Physica E 75 (2016) 370–378

Contents lists available at ScienceDirect

Physica E

journal homepage: www.elsevier.com/locate/physe

Spin polarization induced by an electric field in the presence of weak localization effects

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HIGHLIGHTS

- We evaluate the quantum corrections to the Edelstein and spin Hall effects.
- The corrections to the spin Hall conductivity add up to zero as required by an exact identity.
- The corrections to the Edelstein conductivity are absorbed in the renormalized scattering time.

ARTICLE INFO

Article history:
Received 27 June 2015
Received in revised form
28 September 2015
Accepted 29 September 2015

Keywords:
Spin–orbit coupling
Weak localization
2DEG
Spintronics

ABSTRACT

We evaluate the spin polarization (Edelstein or inverse spin galvanic effect) and the spin Hall current induced by an applied electric field by including the weak localization corrections for a two-dimensional electron gas. We show that the weak localization effects yield logarithmic corrections to both the spin polarization conductivity relating the spin polarization and the electric field and to the spin Hall angle relating the spin and charge currents. The renormalization of both the spin polarization conductivity and the spin Hall angle combine to produce a zero correction to the total spin Hall conductivity as required by an exact identity. Suggestions for the experimental observation of the effect are given.

1. Introduction

Weak localization (WL) is the result of quantum interference corrections to the semiclassical theory of transport [1,2]. It manifests itself in good conductors as a negative or positive correction to the electrical conductivity depending on the symmetry properties of the system. The functional form varies with the effective dimensionality of the sample, behaving as a square root of temperature in three dimensions and logarithmically in two dimensions [3]. In the latter case, the resummation of the logarithmic correction via the renormalization group eventually leads to the Anderson localization transition in $d = 2 + \epsilon$ dimensions [4,5]. In the presence of spin–orbit coupling (SOC), the correction is positive and hence manifests as an antilocalizing behavior [6]. SOC affects WL because it yields a finite spin relaxation time, which introduces a cutoff in the logarithmic singularity associated with the so-called triplet channel of the particle-particle ladder, known as the Cooperon. Since the singlet and the triplet channels contribute to WL with opposite signs, the elimination of the triplet leaves the singlet alone, which then produces the antilocalizing behavior. In metallic conductors and doped semiconductors SOC was traditionally attributed to the electric field of impurities, which do not affect the nature of the electron eigenstates. In the last two decades, however, the two-dimensional electron gas (2DEG) has become one of the most analyzed model systems for electrical transport, due to the possibility of realizing it in semiconducting systems, and more recently at metallic [7] and oxides [8] interfaces. The realization of the 2DEG leads to the breaking of inversion symmetry with respect to the axis, say the z axis, perpendicular to the 2DEG plane, say the x and the y plane. In these circumstances, in the presence of a finite spin–orbit interaction, Bychkov and Rashba have proposed a model Hamiltonian [9], which, besides the standard effective-mass kinetic energy term, contains a spin–orbit coupling term linear in momentum

$$H = \frac{p^2}{2m} + \alpha(\sigma^x p_y - \sigma^y p_x), \tag{1}$$

where $\mathbf{p} = (p_x, p_y)$ is the vector of the components of the

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http://dx.doi.org/10.1016/j.physe.2015.09.046
1386-9477/

momentum operator, m is the effective mass and α a SOC constant with σ^x and σ^y the standard Pauli matrices. The Rashba Hamiltonian equation (1) has been extensively studied over the last 20 years, especially aiming at the development of new spintronic functionalities [10]. In this respect the spin Hall effect (SHE) [11–15] and the current-induced spin polarization effect [16,17] (known also as the Edelstein or inverse spin-galvanic effect) have been the focus of an intensive dedicated research (see [18] for a recent review). These effects, whose precise definition will be given later on, manifest due to the coupling of charge and spin degrees of freedom and hence introduce, besides the standard electrical conductivity, new transport parameters. These are defined as the linear coefficients relating the spin polarization and the spin current to the applied electric field

$$\langle\langle s^y\rangle\rangle = \sigma^{EC}E_x, \quad \langle\langle j_y^z\rangle\rangle = \sigma^{SHC}E_x, \tag{2}$$

where the double brackets indicate the quantum and statistical average. σ^{EC} and σ^{SHC} are referred to as the spin polarization or Edelstein and the spin Hall conductivities, respectively. As for the electrical conductivity, these transport parameters can be studied with the well-known impurity technique. One advantage of this technique [19], based on standard diagrammatic perturbation theory, is the appearance of the semiclassical Drude–Boltzmann theory of transport at the leading approximation in an expansion of the small parameter $\hbar/(\epsilon_F\tau)$, where ϵ_F and τ are the Fermi energy and the elastic scattering time, the only two parameters characterizing a disordered Fermi gas. In such an expansion, WL arises in the next-to-leading approximation in the expansion in $\hbar/(\epsilon_F\tau)$.

WL effects in the presence of the Rashba SOC described by Eq. (1) have been analyzed by several authors, most of the attention having been focused on the electrical conductivity only [20–25], with some works considering the spin conductivity as well [26,27]. It is worth noticing that the latter two works use semiclassical functional integration methods. It is the aim of the present work to extend this analysis to the other transport parameters mentioned above, whose experimental study has developed considerably in the last few years [28,7]. We find that σ^{EC} and the spin Hall angle $\gamma_{SH} = e\sigma^{SHC}_{drift}/\sigma_0$ acquire logarithmic corrections which can be absorbed in terms of the renormalization of the scattering time appearing in the electrical conductivity σ_0. We emphasize that σ^{SHC}_{drift} is not the full spin conductivity σ^{SHC} which would be measured in an experiment [29]. As will be shown in the next section, σ^{SHC} can be expressed in terms of σ^{EC} and σ^{SHC}_{drift}. The renormalizations of both σ^{EC} and γ_{SH} compensate in such a way that σ^{SHC} has no correction as expected on general arguments [30].

The plan of the paper is as follows. In the next section we introduce the disordered Rashba model and review the theory of σ^{EC} and σ^{SHC} to the leading order in the parameter $\hbar/(\epsilon_F\tau)$ within the impurity technique. This is necessary to prepare the ground for the following sections. Section 3 deals with the WL localization corrections in the presence of the Rashba SOC. The evaluation of the electrical conductivity is reviewed as an example. Section 4 presents the calculation of the WL corrections to σ^{EC} and σ^{SHC}. Section 5 provides a discussion of the results obtained, whereas technical points of the calculations are given in the appendices at the end of the paper. From now on, if not otherwise specified, we will work in natural units $\hbar = c = 1$.

2. The disordered Rashba two-dimensional electron gas at leading order in $1/(\epsilon_F\tau)$

In the presence of scattering from impurities, the 2DEG Hamiltonian of Eq. (1) acquires an additional random potential term $U(\mathbf{r})$ defined by the averages

$$\langle U(\mathbf{r})\rangle = 0, \quad \langle U(\mathbf{r})U(\mathbf{r}')\rangle = \frac{1}{2\pi N_0\tau}\delta(\mathbf{r}-\mathbf{r}') \tag{3}$$

where $\mathbf{r} = (x, y)$ and $\mathbf{r}' = (x', y')$ are the coordinate operators, $N_0 = m/2\pi$ the two-dimensional density of states and τ the elastic scattering time. At leading order in the expansion parameter $1/(\epsilon_F\tau)$, the self-energy is given by the self-consistent Born approximation

$$\Sigma^{R,A}(\mathbf{p}, \epsilon) = \frac{1}{2\pi N_0\tau}\sum_{\mathbf{p}'} G^{R,A}(\mathbf{p}', \epsilon), \tag{4}$$

where $G^{R,A}$ denotes the retarded and advanced Green functions. As discussed in [31,32], in the presence of Rashba SOC the Green function has a nontrivial structure in spin space, whereas the self-energy remains diagonal, $\Sigma = \Sigma^0\sigma^0$, $G = G^0\sigma^0 + G^1\sigma^1 + G^2\sigma^2$. Explicitly we have

$$G^0 = \frac{1}{2}(G_+ + G_-)$$
$$G^1 = \frac{\hat{p}_y}{2}(G_+ - G_-)$$
$$G^2 = -\frac{\hat{p}_x}{2}(G_+ - G_-)$$
$$G_\pm = \left(\epsilon + \mu - p^2/2m \mp \alpha p - \Sigma^0\right)^{-1}, \tag{5}$$

with

$$(\Sigma^0)^{R,A} = \mp\frac{i}{2\tau}. \tag{6}$$

The Edelstein (EC) and spin Hall (SHC) conductivities are defined in terms of the spin polarization and spin Hall current induced by an applied electric field taken along the x axis for definiteness's sake $E_x = -\partial_t A_x$. The corresponding Kubo formulae are

$$\sigma^{EC} = \lim_{\omega\to 0}\frac{\text{Im}\,\langle\langle s^y; j_x\rangle\rangle}{\omega}, \tag{7}$$

and

$$\sigma^{SHC} = \lim_{\omega\to 0}\frac{\text{Im}\,\langle\langle j_y^z; j_x\rangle\rangle}{\omega}, \tag{8}$$

where the bare vertices $s^y = \sigma^y/2$, $j_y^z = \sigma^z p_y/2m$ and $j_x = -e\hat{v}_x$, $\hat{v}_x(\mathbf{p}) = p_x/m - \alpha\sigma^y$ denote the operators for spin polarization, spin current and charge current, respectively. The evaluation of the response functions (7) and (8) involves the standard bubble diagrams of the Green function lines obtained by the self-consistent Born approximation (4) decorated by the insertion of the impurity ladder. This corresponds to the inclusion of the so-called vertex corrections, which lead to renormalized vertices [31].

The expression (7) for the EC becomes

$$\sigma^{EC} = -\frac{e}{2\pi}\sum_{\mathbf{p}}\text{Tr}\left[S^y G^R_{\mathbf{p}}\hat{v}_x(\mathbf{p})G^A_{\mathbf{p}}\right], \tag{9}$$

where the vertex renormalization can be attributed either to the left spin vertex or to the right current vertex and we have dropped the dependence on the frequency argument of the Green function. In the former case, by using the renormalized spin vertex indicated by a capital letter $S^y = ((1+x^2)/x^2)\sigma^y \equiv (\tau_{DP}/2\tau)\sigma^y = (\tau_{DP}/\tau)s^y$, one obtains [16]

$$\sigma_0^{EC} = -e\alpha N_0\tau, \tag{10}$$

where the subscript 0 in σ_0^{EC} indicates the lowest order in the parameter $1/(\epsilon_F\tau)$. We have defined the parameter $x = 2\alpha p_F\tau$ and introduced the D'yakonov–Perel relaxation time $\tau_{DP} = 2\tau((1+x^2)/x^2)$, p_F being the Fermi momentum in the absence