Tunneling spectrum of single-spin oscillations

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We consider tunneling between electrodes via a microscopic system which can be modeled by the two-level Hamiltonian (a localized spin 1/2 or a quantum dot). Using the non-equilibrium Keldysh method we find the dependence of the current-current correlation function on the applied voltage V between leads, temperature T, (constant) magnetic field $\mathbf{B}_0$ acting on the spin, and on the degree and orientation $\mathbf{m}_{R,L}$ of spin polarization of electrons in the two (R, L) leads. We show that the current-current correlation function exhibits a peak caused by Larmor-precession when $\mathbf{m}_{R,L}$ are not parallel to $\mathbf{B}_0$ and find the conditions when the signal-to-noise ratio R for this peak reaches its maximum of order 1, $R \leq 4$. We compare our results with those obtained in the quasiclassical approach and discuss the experimental results observed by STM dynamic probes of spin.

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A potential use of solid state qubits as building blocks of a prospective quantum computer has attracted interest in the problem of measurement of tunneling currents via a microscopic system that can be modeled by a two-level Hamiltonian (such as a quantum dot, or a molecule or an atom with a localized spin [1–4]). A fundamental question that arises is what signatures of the spin dynamics are encoded in the tunneling current and how this current affects the spin dynamics.

Scanning tunneling microscopy (STM) experiments, in the presence of a magnetic field $\mathbf{B}_0$, have reported precession of individual spins at the Larmor frequency, manifested as a narrow peak in the tunneling current power spectrum with a signal-to-noise (shot noise) ratio larger than unity [5, 6]. In the non-relativistic approximation the tunneling electrons are coupled to the spin by exchange interaction so the electrons with spin alligned with $\mathbf{B}_0$ do not couple to the oscillatory components of spin (which are perpendicular to $\mathbf{B}_0$). The experimental results [5, 6] are difficult to understand in the non-relativistic model because electrons in the leads were polarized by the same magnetic field which acted on the spin.

Detection of quantum spin oscillations has been studied by Korotkov and Averin [3] and by Ruskov and Korotkov [4] in the context of electron tunneling via a double quantum dot representing a qubit. This problem bares a strong similarity to the problem of tunneling via a localized spin $1/2$. Using the Bloch equations describing the ensemble-averaged evolution of the density matrix for the coupled qubit-detector system, they obtain oscillations at the Larmor frequency in the spectral density of noise $S_I(\omega)$, with $R \leq 4$ in the case of weak interaction between the qubit and the detector. A similar quasiclassical approach for tunneling of electrons via a single spin was used in Ref. [1]. In such a treatment, the action of the electrons on the spin is replaced by the action of an effective *classical* magnetic field with a shot noise spectrum depending on the tunneling current between the leads. Since the quasiclassical approach does not account for the quantum nature of the tunneling electrons a more elaborate solution of the quantum transport equation for both spin and electrons is desired.

An important step towards a full quantum mechanical treatment of electron tunneling via spin was recently given by Parcollet and Hooley [2]. For the case when electrons are polarized along the direction of the magnetic field acting on spin, they computed the spin magnetization of a quantum dot in the two-leads Kondo model as a function of the temperature of the electrons, magnetic field, and voltage between the leads (which exceeds the Kondo temperature T_K). Their striking result is that, even at zero order in the spin-leads coupling, the magnetization is not given by the thermal equilibrium expression $M_{\text{eq}} = (1/2)\tanh(B/2T)$ when the spin is more strongly coupled to the leads than to any other thermal bath. At long times, the state of the spin is completely determined by the properties of the leads, and since the system is out of equilibrium, the steady state thus achieved is not described by a Gibbs distribution function, but is voltage-dependent. They also showed that the correct way to calculate the perturbative corrections to the spin distribution function is by solving a quantum transport equation self-consistently.

In this article we present the conditions under which a dc steady state can be achieved in which the spectral density of the current-current correlation function $S_I(\omega)$ exhibits a peak at (a renormalized) Larmor frequency, ω_L. We calculate $R(V, T, \mathbf{B}, \mathbf{m}_{R,L})$, where $\mathbf{B}$ is the effective magnetic field acting on spin [1]. For a more detailed discussion, interested readers are refered to [8]. We will extend the treatment of Parcolett and Hooley [2] to arbitrary orientations of $\mathbf{m}_{\mathbf{R}}, \mathbf{m}_L$ with respect to $\mathbf{B}$ and account for direct tunneling of electrons between leads. In the case of polarization of the lead electrons with a nonzero component perpendicular to $\mathbf{B}$, we find a peak in the noise power $S_I(\omega)$ at ω_L. Such a peak was absent in the situation $\mathbf{m} \parallel \mathbf{B}$ discussed in Ref. 2. We show that signal-to noise ratio is small, $R \ll 1$, for weak spin polarization of leads electrons or if spin is coupled to the environment much stronger than to leads. The height of the peak at Larmor frequency in $S_I(\omega)$ depends on a combination of voltage V and the magnetic field B at low temperatures in such a way that the signal-to-noise ratio reaches its maximum at the resonance condition $eV = \hbar\omega_L$. The resonance occurs because individual electrons can transfer energy to the spin excitation and thus the tunneling process has a similarity with the scattering of an individual particle by the spin. This effect was missing in the quassiclassical approach [1, 3, 4].

We note that from a more general perspective the electron tunneling via spin represents an example of indirect quantum measurement. The spin is probed by tunneling electrons whose correlation function is measured in a continuous fashion with a classical apparatus (which is not included in the description). This is a weak continuous measurement, in which the signal-to-noise ratio cannot exceed a certain maximum value, as found by Averin and Korotkov [3].

For the system consisting of a spin coupled to leads in the nonrelativistic approximation (no spin-orbit interaction), we use the Hamiltonian of the two-leads Kondo model [2] with a direct tunneling term included

$$\begin{aligned}
\mathcal{H} &= \mathcal{H}_0 + \mathcal{H}_T, \quad \mathcal{H}_T = \mathcal{H}_{\text{ref}} + \mathcal{H}_{\text{tr}}, \qquad (1)\\
\mathcal{H}_0 &= \sum_{\alpha,\mathbf{k},\sigma,\sigma'} [\epsilon_{\mathbf{k}\alpha}\delta_{\sigma\sigma'} - \frac{1}{2}B_\alpha \vec{\sigma}_{\sigma\sigma'}] c^\dagger_{\alpha\mathbf{k}\sigma} c_{\alpha\mathbf{k}\sigma'} - BS_z,\\
\mathcal{H}_{\text{ref}} &= \sum_{\alpha,\mathbf{k},\mathbf{k}',\sigma,\sigma'} T_{\alpha\alpha}\, \mathbf{S}\cdot(c^\dagger_{\alpha\mathbf{k}\sigma}\vec{\sigma}_{\sigma\sigma'} c_{\alpha\mathbf{k}'\sigma'}),\\
\mathcal{H}_{\text{tr}} &= \sum_{\mathbf{k},\mathbf{k}',\sigma,\sigma'} (T_0\, \delta_{\sigma\sigma'} + T_{RL}\, \mathbf{S}\cdot\vec{\sigma}_{\sigma\sigma'}) c^\dagger_{R\mathbf{k}\sigma} c_{L\mathbf{k}'\sigma'} + \text{H.c.},
\end{aligned}$$

where $c^\dagger_{\alpha\mathbf{k}\sigma}$ ($c_{\alpha\mathbf{k}\sigma}$) creates (annihilates) an electron in the left or right lead (depending on $\alpha \in \{L, R\}$) with momentum $\mathbf{k}$ and spin σ, $\epsilon_{\mathbf{k}\alpha} = \epsilon_{\mathbf{k}} - \mu_\alpha$, $\vec{\sigma}$ represents the three Pauli matrices, while T_{LL}, T_{RR} and $T_{LR} = (T_{RL})^*$ are tunneling matrix elements for the tunneling from leads to the spin-1/2 described by the operator $\mathbf{S} = (S_x, S_y, S_z)$ and T_0 is the direct tunneling matrix element. We describe the leads by a free electron gas with a density of states $\rho(\epsilon)$ of bandwidth D, so that $\epsilon_{\mathbf{k}}$ is the bare energy of the electron with momentum $\mathbf{k}$, the same in both leads. We assume weak tunneling, $|T_0|^2\rho_0^2, |T_{RL}|^2\rho_0^2 \ll 1$, where ρ_0 is the density of states of the leads at the Fermi level. The voltage is given by the difference in chemical potentials, $\mu_R - \mu_L = V$. We also assume $T_0 \gg T_{\alpha,\beta}$ ($\alpha, \beta \in \{R, L\}$). The leads are supposed to be in thermal equilibrium at temperature T. H_{ref} describes the spin-flip scattering of an electron from one lead back into the same lead, and H_{tr} represent both direct and spin-assisted tunneling between leads. In the case of fully polarized leads, the summation in the Hamiltonian is taken only over one type of electron spin. This Hamiltonian is similar to the one describing a quantum dot studied by in Refs. [3, 4].

Since spin operators are not appropriate for diagrammatic techniques (as Wick's theorem does not apply to them), we represent spin-1/2 operators S_μ in terms of three Majorana fermions η_μ, $\mu \in \{x, y, z\}$ satisfying $S_\mu = -(i/2)\epsilon_{\mu\nu\lambda}\eta_\nu\eta_\lambda$, $(\eta_\mu)^\dagger = \eta_\mu$ and $\{\eta_\mu, \eta_\nu\} = \delta_{\mu\nu}$. The Hilbert space for Majorana fermions is eight-dimensional but the representation of spin operator reduces in it, and the four resulting twodimensional irreducible representations of spin are equivalent. Hence, one can restrict the calculations to one twodimensional subspace in the Majorana fermion Hilbert space that represents the twodimensional Hilbert space spanned over the two eigenstates of the operator S_z.

It is assumed that for sufficiently long times the composite (spin-electrons) system reaches a dc (non-equilibrium) steady state (which does not depend upon the initial conditions) and we use Keldysh diagrammatic technique [9] to describe it. The application of this standard method used in non-equilibrium phenomena [10] to tunneling via spin problem is described in Ref.[2]. The time-ordered averages in the Keldysh technique are taken with help of the evolution operator S_C along a closed time contour (its part running from $-\infty$ to $+\infty$ denoted by + and the part from $+\infty$ to $-\infty$ by − (refering to later times)). One defines the four (non-independent) real-time Green's functions $G^{mn}_\psi(t, t') = -i\langle\langle \mathcal{T} S_C \psi(t_m)\psi^\dagger(t'_n)\rangle\rangle$ where $m, n \in \{+, -\}$, $\mathcal{T}$ is the time-ordering operator, and the angular brackets denote averaging over the (pure or mixed) state of the interacting system. These four Green's functions (the "± basis") can be expressed in terms of the retarded $G^R_\psi(t, t')$, advanced $G^A_\psi(t, t') = (G^R_\psi(t, t'))^*$ and Keldysh Green's function $G^K_\psi(t, t') = -i\langle[\psi(t), \psi^\dagger(t')]\rangle$ ("Larkin-Ovchinnikov (LO) basis"). In thermal equilibrium the fluctuation-dissipation theorem gives

$$G^K_\psi(\omega) = h_{\text{eq}}(\omega)(G^A_\psi(\omega) - G^R_\psi(\omega)), \qquad (2)$$

with the thermal distribution function of fermions $h_{\text{eq}}(\omega) = -\tanh(\omega/2T)$. As in Ref.[2] we assume that for sufficiently long times the spin reaches a dc (non-equilibrium) steady state which does not depend on the initial conditions. In a dc steady state correlation functions depend on the time difference and here $h(\omega)$ differs from $h_{\text{eq}}(\omega)$ and it should be determined as a stationary solution of the quantum kinetic equation.

The bare Green's functions for electrons in the presence of applied magnetic field (or internal exchange field in the case of magnetically ordered leads) $\mathbf{B}_\alpha$ which induces spin polarization of electrons in the directions $\mathbf{m}_\alpha$. Since the leads are coupled to a thermal bath, the bare lead-electron Green's functions are characteristic for thermal equilibrium. They have the following matrix form in the electron spin space

$$\begin{aligned}
\mathbf{G}_\alpha(\omega) &= G_{\alpha 1}(\omega)\mathbf{I} + G_{\alpha 2}(\omega)\mathbf{m}\cdot\vec{\sigma}, \qquad (3)\\
G^R_{0\alpha,\sigma}(\omega) &= \int d\epsilon\, \frac{\rho_{\alpha,\sigma}(\epsilon)}{\omega - \epsilon + i0^+}, \quad G_{\alpha 1,2} = \frac{G_{\alpha,+} \pm G_{\alpha,-}}{2},\\
G^K_{0\alpha,\sigma}(\omega) &= 2\pi i\, h_{\text{eq}}(\omega - \mu_\alpha)\, \boldsymbol{\rho}_{\alpha,\sigma}(\omega),
\end{aligned}$$

where $\sigma = \pm$ and $\rho_{\alpha,\sigma}(\epsilon) = \rho(\epsilon - \sigma B_c - \mu_\alpha)$, and ρ is the bare density of states which we assume to be similar in both leads. We consider here only local Green's functions, so the momentum index can be dropped. For fully polarized electrons one has $G_{\alpha 1}(\omega) = \pm G_{\alpha 2}(\omega)$.

The bare propagators for η_μ satisfy $G^0_{xx} = G^0_{yy}$, $G^0_{yx} = -G^0_{xy}$, $G^0_{xz} = G^0_{zx} = G^0_{yz} = G^0_{zy} = 0$, due to the symmetry under rotations about the z-axis. A useful property holds for the Keldysh components of the spin Green's functions: since Majorana fermion operators are Hermitian, it turns out from the properties of the commutator that $G^K_{\mu\nu}(t, t') = -G^K_{\nu\mu}(t', t)$. In the frequency ω representation, for time-translation invariant solutions, this means that $G^K_{\mu\nu}(\omega) = -G^K_{\nu\mu}(-\omega)$, and therefore $G^K_{\mu\mu}(\omega)$ is an odd function of ω.

To write down Dyson's equation, it is convenient to use a basis in which the bare propagator is diagonal. Defining the canonical fermion operators $f = (\eta_x + i\eta_y)/\sqrt{2}$ and $f^\dagger = (\eta_x - i\eta_y)/\sqrt{2}$ with the associated Green's functions

$G_{ff^\dagger}(t) = -i\langle\langle f(t)f^\dagger(0)\rangle\rangle$ and $G_{f^\dagger f} = -i\langle\langle f^\dagger(t)f(0)\rangle\rangle$ one has that the bare propagators in the $(f, f^\dagger, \eta^z)$ basis in the steady state in the lowest order perturbation theory are given by

$$G^{0R}_{zz}(\omega) = \frac{1}{\omega + is},\ G^{0R}_{ff^\dagger, f^\dagger f}(\omega) = \frac{1}{\omega \mp B + is}, \qquad (4)$$

$G^{0K}_{zz}(\omega) = h_z(\omega)(G^{0A}_{zz}(\omega) - G^{0R}_{zz}(\omega))$ and analogous relations hold for $G^{0K}_{ff^\dagger}$, $G^{0K}_{f^\dagger f}$, where $s \to 0$ is a small regulator (a width due to an infinitesimally small coupling to a thermal bath). In the thermal equilibrium state $h_z(\omega), h_f(\omega), h_{f^\dagger}(\omega) = -\tanh(\omega/2T)$.

The vertices for interaction of Majorana fermions with electrons follow from the form of the tunneling Hamiltonian (1)

$$T_{\alpha\beta}\mathbf{S}\cdot\vec{\sigma} = \frac{T_{\alpha\beta}}{2}\sigma_z(f^\dagger f - ff^\dagger) + \frac{T_{\alpha\beta}}{\sqrt{2}}(f\sigma^- - f^\dagger\sigma^+)\eta_z,$$

and at each vertex of an internal point an additional $+$ or $-$ factor is assigned (due to the opposite direction of time integration for the points on the $-$ part of the Schwinger-Keldysh contour). Here $\sigma^\pm = (\sigma_x \pm i\sigma_y)$.

The complete Green's function satisfies Dyson's equation

$$G^{-1}(\omega) = G_0^{-1}(\omega) - \Sigma(\omega) \qquad (5)$$

for the 6x6 matrices (in the tensor product of the x, y, z space and the Keldysh space), where the free propagator is given by Eqs. (4). Since the bare propagator is diagonal in the basis (f, f^+, z), simple multiplication of block-diagonal matrices gives that to lowest order in Σ the complete Green's function in the 2x2 $G_{ff^\dagger}$ block depends only on $G^0_{ff^\dagger}$ and $\Sigma_{f^\dagger f}$. One then obtains the matrices $G^R_{f,f+}$, $G^A_{f,f+}$ and $G^K_{f,f+}$.

Parcollet and Hooley [2] have found that if one starts from the equlibrium distribution functions for Majorana fermions, the perturbation theory breaks down when one takes the limit of zero coupling to the leads prior to taking the limit of zero coupling to the thermal bath. One can, however, use the perturbation theory built on the appropriate bare Green's functions with a correct zeroth-order distribution functions $h_z(\omega)$, and $h_{f,f^\dagger}(\omega)$ which are stable with respect to weak perturbation and should be obtained self-consistently . Hence, assuming that after a long time the system is in a steady state, we can define a distribution function $h_f(\omega)$ obeying the self-consistency equation (with h_z, $h_{f^\dagger}$ defined analogously)

$$G^K_{ff^\dagger}(\omega) = h_f(\omega)(G^A_{ff^\dagger}(\omega) - G^R_{ff^\dagger}(\omega)), \qquad (6)$$

$$h_f(\omega) = \frac{\Sigma^K_{f^\dagger f}(\omega)}{\Sigma^A_{f^\dagger f}(\omega) - \Sigma^R_{f^\dagger f}(\omega)}. \qquad (7)$$

A useful relation $h_f(\omega) = -h_{f^+}(-\omega)$ can be derived using the definitions of Green's functions. One therefore needs to take the zeroth order Green's functions for spin in the form Eq. (4), calculate the first order self-energies $\Sigma^K_{f^\dagger f}(\omega)$ and $\Sigma^A_{f^\dagger f}(\omega) - \Sigma^R_{f^\dagger f}(\omega)$ and solve for h_f from Eq. (7).

The spin-assisted electron tunneling leads to a renormalization of the magnetic field acting on the spin [8] and the decoherence of spin precession. Diagrams shown in Fig. 1a,b lead

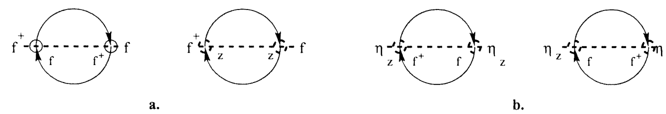

FIG. 1: Diagrams for the self-energy a) Σ_{ff+} and b) Σ_{zz} to second order in $T_{\alpha\beta}$: -the lead electron Green's functions and the spin Green's functions are represented by solid and dashed lines respectively; - the circles represent vertices between fields f, f^+ (solid) and f, η_z (dashed circle).

to the imaginary self-energies of Majorana fermions which depend on T, V, $\mathbf{m}_\alpha$ and $\mathbf{B}$. The expressions for them in the case of full polarization were obtained in [8]. The corresponding expressions in the weak polarization limit, $B_\alpha \ll D$, are proportional to the small parameter $(B_\alpha/D)^2$ (as obtained assuming that near the Fermi energy the electron density of states may be approximated as $\rho(\epsilon) \approx \rho_0(1 + \epsilon/D)$) and have the same form as results for the case of full polarization if one puts $\mathbf{m}_\alpha = 0$ and multiplies by factor 4 to account for both projections of electron spin.

Then the self-consistent equation (7) for the distribution function $h_f(B)$ gives ($b = B/T$, $v = V/T$)

$$h_f(B) = -\frac{2\mathsf{b}(1 - m_{\mathrm{R}z}m_{\mathrm{L}z}) + 2\mathsf{v}(m_{\mathrm{R}z} - m_{\mathrm{L}z}) + \mathsf{b}\,\theta}{\phi^+(1 - m_{\mathrm{R}z}m_{\mathrm{L}z}) + \phi^-(m_{\mathrm{R}z} - m_{\mathrm{L}z}) + \phi(\mathsf{b})\,\theta},$$

with $\theta = [|T_{RR}|^2(1 - m^2_{Rz}) + |T_{LL}|^2(1 - m^2_{Lz})]/|T_{LR}|^2$ and $\phi^\pm = \phi(\mathsf{v}+\mathsf{b}) \pm \phi(\mathsf{v}-\mathsf{b})$, $\phi(x) \equiv x\coth(x/2)$, provided that the denominator is nonzero (otherwise the distribution function is not determined).

The spectral density of the spin-spin correlation function for the S_x component is expressed as $S^{full}_{xx}(\omega) = \frac{1}{2}\int_{-\infty}^{\infty} dt e^{i\omega t}\langle S_x(t)S_x(0) + S_x(0)S_x(t)\rangle$. We will be interested in the part of this correlation function that exhibits a peak at the Larmor frequency. It can be shown that the expression for this part is

$$S_{xx}(\omega) = \frac{1}{8}\int\frac{d\epsilon}{2\pi}(G^A_{yy}(\epsilon + \omega)G^R_{zz}(\epsilon) + G^R_{yy}(\epsilon + \omega)G^A_{zz}(\epsilon)) + (\omega \to -\omega). \qquad (8)$$

The integration in Eq. (8) can be performed in the weak measurements limit $|T_{RL}|^2\rho^2 \ll 1$. We obtain that

$$S_{xx}(\omega) = \frac{1}{4}\left[\frac{\Gamma_t}{(\omega - B)^2 + \Gamma_t^2} + \frac{\Gamma_t}{(\omega + B)^2 + \Gamma_t^2}\right], \qquad (9)$$

where

$$\Gamma_t(B, V, T) = \mathrm{Im}[\Sigma^R_{f^\dagger f}(B) + \Sigma^R_{zz}(0)] \qquad (10)$$
$$= (\pi/4)T\rho_0^2\{(T^{(\mathrm{ex})}_{\mathrm{RL}})^2[2(1 + 2m_{\mathrm{R}z}m_{\mathrm{L}z} - \mathbf{m}_{\mathrm{R}}\cdot\mathbf{m}_{\mathrm{L}})\phi(\mathsf{v}) + (1 - m_{\mathrm{R}z}m_{\mathrm{L}z})(3\phi^+ + 4\mathsf{b}h_f(B)) + (m_{\mathrm{R}z} - m_{\mathrm{L}z}) \times (3\phi^- + 4\mathsf{b}h_f(B))] + \sum_\alpha T^2_{\alpha\alpha}[2(1 + 2m^2_{\alpha z} - |\mathbf{m}_\alpha|^2) + (1 - m^2_{\alpha z})(3\phi(\mathsf{b}) + 2\mathsf{b}h_f(B))]\}.$$

Thus we see that that the spin-spin correlation function of the x-component exhibits oscillations at Larmor frequency, with

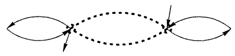

FIG. 2: Diagrams contributing to the peak at Larmor frequency in the current-current correlation function.

the peak depending on V, B and T. The same result follows for $S_{yy}(\omega)$.

The spectral density of the current-current correlation function is defined as $S_I(\omega) = \frac{1}{2}\int_{-\infty}^{\infty} dt e^{i\omega t}\langle \hat{I}(t)\hat{I}(0) + \hat{I}(0)\hat{I}(t)\rangle$. The current operator is given as $\langle \hat{I}\rangle = I_0 + \mathbf{I}_s \cdot \langle \mathbf{S}\rangle$ where $I_0 = 4\pi|T_0|^2\rho_0^2 V$ and for fully polarized electrons $\mathbf{I}_s = 4\pi|T_0 T_{\rm RL}|\rho_0^2 V(\mathbf{m}_{\rm R} + \mathbf{m}_{\rm L})$, whereas for weakly polarized electrons we get that I_s is small as $I_s \propto 4\pi|T_0 T_{\rm RL}|\rho_0^2 V(B_e/D)$. The spectral density of the current-current correlator is

$$S_I(\omega) = S_0 \coth(\frac{eV}{2T}) + S_{Is}\frac{\Gamma_t B^2 \sum_{a=x,y} S_{aa}(m_{\rm Ra} + m_{\rm La})^2}{(\omega^2 - B^2)^2 + 4\Gamma_t^2 B^2}$$

Here the first term is the usual shot-noise, $S_0 = eI$, whereas the second one describes a peak at Larmor frequency (the type of diagrams contributing to it is shown in Fig. (2)). For fully polarized electrons, $S_{Is} = (2\pi|T_0 T_{\rm RL}|\rho_0^2 eV)^2$, whereas in the case of weak polarization we obtain a similar expression but multiplied by a small factor of order $(B_e/D)^2$.

Hence we obtain that the spin precession cannot be seen in the current-current correlation function in the case when the lead electron polarization is parallel to $\mathbf{B}$, in spite of the fact that oscillations exist in the spin-spin correlation functions for the S_x and S_y components. The reason for this is that to observe the oscillations in $S_I(\omega)$ one needs to probe a component of spin perpendicular to $\mathbf{B}$, with the current, that is, the electrons need to have a component of polarization $\mathbf{m}_{\rm R}$ or $\mathbf{m}_{\rm L}$ that couples to S_x or S_y. Another important consequence is that for a weak polarization the spin-dependent contribution $S_{Is}(\omega)$ to the current-current correlation function is very small, proportional to $(B_e/D)^2$.

The signal-to-noise ratio at the peak is

$$R = 8{\rm v}\{2{\rm b}h_f(B)(2+\theta) + 3[\phi^+ + \phi({\rm b})\theta]\}^{-1}.$$

Here we see that the signal-to-noise ratio is reduced by the existence of electron reflection couplings as expected, since this process does not contribute to the current but does affect the width of the peak. The signal-to-noise ratio R in our approach reaches maximum, $R = 4(1 + \tilde{T}^2/2|T_{RL}|^2)^{-1}$ at $V = B$ at $V, B \gg T$, dropping to 4/3 in the limit $V \gg B$. The relaxation rate Γ_t is V-independent at $V < B$ and for $V > B$ the peak width increases linearly with V. Hence, R shows broad resonance when electron energy V coincides with Zeeman splitting (in a kind of a resonance tunneling). The broadening of the resonance is caused by variations in the electron energy change in tunneling due to broad electron bands inside electrodes.

In the quasiclassical approach, the effect of tunneling electrons on spin is replaced by action of the classical effective magnetic field with white noise spectrum of amplitude proportional to $I_0(|T_{RL}|^2/T_0^2)$. This effective field leads to a decoherence rate that is proportional to I_0 and does not depend on B and V in the limit $V \gg T$ (except that $I_0 \propto V$). In such an approach R increases with V and reaches the maximum value 4 in the limit $V \gg B$. Information on the energy spectrum of tunneling electrons is missed in the quasiclassical approch and it fails to describe correct dependence on V, B missing the resonance at $V = B$.

At $V, B \gg T$ for $\mathbf{m}_R = \mathbf{m}_L \perp \mathbf{B}$ the width of the peak at ω_L is $\Gamma = (\pi/4)\rho^2|T_{RL}|^2[6V - 4B\theta + \theta(3B - 2B\theta')]$ at $V > B$ and $\Gamma = (\pi/2)\rho^2|T_{RL}|^2(1+\theta/2)B$ at $B > V$, where $\theta' = (1+\theta/2)/(V/B+\theta/2)$. At high temperatures $T \gg V, B$ the width of the peak is determined by T, $\Gamma_t \propto T$. The peak at Larmor frequency is suppressed in the Zeno regime, when Γ_t exceeds B. This occurs for currents $I_0 > eB(T_0^2/T_{RL}^2)$. We predict that width increases linearly with V at $V > B$.

The important assumption made above is that the coupling of the spin to tunneling electrons is stronger than the one to the environment. This assumption is supported by the value of R of order unity. If Γ is dominated by relaxation due to the environment Γ_e, i.e. $\Gamma_e \gg \Gamma_t$ we would have $R \approx I_s^2/eI_0\Gamma_e \approx \Gamma_t/\Gamma_e \ll 1$. Our assumption may be checked also by studying the voltage dependence of the peak width: a linear dependence is expected for the tunneling-dominated decoherence.

Let us compare our theoretical results with the experimental data of Manassen et al. [5] and Durkan and Welland [6]. In these experiments polarization of electrodes was weak, while R was of order unity. This is in contradiction with our results.

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