This is why we will assume the following property of the tunneling amplitude $T_\kappa(t)$ to be fulfilled:

$$T_L(t)T_R(t) = 0, \quad T_L(t), T_R(t) \neq 0. \tag{13}$$

This assumption allows us to divide the time interval $(0,T)$ into the intervals $(0,\tau)+(\tau,T/2)+(T/2,T/2+\tau)+(T/2+\tau,T)$. We suppose that $T_L(t)\neq 0$ (but $H=0$) only in the time interval $(0,\tau)$ [and, analogously, $T_R(t)\equiv T_L(t+T/2)\neq 0$ in the time interval $(T/2,T/2+\tau)$]. Using this approximation together with the properties (11) of the operator $\hat{L}(t,t')$, we arrive at the following expression for the average current (Appendix B):

$$I = \frac{e}{T}(1-e^{-\Gamma})^2\sum_{n=0}^\infty |(\mathbf{e}_2, \hat{L}(T/2,\tau)\hat{L}^n\mathbf{e}_1)|^2. \tag{14}$$

Here $\hat{L} \equiv \hat{L}(T+\tau,\tau)$ and

$$\Gamma = 2\pi\nu\int_0^\tau dt T_L^2(t) \tag{15}$$

is the probability for an electron to be transferred to the shuttle during the contact time τ. Consequently, in order to calculate the average current it is necessary to investigate the properties of the evolution operator $\hat{L}$. It follows from its definition and our approximation [Eq. (13)] that

$$\hat{L} = e^{-(1+\sigma_3)\Gamma/4+i\sigma_3\Phi_0}\hat{L}(T,T/2+\tau)e^{-(1-\sigma_3)\Gamma/4-i\sigma_3\Phi_0}\hat{L}(T/2,\tau), \tag{16}$$

where $\Phi_0 = \int_0^\tau dtJ(t)$. During the time interval $T/2+\tau < t < T$ (when under our approximation $\Gamma_{L(R)}=0$) the operator $\hat{R}(t)$ is Hermitian and possesses the symmetry properties $\sigma_2\hat{R}^*\sigma_2 = -\hat{R}$. As a consequence, the operator $\hat{U}\equiv\hat{L}(T,T/2+\tau)$ is unitary and completely determined by the probability amplitude $\gamma e^{i\varphi}$ for a spin-flip transition. It can be written as

$$\hat{U} = \begin{pmatrix} \sqrt{1-\gamma^2} & i\gamma e^{i\varphi} \\ i\gamma e^{-i\varphi} & \sqrt{1-\gamma^2} \end{pmatrix}, \tag{17}$$

where the modulus γ and phase φ depend on both the exchange interaction $J(t)$ and the magnetic field H. In addition, the symmetry properties $\sigma_3\hat{R}(T-t)\sigma_3 = -\hat{R}(t)$ gives the relation

$$\hat{L}(T/2,\tau) = \sigma_1 \hat{L}(T,T/2-\tau)\sigma_1. \quad (18)$$

As a result, with the help of Eqs. (16) and (18), the operator $\hat{L}$ can be expressed as

$$\hat{L} = e^{-\Gamma/2}(e^{-\sigma_3\Gamma/4+i\Phi_0\sigma_3}\hat{U}\sigma_1)^2. \quad (19)$$

Proceeding with the analysis we (i) calculate the eigenvalues λ_i and eigenvectors $\mathbf{b}_i$ of the operator $\hat{L}$ of Eq. (19), $\hat{L}\mathbf{b}_i=\lambda_i\mathbf{b}_i$, and (ii) substitute the expansion $\mathbf{e}_i=a_{ij}\mathbf{b}_j$ into Eq. (14) and calculate the average current. The result is

$$I = \frac{e}{T}\sinh\Gamma/2\frac{\cosh\Gamma/2+\cos 2\vartheta}{1+\cos 2\vartheta\cosh\Gamma/2+\eta\sinh^2\Gamma/2}, \quad (20)$$

where $\vartheta=\varphi+\Phi_0$ and $\eta=(1+\gamma^2)/2\gamma^2$. Equation (20) for the average current is our main result, which now has to be analyzed further.

III. CALCULATION OF THE CURRENT IN THE LIMITS OF STRONG AND WEAK EXCHANGE COUPLING

Although the result (20) for the tunnel current is compact, it is, in general, a rather complicated problem to find the magnetic-field dependence of the coefficient η, which depends on the probability amplitude γ for flipping the spin of shuttled electrons. Three different timescales are involved in the spin dynamics of a shuttled electron. They correspond to three characteristic frequencies: (i) the frequency of spin rotation, determined by the tunnel exchange interaction with the magnetic leads; (ii) the frequency of spin rotation in the external magnetic field, and (iii) the frequency of shuttle vibrations. Different regimes occur depending on the relation between these time scales. Here we will consider two limiting cases, where a simple solution of the problem can be found. Those are the limits of weak, $J_0 \ll \mu H$, and strong, $J_0 \gg \mu H$, exchange interactions between the dot and the leads. Here $J_0=\max J(t)$.

A. Strong exchange interaction between dot and leads

A strong magnetic coupling to the leads, $J_0 \gg \mu H$, preserves the electron-spin polarization, preventing spin-flips of shuttled electrons due to an external magnetic field. However, if the magnetization of the two leads are in opposite directions, the exchange coupling to the leads have a different sign. Therefore, the exchange couplings to the two leads tend to cancel out when the dot is in the middle of the junction. Hence the strong exchange interaction affecting a dot electron depends on time and periodically changes sign, being arbitrary small close to the time of sign reversal. In Fig. 2 the on-dot electronic energy levels for spins parallel and antiparallel to the lead magnetization are presented as a function of time. The effect of an external magnetic field is in the limit $J_0 \gg \mu H$ negligible almost everywhere, except in the vicinity of the level crossing. At this "time point," which we denote t_{LZ}, the external magnetic field removes the degeneracy and a gap is formed in the spectrum. The probability of electronic spin-flip in this case is determined by the probability of a Landau–Zener reflection from the gap formed by the

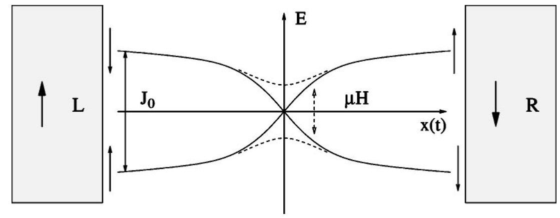

FIG. 2. On-dot energy levels for spin-up and spin-down electron states as a function of the position of the dot. The level crossing in the middle of the device is removed by an external magnetic field, so that a gap is formed in the energy spectrum (dashed lines).

magnetic field (in this case a Landau–Zener transition across the gap is a mechanism for backscattering of the electron, as this is the channel where the electronic spin is preserved). The matrix operator $\hat{U}$, parametrized by γ and φ in Eq. (17), can readily be expressed in terms of the Landau–Zener scattering amplitudes. The phase $\varphi=\varphi_0+\Phi_1$, where φ_0 is the Landau–Zener phase shift and $\Phi_1=\int_\tau^{T/2}dtJ(t)$, while γ^2 is given by the probability of Landau-Zener backscattering,

$$\gamma^2 = 1-\exp\left[-\frac{\pi(g\mu H)^2}{8|J'(t_{LZ})|}\right]. \quad (21)$$

The magnetic-field dependence of the current, calculated from Eq. (20) using these results, is shown in Fig. 3. The width δH of the dip at small fields in the function $I=I(H)$ can be found directly from Eqs. (20) and (21). Restoring dimensions one finds that

$$\delta H \simeq \frac{1}{\mu}\sqrt{J_0\hbar\omega}. \quad (22)$$

B. Weak exchange interaction

In the limit $J_0 \ll \mu H$ one may neglect the influence of the magnetic leads on the on-dot electron-spin dynamics. In this

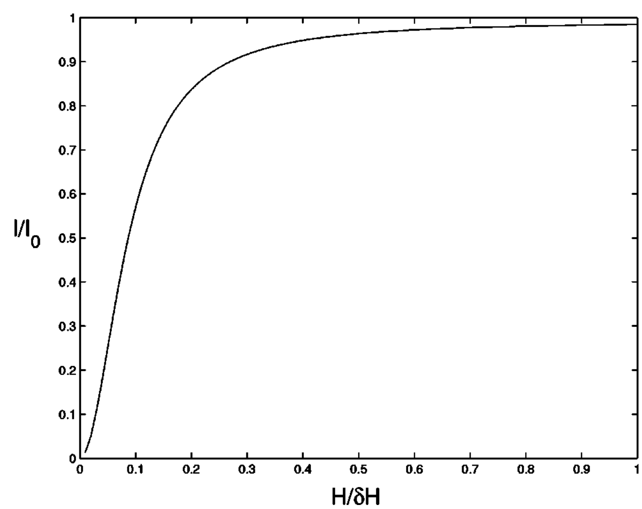

FIG. 3. The current I in units of $I_0 \simeq e\Gamma/2T$ plotted as a function of normalized magnetic field $H/\delta H$ for the limiting case of a strong exchange coupling between dot and leads. The values $\Gamma=0.3$ and $\int_0^T dtJ(t)=\pi/6$ were used for this plot; δH is defined by Eq. (22).

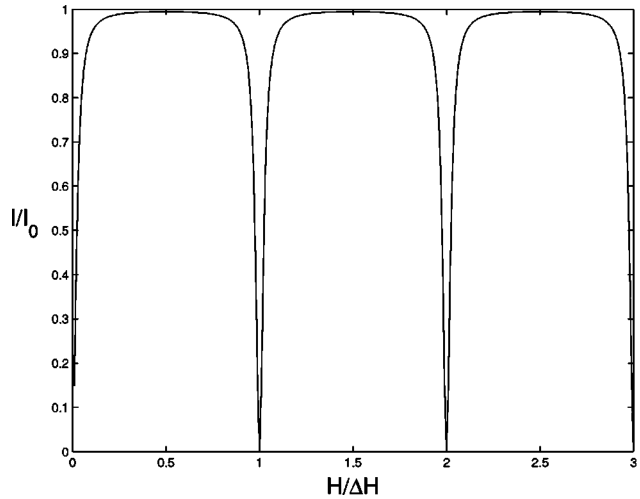

FIG. 4. The current I in units of $I_0 \simeq e\Gamma/2T$ plotted as a function of normalized magnetic field $H/\Delta H$ for the limiting case of a weak exchange coupling between dot and leads. ΔH is defined by Eq. (25) and the value $\Gamma=0.3$ was used.

case the matrix $\hat{U}$ given by Eq. (17) can easily be calculated and Eq. (20) reduces to

$$I = \frac{2e}{T}\frac{\sin^2 \vartheta/2 \tanh \Gamma/4}{\sin^2 \vartheta/2 + \tanh^2 \Gamma/4}, \tag{23}$$

where $\vartheta = g\mu\int_\tau^{T/2} dtH$ is the rotation angle of the spin in the external field.

Two different scales for the external magnetic field determine the magnetotransmittance in this limit. One scale is associated with the width in magnetic field of the resonant behavior of the transmittance [see the denominator in Eq. (23)]. This scale is

$$\delta H \simeq \Gamma\frac{\hbar\omega}{g\mu}, \tag{24}$$

where $\omega = 2\pi/T$ is the shuttle vibration frequency. The second scale,

$$\Delta H \simeq \frac{\hbar\omega}{g\mu}, \tag{25}$$

comes from the periodic function $\sin^2 \vartheta/2$ that enters Eq. (23) (the estimations of both δH and ΔH were done under the natural asumption that $\tau \ll T$). The magnetic-field dependence of the current is presented in Fig. 4. Dips in the transmittance of width δH appear periodically as the magnetic field is varied, the period being ΔH. This amounts to a giant magnetotransmittance effect. It is interesting to notice that by measuring the period of the variations in $I(H)$ one can in principle determine the shuttle vibration frequency. This amounts to a dc method for spectroscopy of the nanomechanical vibrations. Equation (25) gives a simple relation between the vibration frequency and the period of the current variations. The physical meaning of this relation is very simple: every time when $\omega/\Omega = n+1/2$, where Ω is the spin precession frequency in the applied magnetic field, the shuttled electron is able to fully flip its spin to remove the "spin-blockade" of tunneling between spin polarized leads having their magnetization in opposite directions.

IV. CONCLUSIONS

The analysis presented above demonstrates the possibility of a giant magnetotransmittance effect caused by shuttling of spin-polarized electrons between magnetic source and drain electrodes. The sensitivity of the shuttle current to an external magnetic field is determined, according to Eq. (24), by the transparency of the tunnel barriers. By diminishing the tunneling transmittance one can increase the sensitivity of the device to an external magnetic field. The necessity to have a measurable current determines the limit of this sensitivity. In the low transparency limit, $\Gamma \ll 1$, the current through the device can be estimated as $I \simeq e\Gamma\omega$. If one denotes the critical field that determines the sensitivity of the device by H_{cr}, one finds from Eq. (24) that $H_{cr} \simeq \delta H$. The critical field can now be expressed in terms of the current transmitted through the device as

$$H_{cr}(\mathrm{Oe}) \simeq \frac{\hbar I}{e\mu g} \simeq 10^2 (g_0/g) I\ (\mathrm{nA}), \tag{26}$$

where $g_0(=2)$ is the gyromagnetic ratio for the free electrons. For $I \simeq 10^{-1}-10^{-2}$ nA and $g_0/g \simeq 1$, this gives a range $H_{cr} \simeq 1-10$ Oe. A further increase in sensitivity would follow if one could use a shuttle with several (N) electronic levels involved in the tunneling process. The critical magnetic field would then be inversely proportional to the number of levels, $H_{cr}(N) = H_{cr}(N=1)/N$.

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APPENDIX A

The Shrödinger equation results in equations for the coefficients $c_\kappa^{\alpha\beta}, c_\sigma^\alpha$

$$i\frac{\partial c_\uparrow^\alpha}{\partial t} = -J(t)c_\uparrow^\alpha - (g\mu H/2)c_\downarrow^\alpha + T_L(t)\sum_\beta c_L^{\alpha\beta}, \tag{A1}$$

$$i\frac{\partial c_\downarrow^\alpha}{\partial t} = J(t)c_\uparrow^\alpha - (g\mu H/2)c_\downarrow^\alpha + T_R(t)\sum_\beta c_R^{\alpha\beta},$$

$$i\frac{\partial c_L^{\alpha\beta}}{\partial t} = \varepsilon_\beta c_L^{\alpha\beta} + T_L(t)c_\uparrow^\alpha,$$

$$i\frac{\partial c_R^{\alpha\beta}}{\partial t} = \varepsilon_\beta c_R^{\alpha\beta} + T_R(t)c_\downarrow^\alpha.$$

As it follows from the last two equations (together with the initial conditions):

$$c_R^{\alpha\beta}(t) = -i\int_{-\infty}^{t} dt' e^{i\varepsilon_\beta(t'-t)}T_R(t')c_\downarrow^\alpha(t'), \qquad (A2)$$

$$c_L^{\alpha\beta}(t) = e^{-i\varepsilon_\beta t}\delta_{\alpha\beta} - i\int_{-\infty}^{t} dt' e^{i\varepsilon_\beta(t'-t)}T_L(t')c_\uparrow^\alpha(t').$$

Therefore, for the $\Sigma_\beta c_R^{\alpha\beta}(t)$ one gets

$$\sum_\beta c_R^{\alpha\beta}(t) = -i\int_{-\infty}^{t} dt' T_R(t')c_\downarrow^\alpha(t')\sum_\beta e^{i\varepsilon_\beta(t'-t)}.$$

In the wide-band approximation we suppose that $\nu(\varepsilon) = \text{const}$, therefore, $\Sigma_\beta e^{i\varepsilon_\beta(t'-t)} = 2\pi\nu\delta(t'-t)$ and

$$\sum_\beta c_R^{\alpha\beta}(t) = -i\pi\nu T_R(t)c_\downarrow^\alpha(t). \qquad (A3)$$

Analogously,

$$\sum_\beta c_L^{\alpha\beta}(t) = e^{-i\varepsilon_\alpha t} - i\pi\nu T_L(t)c_\uparrow^\alpha(t). \qquad (A4)$$

Substituting the expressions Eqs. (A3) and (A4) into the first two equations (A1), one gets Eq. (7) for the bivector $\mathbf{c}^\alpha$.

APPENDIX B

Under our approximation we can change the integration limits in Eq. (12)

$$I = \frac{e}{T}\int_{T/2}^{T/2+\tau} dt\Gamma_R(t)\int_{-\infty}^{\tau} dt'\Gamma_L(t')|\hat{L}_{21}(t,t')|^2. \qquad (B1)$$

Beside this, in the time moments $T/2 < t < T/2+\tau$, $\hat{L}(t,T/2)$ is a diagonal matrix and therefore $\hat{L}_{21}(t,t') = \hat{L}_{22}(t,T/2)\hat{L}_{21}(T/2,t')$. As a consequence, the integral in the expression for the average current [Eq. (B1)] is factorized

$$I = \frac{e}{T}\int_{T/2}^{T/2+\tau} dt\Gamma_R(t)|\hat{L}_{22}(t,T/2)|^2\int_{-\infty}^{\tau} dt'\Gamma_L(t')|\hat{L}_{21}(T/2,t')|^2. \qquad (B2)$$

The first integral in Eq. (B2) is easy to calculate,

$$\int_{T/2}^{T/2+\tau} dt\Gamma_R(t)|\hat{L}_{22}(t,T/2)|^2 = 1 - e^{-\Gamma}, \qquad (B3)$$

where quantity Γ is defined in Eq. (15). The second integral in Eq. (B2) can be transformed in the following manner:

$$\int_{-\infty}^{\tau} dt\Gamma_L(t)|\hat{L}_{21}(T/2,t)|^2 = \sum_{n=0}^{\infty}\int_0^{\tau} dt\Gamma_L(t)|(\mathbf{e}_2,\hat{L}(T/2,\tau)\hat{L}(\tau,t-nT)\mathbf{e}_1)|^2. \qquad (B4)$$

For the quantity $\hat{L}(\tau,t-nT) \equiv \hat{L}(\tau+nT,t)$ one has

$$\hat{L}(\tau+nT,t) = \left[\prod_{k=1}^{n}\hat{L}(\tau+kT,\tau+(k-1)T)\right]\hat{L}(\tau,t) = \hat{L}^n(\tau+T,\tau)\hat{L}(\tau,t). \qquad (B5)$$

Substituting this expression in Eq. (B4) and calculating the integral in the same manner, as in Eq. (B3), one gets the Eq. (14) for the average current.

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