Dimensional Control of Antilocalization and Spin Relaxation in Quantum Wires

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(Received 12 May 2006; published 27 April 2007)

The weak localization correction to the conductivity of quantum wires with linear Rashba-Dresselhaus spin-orbit coupling is derived analytically as function of wire width W. The spin relaxation rate is found to decrease as W becomes smaller than the spin-precession length L_{SO}. As a result, the sign of the conductivity correction switches to weak localization, positive magnetoconductivity for wire widths smaller than L_{SO}. A relaxation rate due to the cubic Dresselhaus coupling γ with a corresponding length scale L_γ remains, however, even in narrow wires $W \ll L_{SO}$. At low temperature, an antilocalization peak with negative magnetoconductivity is therefore recovered when the dephasing length exceeds L_γ.

DOI: 10.1103/PhysRevLett.98.176808 PACS numbers: 73.21.Hb, 72.10.Fk, 72.15.Rn, 73.20.Fz

Quantum interference of electrons in low-dimensional, disordered conductors results in corrections to the electrical conductivity $\Delta\sigma$. This weak localization (WL) effect is a very sensitive tool to study dephasing and symmetry breaking mechanisms in conductors [1]. The entanglement of spin and charge by spin-orbit interaction (SOI) reverses the effect to weak antilocalization (WAL), enhancing the conductivity. Because electron momenta are randomized by disorder, SOI results in randomization of their spin, Dyakonov-Perel spin relaxation (DPSR), with a rate of $1/\tau_s$ [2]. Since this mechanism breaks down in single channel wires whose width W is of the order of the Fermi wavelength λ_F [3,4], one may ask if spin relaxation is suppressed already in wider wires. As DPSR is caused by elastic momentum scattering, one could expect that it is suppressed in ballistic wires, when the elastic mean free path l_e exceeds W. In this Letter, we show, for any combination of linear Rashba and Dresselhaus SOI, that $1/\tau_s$ is strongly reduced in even wider wires: as soon as wire width W is smaller than the bulk spin-precession length L_{SO}. This explains the reduction of spin relaxation in quantum wires, recently observed with optical [5] and WL measurements [6]. L_{SO} can be as large as several μm and exceed both l_e and λ_F. In clean, ballistic 2D electron systems (2DES), L_{SO} is the length on which the electron spin precesses a full cycle. This length scale is not changed as wire width W is reduced below L_{SO}: the SOI itself remains unchanged, as long as there are many transverse channels. Therefore, while the spin of conduction electrons relaxes only on the enhanced length scale $L_s(W) = \sqrt{D\tau_s}$ ($D = v_F^2\tau/2$ diffusion constant, v_F, Fermi velocity), the spin can precess coherently as it moves along the wire on the length scale L_{SO}. Thus, this dimensional reduction of spin relaxation rate $1/\tau_s(W)$ can be very useful for the realization of spintronic devices, which rely on coherent spin evolution [7,8].

Hikami, Larkin, and Nagaoka [9] derived WAL for conductors with impurities of heavy elements. As conduction electrons scatter from them, the SOI randomizes their spin. The resulting spin relaxation suppresses interference of time reversed paths in spin triplet configurations, while singlet interference remains unaffected. The remaining singlet interference reduces the return probability, resulting in WAL. Weak magnetic fields suppress the singlet contributions, yielding negative magnetoconductivity. If the host lattice provides SOI, the conductivity has to be calculated in the basis of eigenstates of the Hamiltonian with SOI,

$$H_0 = (\hbar^2/2m_e)\mathbf{k}^2 + \hbar\boldsymbol{\sigma}\boldsymbol{\Omega} \tag{1}$$

(m_e, effective electron mass), $\boldsymbol{\Omega}^T = (\Omega_x, \Omega_y)$, are precession frequencies of the electron spin around the x- and y-axis. $\boldsymbol{\sigma}$ is a vector, with components σ_i, $i = x, y$, the Pauli matrices. The breaking of inversion symmetry causes a SOI, given by [10]

$$\boldsymbol{\Omega}_{\mathbf{D}} = \alpha_1(-\hat{e}_x k_x + \hat{e}_y k_y)/\hbar + \gamma(\hat{e}_x k_x k_y^2 - \hat{e}_y k_y k_x^2)/\hbar. \tag{2}$$

$\alpha_1 = \gamma\langle k_z^2\rangle$, the linear Dresselhaus parameter, measures the strength of the term linear in momenta k_x, k_y in the plane of the 2DES. When $\langle k_z^2\rangle \sim 1/a^2 \geq k_F^2$ (a, thickness of the 2DES, k_F, Fermi wave number), that term exceeds the cubic Dresselhaus terms with coupling γ. Asymmetric confinement of the 2DES yields the Rashba term (α_2, Rashba parameter) [11],

$$\boldsymbol{\Omega}_{\mathbf{R}} = \alpha_2(\hat{e}_x k_y - \hat{e}_y k_x)/\hbar. \tag{3}$$

The quantum correction to the conductivity $\Delta\sigma$ arises from the fact that the quantum return probability to a given point $\mathbf{x_0}$ after a time t, $P(t)$, differs from the classical return probability, due to quantum interference. Therefore, $\Delta\sigma$ is proportional to a time integral over the quantum mechanical return probability $P(t) = \lambda_F^d(t)n(\mathbf{x_0}, t)$ (d, dimension of diffusion, n, electron density). For uncorrelated disorder potential, $V(x)$, with $\langle V\rangle = 0$ and $\langle V(x)V(x')\rangle = \delta(x - x')/2\pi\nu\tau$ [$\nu = m/(2\pi\hbar^2)$, average density of states per spin channel, τ, elastic mean free time], we can perform the disorder average. Going to momentum ($\mathbf{Q}$) and frequency (ω) representation, and summing up ladder diagrams to take into account the diffusive motion, yields the quantum correction to the static conductivity [9],

$$\Delta\sigma = -\frac{2e^2}{h}\frac{\hbar D}{\text{Vol.}}\sum_{\mathbf{Q}}\sum_{\alpha,\beta=\pm} C_{\alpha\beta\beta\alpha,\omega=0}(\mathbf{Q}), \tag{4}$$

where the sum is over momenta $\mathbf{Q}$, and spin indices α, $\beta = \pm$ of the time reversed paths, respectively, and the Cooperon propagator $\hat{C}$ is neglecting the Zeeman coupling for $\epsilon_F\tau \gg 1$ (ϵ_F, Fermi energy), given by

$$\hat{C}(\mathbf{Q})^{-1} = \frac{\hbar}{\tau} - \int\frac{d\Omega}{\Omega}\frac{\hbar/\tau}{1 + i\frac{\tau}{\hbar}\mathbf{v}(\hbar\mathbf{Q} + 2e\mathbf{A} + 2m_e\hat{a}\mathbf{S})}. \tag{5}$$

The integral is over all angles of velocity $\mathbf{v}$ on the Fermi surface (Ω, total angle; e, electron charge; $\mathbf{A}$, vector potential). $\mathbf{S}$ is the sum of the spin vectors of time reversed paths: $\mathbf{S} = (\boldsymbol{\sigma} + \boldsymbol{\sigma}')/2$, so that $\hat{C}$ is a 4 by 4 matrix in spin space. $\hat{a}$ is the 2 by 2 matrix

$$\hat{a} = \frac{1}{\hbar}\begin{pmatrix} -\alpha_1 + \gamma k_y^2 & -\alpha_2 \\ \alpha_2 & \alpha_1 - \gamma k_x^2 \end{pmatrix}. \tag{6}$$

In 2D, the angular integral is continuous from 0 to 2π, yielding to lowest order in $(\mathbf{Q} + 2e\mathbf{A} + 2m\hat{a}\mathbf{S})$,

$$\hat{C}(\mathbf{Q}) = \frac{\hbar}{D(\hbar\mathbf{Q} + 2e\mathbf{A} + 2e\mathbf{A}_{\mathbf{S}})^2 + H_\gamma}. \tag{7}$$

The effective vector potential due to spin-orbit interaction, $\mathbf{A}_{\mathbf{S}} = m_e\hat{\alpha}\mathbf{S}/2$, ($\hat{\alpha} = \langle\hat{a}\rangle$) couples to total spin $\mathbf{S}$. The cubic Dresselhaus coupling reduces the effect of the linear one to $\alpha_1 - m_e\gamma\epsilon_F/2$. Furthermore, it gives rise to the spin relaxation term in Eq. (7),

$$H_\gamma = D\frac{m_e^2\epsilon_F^2\gamma^2}{\hbar^2}(S_x^2 + S_y^2). \tag{8}$$

In the representation of singlet, $|S=0; m=0\rangle = (|+\rangle|-\rangle - |-\rangle|+\rangle)/\sqrt{2}$ and triplet states $|S=1; m=0,\pm\rangle = (|+\rangle|-\rangle + |-\rangle|+\rangle)/\sqrt{2}$, $|S=1, m=1\rangle = |+\rangle|+\rangle$, and $|S=1, m=-1\rangle$, the 4 by 4 matrix $\hat{C}$ decouples into a singlet and triplet sector. Thus, introducing in Eq. (4), $1 = \sum_{S,m}|S,m\rangle\langle S,m|$, the quantum conductivity becomes a sum of singlet and triplet terms,

$$\Delta\sigma = -2\frac{e^2}{h}\frac{\hbar D}{\text{Vol.}}\sum_{\mathbf{Q}}\Bigg(-\frac{\hbar}{D(\hbar\mathbf{Q} + 2e\mathbf{A})^2} + \sum_{m=0,\pm1}\langle S=1,m|\hat{C}(\mathbf{Q})|S=1,m\rangle\Bigg). \tag{9}$$

It remains to diagonalize $\hat{C}$. For general SOI and magnetic fields, this results in cumbersome expressions. Exact analytical solutions are known in special cases [12,13]. In 2D one can treat the magnetic field nonperturbatively, in the basis of Landau bands [9]. In wires with widths smaller than cyclotron length $k_F l_B^2$ (l_B, the magnetic length, defined by $Bl_B^2 = \hbar/e$), the Landau basis is not suitable. However, one can define a rate with which the magnetic field breaks time reversal invariance, $1/\tau_B$, since in a magnetic field, the electrons acquire a magnetic phase. Averaging over all closed paths, this rate cuts off the divergence in Eq. (9), arising at small wave vectors $\mathbf{Q}^2 < 1/D\tau_B$. In 2D systems, τ_B is the time an electron diffuses along a path enclosing one magnetic flux quantum, $D\tau_B = l_B^2$. In wires of finite width W, the area which an electron can enclose in a time τ_B is $W\sqrt{D\tau_B}$. Thus, $1/\tau_B = De^2W^2B^2/(3\hbar^2)$. For arbitrary magnetic field, one gets $1/\tau_B = D(2e)^2B^2\langle y^2\rangle/\hbar^2$ with the expectation value of the square of the transverse position $\langle y^2\rangle$. This yields $1/\tau_B = D/l_B^2[1 - 1/(1 + W^2/3l_B^2)]$. Thus, we can diagonalize the Cooperon propagator as given by Eq. (7) without magnetic field and add $1/\tau_B$ together with dephasing rate $1/\tau_\varphi$ to the denumerator of $\hat{C}(\mathbf{Q})$, when calculating the conductivity correction, Eq. (9). In 2D, the Cooperon propagator can be diagonalized for pure Rashba coupling $\alpha_1 = 0$, $\gamma = 0$, or pure Dresselhaus coupling $\alpha_2 = 0$ [12,13]. Keeping only Rashba coupling α_2, diagonalization yields the triplet Cooperon Eigenvalues,

$$E_{T0}/(D\hbar) = \mathbf{Q}^2 + Q_{\text{SO}}^2,$$
$$E_{T\pm}/(D\hbar) = \mathbf{Q}^2 + \frac{3}{2}Q_{\text{SO}}^2 \pm \frac{1}{2}Q_{\text{SO}}^2\sqrt{1 + 16\frac{\mathbf{Q}^2}{Q_{\text{SO}}^2}}, \tag{10}$$

where $Q_{\text{SO}} = 2m_e\alpha_2/\hbar^2$. If we use the approximation,

$$E_{T\pm}/(D\hbar) \approx (Q \pm Q_{\text{SO}})^2 + Q_{\text{SO}}^2/2, \tag{11}$$

which is plotted for comparison with the exact dispersion, Eq. (10) in Fig. 1, we can integrate analytically over the 2D momenta Q and get the 2D quantum correction

$$\Delta\sigma = -\frac{1}{2\pi}\ln\frac{H_\varphi}{H_\varphi + H_s} + \frac{1}{\pi}\ln\frac{H_\varphi + H_s/2}{H_\tau}, \tag{12}$$

in units of e^2/h. All parameters are rescaled to dimensions of magnetic fields: $H_\varphi = \hbar/(4eD\tau_\varphi)$, $H_\tau = \hbar/(4eD\tau)$, and spin relaxation field $H_s = \hbar/(4eD\tau_{Sxx})$ [13]. The 2D spin relaxation rate of one spin component is for pure Rashba coupling, $1/\tau_{Sxx} = 1/\tau_s = 2k_F^2\alpha_2^2\tau/\hbar^2$ [12,13], and is related to spin-orbit gap $\Delta_{\text{SO}} = \hbar v_F Q_{\text{SO}}$, by $1/\tau_s = (\Delta_{\text{SO}}/\hbar)^2\tau/d$. We see that the largest contribution to the weak localization correction in Eq. (9) does come from the smallest Cooperon eigenvalues. Therefore, the minima of the eigenvalues, seen in Fig. 1, do cut off the logarithmic

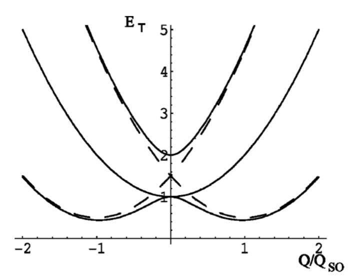

FIG. 1. Dispersion of triplet Cooperon modes in 2D in units of $\hbar DQ_{\text{SO}}^2$, Eq. (10) (solid lines) and Eq. (11) (dashed lines).

divergency, and we may call their values spin relaxation gaps, accordingly. These gaps are thus direct measures of the spin relaxation rate. We note that the two lowest minima of the triplet modes are shifted to nonzero wave vectors, $Q = \pm Q_{\rm SO}$. Thus, the spin relaxation gap is by a factor $1/2$ smaller than the eigenvalue at $Q = 0$ [12].

Without spin-orbit interaction, the WL of wires with width $W < L_\varphi$ is dominated by the transverse zero mode $Q_y = 0$. Integrating over the longitudinal momenta Q_x then yields the quasi-1D WL [14]. However, in the presence of SOI, setting simply $Q_y = 0$ is not correct. Rather, one has to solve the Cooperon equation with modified boundary conditions (BC) in transverse direction [4,15],

$$(-i\partial_y + 2eA_{Sy})C(x,y)|_{y=\pm W/2} = 0, \qquad (13)$$

for all x. Clearly, the transverse zero mode $Q_y = 0$ does not satisfy this BC. We therefore apply a non-Abelian gauge transformation to simplify this BC [15]. For quantum wires of length $L \gg L_\varphi$, we need a gauge transformation acting in transverse direction, only: $\hat{C} \to \tilde{C} = U\hat{C}\,\bar{U}$, with $U = \exp(i2eA_{Sy}y/\hbar)$, which simplifies the BC to, $-i\partial_y\tilde{C}(x,y)|_{y=\pm W/2} = 0$. For $W < L_\varphi$, transverse nonzero modes contribute terms to the conductivity which are only of order W/nL_φ, n integer. Therefore, it is sufficient to diagonalize the 0-mode expectation value of the gauge transformed inverse Cooperon propagator, $\tilde{H}_{\rm 1D} = \langle 0|\tilde{C}^{-1}|0\rangle$. Additional terms are created in $\tilde{H}_{\rm 1D}$ due to the non-Abelian nature of the transformation. We diagonalize $\tilde{H}_{\rm 1D}$, neglecting the small term due to cubic Dresselhaus coupling γ. We introduce the notation, $Q_{\rm SO}^2 = Q_D^2 + Q_R^2$, where Q_D depends on Dresselhaus SOI, $Q_D = m_e(2\alpha_1 - m_e\epsilon_F\gamma)/\hbar$. Q_R depends on Rashba coupling: $Q_R = 2m_e\alpha_2/\hbar$. We finally find the quasi-1D triplet eigenvalues,

$$\frac{E_{T0}}{\hbar D} = Q_x^2 + Q_{\rm SO}^2\delta_{\rm SO}^2\left(\frac{1}{2}t_{\rm SO}\delta_{\rm SO}^2 + 2c_{\rm SO}(1-\delta_{\rm SO}^2)\right),$$
$$\frac{E_{T\pm}}{\hbar D} = Q_x^2 + \frac{1}{4}Q_{\rm SO}^2\Bigg(4 - t_{\rm SO}\delta_{\rm SO}^4 - 4c_{\rm SO}\delta_{\rm SO}^2(1-\delta_{\rm SO}^2) \pm 2\sqrt{h(\delta_{\rm SO}) + \frac{16Q_x^2}{Q_{\rm SO}^2}[1 + c_{\rm SO}(c_{\rm SO}-2)\delta_{\rm SO}^2]}\Bigg), \qquad (14)$$

where $\delta_{\rm SO} = (Q_R^2 - Q_D^2)/Q_{\rm SO}^2$, and

$$c_{\rm SO} = 1 - \frac{2\sin(Q_{\rm SO}W/2)}{Q_{\rm SO}W}, \qquad t_{\rm SO} = 1 - \frac{\sin(Q_{\rm SO}W)}{Q_{\rm SO}W}. \qquad (15)$$

Here, $h(\delta_{\rm SO}) = t_{\rm SO}\delta_{\rm SO}^8/4 + \delta_{\rm SO}^2(1-\delta_{\rm SO}^2)[4c_{\rm SO}^2(1 - 3\delta_{\rm SO}^2 + 3\delta_{\rm SO}^4) + t_{\rm SO}^2\delta_{\rm SO}^2(1+\delta_{\rm SO}^2) - 6c_{\rm SO}t_{\rm SO}\delta_{\rm SO}^4]$. In Fig. 2, the gap of E_{T0} and the dispersion of the other two triplet modes are plotted for pure Rashba coupling $\delta_{\rm SO} = 1$, as function of width W, scaled with $Q_{\rm SO}$. In Fig. 3, the magnetoconductivity is plotted for pure Rashba coupling $\delta_{\rm SO} = 1$ as function of width W and magnetic field B.

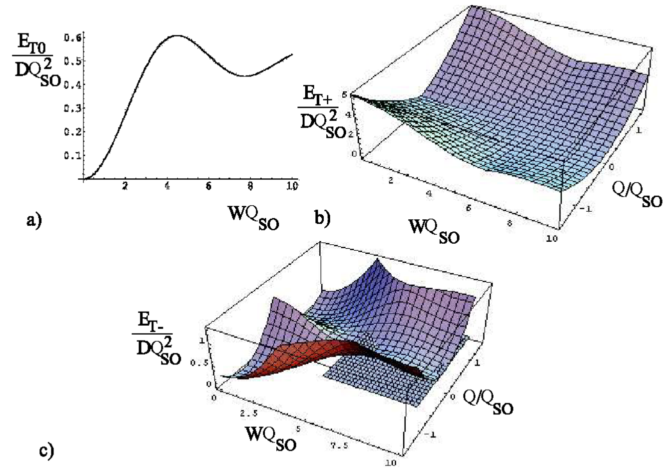

FIG. 2 (color online). For pure Rashba coupling $\delta_{\rm SO} = 1$: (a) Gap of Triplet mode E_{T0} as function of wire width W (in units of $L_{\rm SO} = 1/Q_{\rm SO}$). (b) Dispersion of Triplet mode E_{T+} and (c) of E_{T-} as function of W and momentum Q (scaled with $Q_{\rm SO}$) and $E/(\hbar DQ_{\rm SO}^2) = 1/2$ for comparison.

Inserting Eq. (14) into WL correction Eq. (9), the integral over momentum Q_x is done numerically. We note a change of sign from WAL to WL as $Q_{\rm SO}W$ becomes smaller than 1. In the crossover regime, $Q_{\rm SO}W \approx 1$ very weak magnetoconductivity is found. In the limit $WQ_{\rm SO} \gg 1$, the gaps of triplet mode dispersions, Eq. (14) coincide with 2D gap values $\hbar DQ_{\rm SO}^2(1/2, 1/2, 1)$ of Eqs. (10) (Note that spin quantization axis is rotated by the gauge transformation). For $WQ_{\rm SO} < 1$, the spin relaxation gap of triplet mode E_{T0} is to first order in $t_{\rm SO}$ and $c_{\rm SO}$: $\Delta_0 = DQ_{\rm SO}^2[2c_{\rm SO}\delta_{\rm SO}^2(1-\delta_{\rm SO}^2) + t_{\rm SO}\delta_{\rm SO}^4/2]$ and the gap of $E_{T\pm}$ is $\Delta_\pm = \Delta_0/2 + DQ_{\rm SO}^2(2c_{\rm SO} - t_{\rm SO}/2)\delta_{\rm SO}^4$. For $WQ_{\rm SO} \ll 1$, we can integrate over Q_x analytically, and get

$$\Delta\sigma = \frac{\sqrt{H_W}}{\sqrt{H_\varphi + B^*(W)/4}} - \frac{\sqrt{H_W}}{\sqrt{H_\varphi + B^*(W)/4 + H_s(W)}} - 2\frac{\sqrt{H_W}}{\sqrt{H_\varphi + B^*(W)/4 + H_s(W)/2}}, \qquad (16)$$

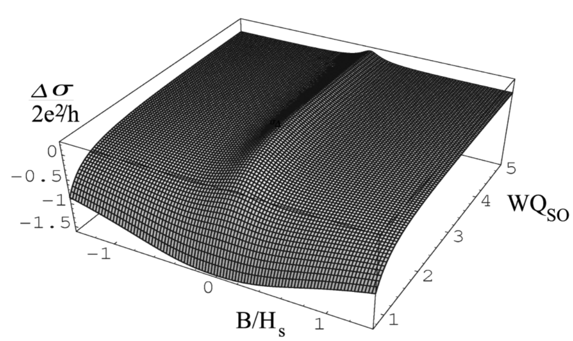

FIG. 3. The quantum conductivity correction in units of $2e^2/h$ as function of magnetic field B (scaled with bulk relaxation field H_s) and width W (scaled with spin-precession length $L_{\rm SO}$) for pure Rashba coupling, $\delta_{\rm SO} = 1$.

in units of e^2/h. We defined $H_W = \hbar/(4eW^2)$, and the effective external magnetic field,

$$B^*(W) = \left[1 - 1\Big/\left(1 + \frac{W^2}{3l_B^2}\right)\right]B. \tag{17}$$

The spin relaxation field $H_s(W)$ is for $W < L_{SO}$,

$$H_s(W) = \frac{1}{12}\left(\frac{W}{L_{SO}}\right)^2 \delta_{SO}^2 H_s. \tag{18}$$

The similarity to the effective magnetic field, Eq. (17), could be expected, since linear SOI enters the Cooperon, Eq. (7), via an effective magnetic vector potential [16]. Cubic Dresselhaus SOI gives rise to additional spin relaxation, Eq. (8), which has no analogy to magnetic field and is therefore not suppressed. When W is larger than spin-precession length L_{SO}, higher transverse modes become relevant, which may remove the oscillatory behavior of triplet eigenvalues as function of W seen in the zero mode approximation, Fig. 2 [17]. One can expect that in ballistic wires, $l_e > W$, the spin relaxation rate is suppressed in analogy to the flux cancellation effect, yielding the weaker rate, $1/\tau_s = (W/Cl_e)(DW^2/12L_{SO}^4)$, where $C = 10.8$ [18].

In conclusion, for wire widths W smaller than spin-precession length L_{SO}, spin relaxation due to linear Rashba and Dresselhaus SOI is suppressed. The spin relaxes then due to cubic SOI, only. The total spin relaxation rate as function of wire width is for $W < L_{SO}$,

$$\frac{1}{\tau_s}(W) = \frac{1}{12}\left(\frac{W}{L_{SO}}\right)^2 \delta_{SO}^2 \frac{1}{\tau_s} + D\frac{(m_e^2 \epsilon_F \gamma)^2}{\hbar^3}, \tag{19}$$

where $1/\tau_s = 2p_F^2[\alpha_2^2 + (\alpha_1 - m_e\gamma\epsilon_F/2)^2]\tau$ is the 2D spin relaxation rate. Using the analogy to a magnetic field, the enhancement of spin relaxation length $L_s = \sqrt{D\tau_s(W)}$ can be understood qualitatively: In a wire an electron covers by diffusion in time τ_s an area WL_s. Requiring that to be equal to L_{SO}^2 yields $1/L_s^2 \sim (W/L_{SO})^2/L_{SO}^2$, in agreement with Eq. (19). Reduction of spin relaxation has recently been observed in optical measurements of n-doped InGaAs quantum wires [5], where $\delta_{SO} \approx 1$, and in WL measurements in InGaAs, GaAs, and GaN wires [6]. Reference [5] reports saturation of spin relaxation in narrow wires, $W \ll L_{SO}$, attributed to cubic Dresselhaus coupling, in full agreement with Eq. (19). Thus, when dephasing length L_φ exceeds $L_\gamma = \hbar/m_e^2\epsilon_F\gamma$, a WAL peak should reappear at small magnetic fields, $l_B > L_\gamma$.

We thank V. L. Fal'ko and F. E. Meijer for stimulating discussions, I. Aleiner, K. Dittmer, C. Marcus, J. Ohe, T. Ohtsuki, K. Slevin, and A. Wirthmann for helpful discussions, and A. Chudnovskiy, K. Patton and E. Mucciolo for useful comments. We gratefully acknowledge the hospitality of MPIPKS in Dresden, the Physics Department of Sophia University, Tokyo, and Aspen Center for Physics. This work was supported by No. DFG-SFB508 B9.

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