Collective Excitation and Spectral Function of Quantum Wires in the Presence of Rashba Spin-Orbit interactions and a Parallel Magnetic Field

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The effects of the electron-electron interaction on the spin-filter proposed by Stˇreda and ˇSeba are studied. Using bosonization methods, we obtain the velocity of the collective mode and the spectral function exactly at low energy. These results can be experimentally verified by tunneling experiments.

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I. INTRODUCTION

Spintronics is a rapidly growing field, whose main objective is the active and versatile control over the spin degrees of freedom of electrons [1,2]. The devices based on spintronics have such merits as faster switching times and lower power consumption, to mention only a few, over conventional semiconductors devices based on charge degrees of freedom. The key point in the implementation of spintronic devices is the efficient generation of spin polarization. There exist diverse methods to generate spin polarization, and in this paper, we will be interested in one such methods that called spin filter [1]. Among many possible realizations of a spin filter, we focus on a mechanism that is closely related with the idea of spin field-effect transistor proposed by Datta and Das in 1990 [3]. The idea of a spin field-effect transistor hinges on the precession of the electron spin induced by Rashba spin-orbit interaction (SOI) [4,5]. The Rashba SOI originates from the confinement potential, which necessarily exists in all low-dimensional artificial semiconductor structures. Furthermore, the Rashba SOI can be controlled easily by changing the gate voltage, for example [6–9].

Stˇreda and ˇSeba have shown that by combining a strong Rashba SOI with a magnetic field parallel to the one-dimensional electron system (quantum wire), an energy interval in the energy spectrum emerges where the orientation of spin states is controlled by the direction of the electron velocity [10]. However, they did not study the effects of the electron-electron interaction, which are known to influence the physical properties of a one-dimensional electron system at low energy [11]. Recently, the interplay of the Rashba SOI and the electron-electron interaction in quantum wires has been studied by several groups [12,13]. However, none of these studies dealt with the case where the applied magnetic field was parallel to quantum wire as in the case of Stˇreda and ˇSeba.

In this paper, we study the influences of electron-electron interaction on the spin filter system proposed by Stˇreda and ˇSeba. We obtain the exact dispersion relation of collective mode and the exact electron spectral function at low energy by employing bosonization methods [11, 14]. The main results of this paper are Eqs. (30) and (36). In contrast to the ordinary Luttinger liquid where the collective excitations consist of two independent charge and spin modes, the charge and spin excitations are mixed for the collective excitations of the system we considered. The velocity of this collective excitation depends sensitively on the Rashba SOI and the magnetic field. The exact electron Green function explicitly exhibits how the electron-electron interaction modifies the spin-filtering mechanisms.

II. MODEL

The system we study is depicted in Fig. 1. The full Hamiltonian is

\[ \mathcal{H} = \mathcal{H}^{(0)} + \mathcal{H}_{\text{int}}, \]  

(1)
where $\mathcal{H}^{(0)}$ is the non-interacting Hamiltonian and $\mathcal{H}_{\text{int}}$ is the Hamiltonian for the electron-electron interaction which will be considered in Sec. III. The non-interacting Hamiltonian is given by (electron charge $= -e$)

$$\mathcal{H}^{(0)} = \sum_{\sigma = \uparrow, \downarrow} c_{\sigma} \mathcal{H}_{1} c_{\sigma},$$

$$\mathcal{H}_{1} = \frac{1}{2m^{*}} \left( \mathbf{p} + \frac{e}{c} \mathbf{A} \right)^{2} + W(\mathbf{r}) + \mathcal{H}_{R} + \mathcal{H}_{D} + \mathcal{H}_{Z},$$

where $c_{\sigma}$ is the electron operator and $m^{*}$ is the effective band mass. For InAs, the effective electron mass is $m^{*} \approx 0.024m_{e}$ [15,16]. $W(\mathbf{r}) = W(z)$ is the transverse confining potential enforcing quasi-one dimensional motion of electrons along the $x$-axis. The general spin-orbit interaction derived from the Dirac equation is of the form

$$H_{SO} = -\frac{1}{2} \mathbf{\tilde{\mu}} \cdot \mathbf{B}_{\text{eff}} = - \frac{|e|\hbar^{2}}{4m_{e}^{*}c^{2}} \mathbf{\tilde{\sigma}} \cdot \mathbf{\tilde{k}} \times \mathbf{\tilde{E}},$$

$$\mathbf{\tilde{\mu}} = \frac{-|e|\hbar}{m_{e}c} \mathbf{\tilde{S}}, \quad \mathbf{B}_{\text{eff}} = - \mathbf{\tilde{\sigma}} \times \mathbf{\tilde{E}}, \quad \mathbf{\tilde{E}} = \frac{\hbar \mathbf{\tilde{v}}}{m_{e}},$$

where $\mathbf{\tilde{\mu}}$ is the Thomas precession factor. The vector potential $\mathbf{A}$ describes the applied magnetic field parallel to the quantum wire along $x$-axis.

$$\mathbf{B} = -B\mathbf{\tilde{x}}, \quad \mathbf{A} = -By\mathbf{\tilde{z}}, \quad B > 0,$$

where we have chosen the Landau gauge. We will assume that the energy gap between subbands is sufficiently large so that the interband transition can be neglected. We will also assume that only the lowest subband is filled. Then, the matrix elements $p_{y}, p_{z} + eA_{z}/c$ can be neglected [3].

$\mathcal{H}_{R}$ is the Hamiltonian for the Rashba SOI. The Rashba electric field is given by $\mathbf{E}_{R} = E_{0}\mathbf{\hat{y}}$, $E_{0} > 0$. From Eq. (4),

$$\mathcal{H}_{R} = -\eta_{R} \left( \sigma_{z} k_{x} - \sigma_{x} k_{z} \right), \quad \eta_{R} = \frac{|e|\hbar E_{0}}{4m_{e}^{*}c^{2}} > 0,$$

where $m_{e}^{*}$ is the electron mass in vacuum. For InAs-based quantum wells, [15]

$$\eta_{R} = \alpha_{R}(eE_{0}), \quad \alpha_{R} \sim 110 \text{Å}^{2}.$$

The practical range of the Rashba electric field is $10^{4} - 10^{8}$ V/cm, so the parameter $\eta_{R}$ is in the range $10^{-8} - 10^{-10}$ eV/cm. In fact, there is no simple way to calculate $\eta_{R}$ because it depends both on the electric field inside the semiconductor heterostructure and on the detailed boundary conditions at the interface.

The bulk Hamiltonian of the Dresselhaus SOI is given by [17]

$$\mathcal{H}_{\text{bulk},D} = \gamma_{c} \left[ \sigma_{x} k_{x} \left( k_{y}^{2} - k_{z}^{2} \right) + \sigma_{y} k_{y} \left( k_{x}^{2} - k_{z}^{2} \right) + \sigma_{z} k_{z} \left( k_{x}^{2} - k_{y}^{2} \right) \right].$$

To obtain the effective Hamiltonian of the quantum wire, we have to take the average of the above bulk Hamiltonian with respect to the ground-state wave function of transverse $(y, z)$ degrees of freedom. In our geometry, the Rashba electric field is applied in the $y$-direction, and the lateral confining potential enforcing quasi one-dimensional motion is applied in the $z$-direction. Clearly, $(k_{z}) = 0$ because the subband wavefunction along the $z$-axis has even parity. The subband wavefunction along the $y$-axis is a real function: therefore, the expectation value $(k_{y}) = 0$, too. However, we have to note that $(y) \neq 0$ because the inversion symmetry is lacking in the $y$-direction. The effective Hamiltonian for the quantum wire is then

$$\mathcal{H}_{D} = \gamma_{c} \sigma_{x} k_{x} \left( k_{y}^{2} - \langle k_{y}^{2} \rangle \right) = \eta_{D} \sigma_{x} k_{x},$$

where $\eta_{D} = \gamma_{c} \left( \langle k_{y}^{2} \rangle - \langle k_{y}^{2} \rangle \right)$. Quantum wires can be tailor made such that $\eta_{D}$ becomes very small. Henceforth, we will assume this condition and consider only the Rashba SOI.

Finally, $\mathcal{H}_{Z}$ is the Zeeman coupling term:

$$\mathcal{H}_{Z} = -\mathbf{\mu} \cdot \mathbf{B} = -\frac{g_{0}}{2} \mu_{B} \sigma_{z} B \equiv -\epsilon_{z} \sigma_{z},$$

$$\epsilon_{z} = \frac{g_{0}}{2} \mu_{B} B > 0,$$

$g_{0}$ is the bulk $g$-factor. The $g$-factor of InAs is approximately 15 [15].

After the projection to the lowest subband, the one-dimensional one-particle Hamiltonian can be expressed as

$$\mathcal{H}_{1} = E_{0} + \frac{\hbar^{2} k^{2}}{2m^{*}} - \eta_{R} k \sigma_{z} - \epsilon_{z} \sigma_{z}.$$

$E_{0}$ is the energy at the bottom of the lowest 1D subband, which is taken to be zero. By the diagonalization of the Hamiltonian in Eq. (12), the energy eigenvalues and the corresponding normalized eigenvectors are obtained as follows:

$$E_{+}(k) = \frac{\hbar^{2}}{2m^{*}} k^{2} + \sqrt{\epsilon_{z}^{2} + \eta_{R}^{2} k^{2}},$$

$$E_{-}(k) = \frac{\hbar^{2}}{2m^{*}} k^{2} - \sqrt{\epsilon_{z}^{2} + \eta_{R}^{2} k^{2}},$$

$$\xi_{+} = \left( \frac{u_{k}^{+}}{v_{k}^{+}} \right), \quad \xi_{-} = \left( \frac{u_{k}^{-}}{v_{k}^{-}} \right),$$
where

$$u_k^+ = \frac{(-\eta_R k + D)}{\sqrt{(-\eta_R k + D)^2 + \epsilon_Z^2}},$$

$$v_k^- = \frac{-\epsilon_Z}{\sqrt{(-\eta_R k + D)^2 + \epsilon_Z^2}},$$

$$u_k^- = \frac{(-\eta_R k + D)}{\sqrt{(-\eta_R k + D)^2 + \epsilon_Z^2}},$$

$$v_k^+ = \frac{-\epsilon_Z}{\sqrt{(-\eta_R k + D)^2 + \epsilon_Z^2}},$$

(15)

with

$$D \equiv \sqrt{(\eta_R k)^2 + \epsilon_Z^2}. \quad (16)$$

$u_k^+$ and $v_k^+$ represent the amplitudes for the spin to point in the $+z$ and the $-z$ direction, respectively. The energy eigenvalues are depicted in Fig. 2. Fig. 3 clearly demonstrates that the spin of left-moving quasiparticles is mostly polarized in the $+z$ direction while that of right-moving quasiparticles is mostly polarized in the $-z$ direction. This is the crux of the antisymmetric spin filtering proposed by Středa and Seba [10]. Nevertheless, still substantial spin-mixing exists for moderate values of the Fermi momentum.

Let $b_k$ and $a_k$ be the quasiparticle operators corresponding to $E_+(k)$ and $E_-(k)$, respectively. They can be explicitly expressed in terms of electron operators as follows:

$$b^\dagger_k = c^\dagger_{k1}u_k^+ + c^\dagger_{k2}v_k^+, \quad a^\dagger_k = c^\dagger_{k1}u_k^- + c^\dagger_{k2}v_k^-,$$

$$c^\dagger_{k1} = b^\dagger_k v_k^- + a^\dagger_k u_k^-, \quad c^\dagger_{k2} = b^\dagger_k u_k^- + a^\dagger_k v_k^-.$$  

(17)

If the electrons are filled such that the Fermi energy is located in the energy gap between the $a$ and the $b$ bands at $k = 0$, we can safely neglect the $b$-type quasiparticles in the low-energy region. Then, Eq. (17) can be simplified to

$$a^\dagger_k = c^\dagger_{k1}u_k^- + c^\dagger_{k2}v_k^-,$$

$$c^\dagger_{k1} \sim a^\dagger_k u_k^-,$$  

$$c^\dagger_{k2} \sim a^\dagger_k v_k^-.$$  

(18)

III. ELECTRON-ELECTRON INTERACTION AND BOSONIZED HAMILTONIAN

A general electron-electron interaction in a quantum wire is given by

$$\mathcal{H}_{\text{int}} = \frac{1}{2} \sum_{k_1, k_2, q, \sigma} V_q c_{k_1, \sigma}^\dagger c_{k_2, \sigma}^\dagger c_{k_2, q, \sigma} c_{k_1, -q, \sigma},$$  

(19)

where $V_q$ is the interaction matrix element. For the long-range Coulomb interaction the interaction, matrix element is

$$V_q = \frac{2e^2}{\epsilon} K_0(|q|w) \rightarrow \frac{2e^2}{\epsilon} \ln \frac{1}{|q|w} \quad \text{for} \quad |q|w \ll 1. \quad (20)$$

$K_0$ is the modified Bessel function, and $w$ is the cutoff length scale, which is the order of the width of the quantum wire. $\epsilon$ is the bulk dielectric constant. For a short-range interaction, such as screened Coulomb interaction, the matrix element $V_q$ can be taken to be independent of the momentum transfer $q$. Projecting the Hamiltonian in Eq. (19) to the $a$-band with the use of Eq. (18), we obtain

$$\mathcal{H}_{\text{int}} = \frac{1}{2} \sum_{k_1, k_2, q} \langle k_1, k_2 | \hat{V} | k_1 - q, k_2 + q \rangle$$

$$\times a_{k_1}^\dagger a_{k_2} a_{k_2 - q} a_{k_1 - p},$$  

(21)

where

$$\langle k_1, k_2 | \hat{V} | k_1 - q, k_2 + q \rangle = V_p \langle \xi^\dagger_1 (k_1) \xi_-(k_1 - q) | \xi^\dagger_2 (k_2) \xi_-(k_2 + q) \rangle$$  

(22)

is the projected interaction matrix element in the low-energy Hilbert space. The explicit expression of the eigenvector $\xi_-$ is given by Eqs. (14) and (15). At low energy, only the electron states near the left and the right Fermi points need to be considered. Following the usual procedures of g-ology and the bosonization method [11],
we can express the interaction Hamiltonian in Eq. (21) within g-ology scheme. The details can be found in Ref. [18]. In this paper, a commensurate filling is not considered, so that the umklapp processes ($g_1$) can be neglected. We note further that for fermions of a single species (like α-quasiparticle here), the $g_1$ process is identical to the $g_2$ process [11].

It turns out that the effective low-energy Hamiltonian is given by [18]

$$\mathcal{H} = \pi v_F \int dx \left[ : \rho_R^2(\phi) : + : \rho_L^2(\phi) : \right]$$

$$+ \frac{1}{2N} \sum_q V_q \left[ \rho_R(q)\rho_R(-q) + \rho_L(-q)\rho_L(q) \right]$$

$$+ \frac{1}{N} \sum_q \left[ V_q - g V_{2k_F} \right] \rho_R(q)\rho_L(-q). \quad (23)$$

where $\rho_{R/L} = \sum_{\alpha} \psi_{R/L(\alpha)}^\dagger (p + q) \psi_{R/L(\alpha)} (p)$ is the density operator of right/left moving fermions, and $g = \frac{\epsilon_g^2}{\alpha_Z^2 + (\eta R k_F)^2}$. (24)

By introducing phase fields

$$\theta(x) = \frac{1}{2} \left[ \phi_R(x) + \phi_L(x) \right]$$

$$\phi(x) = \frac{1}{2} \left[ \phi_R(x) - \phi_L(x) \right], \quad (25)$$

where $\rho_{R/L}(x) = \frac{1}{2\pi} \partial_x \rho_{R/L}(x)$, the Hamiltonian in Eq. (23) can be immediately bosonized [11]. In momentum-frequency space,

$$S = \frac{1}{2\pi} \int \frac{d\omega dq}{(2\pi)^2} \left[ \theta(-q, -\omega) \phi(-q, -\omega) \right]$$

$$\times \left[ v_\theta(q) q^2 i \omega \quad v_\phi(q) q^2 \right] \left[ \theta(q, \omega) \phi(q, \omega) \right] \quad (26)$$

where

$$v_\theta(q) = v_\theta(q) = v_F \left( 1 + \frac{V_q}{\pi v_F} - g \frac{V_{2k_F}}{2\pi v_F} \right),$$

$$v_\phi(q) = v_\phi(q) = v_F \left( 1 + \frac{g V_{2k_F}}{2\pi v_F} \right). \quad (27)$$

The physical implication of the result in Eq. (27) are discussed below.

IV. PHYSICAL PROPERTIES

1. Collective Excitation

The dispersion relation of the collective excitation can be obtained from the kernel of the action in Eq. (26):

$$\det \begin{bmatrix} v_\theta(q) q^2 & i \omega \\ i \omega & v_\phi(q) q^2 \end{bmatrix} = 0 \quad (28)$$

After analytic continuation $i \omega \to \omega$, we find

$$\omega = \left[ v_\theta(q) v_\phi(q) \right]^{1/2} = v_0 q. \quad (29)$$

$v_0$ is the velocity of collective excitation. From Eq. (27), one can write

$$v_\omega = v_F \left[ 1 + \frac{V_q}{\pi v_F} + \frac{g V_{2k_F}/2}{(\pi v_F)^2} \right]^{1/2}. \quad (30)$$

The quantity in the bracket of Eq. (30) represents the renormalization effect due to the electron-electron interaction. The velocity of collective excitation can be controlled by band filling, the Rashba SOI, and the magnetic field through the dependences on $v_F$ and $g$.

2. Electron Spectral Function

The electron Green function in imaginary time is defined by

$$G_{\sigma \sigma'}(k, \tau) = -\langle T_\tau c_{\sigma}(\tau) c_{\sigma'}^\dagger (0) \rangle, \quad (31)$$

$T_\tau$ denotes the time ordering in imaginary time. At low energy the electron Green’s function can be expressed by using that of α-quasiparticle using the relation in Eq. (18):

$$G_{\sigma \sigma'}(k, \tau) = \xi_{-\sigma}(k) f_{-\sigma'}(k) G_{a}(k, \tau),$$

$$G_{a}(k, \tau) = -\langle T_\tau a(k, \tau) a_k^\dagger (0) \rangle. \quad \text{(32)}$$

In frequency space,

$$G_{\sigma \sigma'}(k, i \epsilon) = \xi_{-\sigma}(k) f_{-\sigma'}(k) G_{a}(k, i \epsilon). \quad \text{(33)}$$

On the other hand,

$$G_{a}(k, \tau) = \int dx e^{-ik_F x} G_{a}(x, \tau),$$

$$G_{a}(x, \tau) = -\langle T_\tau a(x, \tau) a(0, 0) \rangle. \quad \text{(34)}$$

Within the linearization approximation, one can write

$$a(x) = e^{ik_F x} \psi_R(x) + e^{-ik_F x} \psi_L(x). \quad \text{(35)}$$

By the standard bosonization formula [11]

$$\psi_R(x) = \frac{K_R}{\sqrt{2\pi(1/\Lambda)}} e^{i\phi + i\theta},$$

$$\psi_L(x) = \frac{K_L}{\sqrt{2\pi(1/\Lambda)}} e^{i\phi - i\theta},$$

where $1/\Lambda$ is a length scale of the order of the lattice spacing and $K_{R/L}$ are the Klein factors. Defining

$$D_{R/L}(x, \tau) = -(T_\tau \psi_R/L(x, \tau) \psi_R/L(0, 0)), \quad \text{(34)}$$

we can write

$$G_{a}(x, \tau) = e^{ik_F x} D_R(x, \tau) + e^{-ik_F x} D_L(x, \tau). \quad \text{(35)}$$
The Green’s functions \( D_{R/L}(x, \tau) \) can be computed exactly in a bosonization scheme. If the interaction matrix element is independent of the momentum transfer (as in Luttinger model),

\[
D_{R/L}(x, \tau) = -\frac{\text{sgn}(\tau)}{2\pi} \frac{1}{\tau_0 + i x} \left( \Lambda^2 ([\tau \tau_0]^2 + x^2) \right)^{\alpha/2},
\]

(36)

where \( \tau_0 = \sqrt{\tau_0^2 + \tau^2} \) is independent of momentum. The anomalous exponent \( \alpha \) is given by

\[
\alpha = \frac{1}{2} \left[ \sqrt{\frac{\tau_0^2}{\tau_0^2 + \tau^2}} + \sqrt{\frac{\tau_0^2}{\tau_0^2 + \tau^2} - 2} \right].
\]

(37)

Clearly this exponent can be also controlled by using the Rashba SOI and the magnetic field. In momentum-frequency space,

\[
D_{R/L}(k, i\epsilon) = \int dx d\tau e^{-i k x - i \epsilon \tau} D_{R/L}(x, \tau).
\]

(38)

If \( k \) is near the right Fermi point \( k_F \), the contribution from \( D_L \) to \( G_a \) can be neglected, so

\[
G_a(k, i\epsilon) \sim \int dx e^{-(i k - k_F) x} D_R(x, \tau),
\]

\[
= D_R(k - k_F, \epsilon), \quad |k - k_F| < \Lambda.
\]

(39)

The case where \( k \) is near the left Fermi point can be treated similarly.

The spectral function of \( a \)-quasi-particle at low energy is given by

\[
A^{(R/L)}_a(k, \epsilon) = -\text{Im} D_{R/L}(k \pm k_F, i\epsilon \rightarrow \epsilon + \delta).
\]

(40)

The Fourier transform of \( D_{R/L}(k, \tau) \) gives \[11, 14\]

\[
A^{(R)}_a(k, \epsilon) = \frac{c}{2v_0 \Lambda} \theta(u) \theta(-u')
\]

\[
\times \left( -u' \right)^{\alpha/2} \left( u \right)^{\alpha/2 - 1} \frac{\gamma}{\Lambda} e^{-\epsilon/v_0 \Lambda}
\]

\[
+ \frac{c}{2v_0 \Lambda} \theta(u) \theta(u') \left( u \right)^{\alpha/2 - 1} \left( -u' \right)^{\alpha/2} \frac{\gamma}{\Lambda} e^{\epsilon/v_0 \Lambda},
\]

\[
u = \frac{\eta - v_0 + \epsilon}{2v_0}, \quad u' = \frac{\eta - v_0 - \epsilon}{2v_0},
\]

\[q_+ = k - k_F, \quad c = \frac{1}{\Gamma(\alpha/2)\Gamma(\alpha/2 + 1)}.
\]

(41)

The left moving part is

\[
A^{(L)}_a(k, \epsilon) = \frac{c}{2v_0 \Lambda} \theta(u) \theta(u') \left( u \right)^{\alpha/2} \left( -u' \right)^{\alpha/2 - 1} \frac{\gamma}{\Lambda} e^{-\epsilon/v_0 \Lambda}
\]

\[
+ \frac{c}{2v_0 \Lambda} \theta(u) \theta(-u') \left( -u' \right)^{\alpha/2} \left( u \right)^{\alpha/2 - 1} \frac{\gamma}{\Lambda} e^{\epsilon/v_0 \Lambda},
\]

\[
u = \frac{\eta - v_0 + \epsilon}{2v_0}, \quad u' = \frac{\eta - v_0 - \epsilon}{2v_0}, \quad q_- = -k - k_F.
\]

(42)

From Eq. (40), the density of states of electrons can be extracted:

\[
N_{\sigma \sigma'}(\omega) = -\text{Im} \sum_k G_{\sigma \sigma'}(k, i\epsilon \rightarrow \omega + i\delta).
\]

(43)

The results in Eqs. (41) and (42) exhibit the typical power law dependences of Luttinger liquids. For ordinary Luttinger liquids the exponents of the power law are determined only by the electron-electron interaction, so it is difficult to control the exponents artificially. In contrast, the exponents of our spin-filter depends on external parameters such as Zeeman and Rashba coupling. Through the summation over \( k \), the spin-mixing matrix elements \( u_k \) and \( v_k \) influence the frequency dependence of density of states.

V. DISCUSSION AND CONCLUSIONS

There is a mechanism of spin filtering that employs the interplay of the Rashba effect and wave-number selectivity due to momentum-resolved tunneling between parallel quantum wires [19, 20]. Two parallel quantum wires of a very high quality can be fabricated, [21] which permits the momentum along quantum wire to be approximately conserved during the tunneling between them. These tunneling experiments can measure, in principle, the velocity of collective excitation and the momentum-dependent spectral function of a one-dimensional spin-filter. Through experiments, our theoretical predictions can be verified. In conclusion, we have calculated the velocity of collective excitation and the spectral function of a one-dimensional spin-filter by using a bosonization technique and found a non-trivial dependence on the Zeeman and Rashba coupling.

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