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A drift-diffusion model for spin-polarized transport in a two-dimensional non-degenerate electron gas controlled by spin–orbit interaction

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Abstract

We apply the Wigner function formalism to derive drift-diffusion transport equations for spin-polarized electrons in a III–V semiconductor single quantum well. The electron spin dynamics is controlled by the spin–orbit interaction which is linear in the momentum. In the transport regime studied, the electron momentum scattering rate is appreciably faster than the spin dynamics. A set of transport equations is defined in terms of a particle density, a spin density, and the respective fluxes. The model developed allows study of the coherent dynamics of a non-equilibrium spin polarization. As an example, we consider a stationary transport regime for a heterostructure grown along the (0, 0, 1) crystallographic direction. Due to the interplay of the Rashba and Dresselhaus spin–orbit terms, the spin dynamics strongly depends on the transport direction. The model is consistent with the results of pulse–probe measurements of the spin coherence in strained semiconductor layers. It can be useful in studying properties of spin-polarized transport and modelling spintronic devices operating in the diffusive transport regime.

1. Introduction

The spin-dependent properties of the electron transport in semiconductors have recently attracted significant attention from the scientific community in connection with the developing fields of magnetoelectronics and spintronics [1–5]. In comparison with magnetoelectronic devices utilizing giant magnetoresistance and tunnelling magnetoresistance effects in layered ferromagnetic metal structures [6], semiconductor spintronic devices promise to be more universal in application due to the possibility of adjusting the potential variation and spin polarization in the active region of the spin devices by means of external voltages and doping profiles [7–9]. Different designs of transistors and spin-filtering devices utilizing control of the spin polarization in semiconductor structures have been proposed [10–19].

It is of interest in spintronics to build a device that is tolerant of undesirable effects of the environment and working at room temperature [1, 2, 4, 5]. However, the crucial phenomenon for spintronic devices is the loss of the non-equilibrium spin polarization owing to spin–environment interactions. The functionality of most of the proposed devices is sensitive to the temperature, impurities, and internal and external fields. Detailed examination of a spin transport problem in semiconductor structures is required for modelling of realistic processes in such devices.

In this work we study spin-polarized transport in a two-dimensional non-degenerate electron gas (2DEG) in a III–V semiconductor heterostructure in the collision dominated regime [20], where the electron momentum scattering rate is appreciably faster than the spin dynamics. In general, this approximation is applicable for transport in heterostructures with a weak spin–orbit coupling at high temperature. For example, in GaAs/AlGaAs heterostructures for $T \sim 100$ K or higher the electron momentum scattering is mostly governed by the emission of polar optical phonons [21]. This is true even at lower temperatures if a moderate or strong electric field is applied. The characteristic scattering rate for this process is of the order of several 10^{12} s^{-1} [21]. The spin evolution in III–V semiconductors without external magnetic fields is mostly controlled by the spin–orbit interaction [20, 22]. It can be characterized by a spin precession frequency, Ω. According to recent measurements, for a Ga/AlGaAs heterostructure $\Omega \sim 10^{10}$–10^{11} s^{-1} [23, 24]. This shows that during one period of spin precession an electron experiences many collisions. The question is whether this spin transport regime can be useful for spintronics. Since the seminal proposal of a spin field effect transistor (spin-FET) by Datta and Das [10], utilization of the spin–orbit interaction in spintronic devices has remained an attractive idea [13–17]. A comprehensive review of the spin–orbit coupling effects for the purposes of spintronics can be found in [25]. For such devices, due to the electric field dependence of the spin–orbit coupling constants, a conventional electric gate can be used to control the coherent electron spin dynamics [10, 13–17]. However, the same spin–orbit interaction mechanism leads to spin dephasing due to the randomization of the electron momentum (the Dyakonov–Perel spin relaxation mechanism [20, 22, 26]). To avoid undesirable effects of this coupling, device models [10, 13–15] were proposed for the ballistic transport regime. In the recent work by Schliemann *et al* [16] it was shown that in some cases the spatial electron motion and the spin evolution can be decoupled owing to the symmetry of the Rashba [27] and Dresselhaus [28] spin–orbit interaction terms. The effect of momentum scattering on the spin coherent dynamics can be diminished and the spin-FET design [10] is applicable in the diffusive transport regime [16, 29]. Moreover, Schliemann *et al* have proposed a non-ballistic spin-FET [16], in which the spin dephasing is controlled by an external gate. Another design of a spin-FET utilizing only non-magnetic materials and operating with a spin dephasing rate in semiconductor heterostructures was proposed in [17]. These devices can be operational in the semiclassical diffusive transport regime that is of interest in this work.

The general drift-diffusion approach to spin-polarized transport is based on the two-state (spin-up and spin-down) model originally developed for ferromagnetic metals [30, 31]. The phenomenological model for non-collinear spin transport including the effects of spin–orbit coupling has been developed for the regime where the spin dynamics is significantly faster than the momentum scattering rate [32, 33]. Although these models can be useful for investigation of a broad class of transport problems in semiconductors [7–9, 32–34], they do not include the effects of a spin phase memory and are inapplicable in problems where quantum superposition of spin-up and spin-down states plays an important role [10]. The spin density matrix or spin polarization vector description of a spin state [35–37] is more appropriate for this case.

In our model we use the Wigner function representation for an electron with spin [38]. This approach was utilized before for different transport problems including the effects of

quantum potentials [39], quantum collisions [40], and electron transport in magnetic fields [41]. Recently, it has been applied to study spin-polarized electron transport in semiconductor heterostructures in the ballistic regime [42]. We consider the semiclassical transport regime where collisions with phonons and impurities control the transport properties. We show that in this model the Wigner function transport equation can be reduced to a set of drift-diffusion equations for a particle density, a particle current, a spin density, and a spin current.

2. Model

In most spintronic devices utilizing spin–orbit interaction in semiconductor heterostructures to control the spin dynamics [10, 13–17], electrons are confined by the effective potential in the direction orthogonal to the semiconductor interface and propagate in the plane of the heterostructure. The effective mass Hamiltonian for an in-plane electron motion in the one-subband approximation can be written as

$$\begin{aligned} H &= \frac{\mathbf{p}^2}{2m^*} + V(\mathbf{r}) + H_{\mathrm{SO}}, \\ H_{\mathrm{SO}} &= \mathbf{p}\mathbf{A}\boldsymbol{\sigma}/\hbar. \end{aligned} \tag{1}$$

It is assumed that the electron motion in the direction of quantization can be decoupled from the motion in the plane of a quantum well (QW) and that the electron kinetic energy is small in comparison with the subband splitting. The shape of the conduction band is assumed parabolic. The operators for the electron momentum, $\mathbf{p}$, and the spatial position, $\mathbf{r}$, are defined as two-dimensional vectors in the plane of the QW, while the spin operator, $\boldsymbol{\sigma}$, is a three-dimensional vector. The potential, $V(\mathbf{r})$, corresponds to the interaction with an electric field oriented in the plane of the QW. The spin–orbit interaction term, H_{SO}, is written in a general dyadic form linear in the electron momentum. This term is assumed small in comparison with other terms in the Hamiltonian, H. The matrix elements $A_{j\alpha}$ are constants of the spin–orbit interaction, coupling the jth component of the momentum with the αth component of the spin. Here, and in the following, we use Latin letters to index vector or matrix components in spatial dimensions and Greek letters to index components in the spin space. We set the z axis of the spatial coordinate system in the direction orthogonal to the QW plane, while the orientation of the spin coordinate system is left unspecified. An arbitrary rotation of the spin coordinate system will affect the form of the spin–orbit coupling matrix $\mathbf{A}$, but not the general representation of equation (1).

The quantum state of an electron with spin can be described by the density matrix operator, $\rho(\mathbf{r}, \mathbf{r}', s, s', t)$, which is dependent on two coordinate variables and two spin variables. After transformation to the new spatial representation,

$$\begin{aligned} \mathbf{R} &= (\mathbf{r} + \mathbf{r}')/2, \\ \Delta\mathbf{r} &= \mathbf{r} - \mathbf{r}', \end{aligned} \tag{2}$$

the equation for the density matrix will be

$$\begin{aligned} \mathrm{i}\hbar\frac{\partial\rho}{\partial t} = &-\frac{\hbar^2}{m^*}\sum_j \frac{\partial^2\rho}{\partial R_j \partial \Delta r_j} + (V(\mathbf{R} + \Delta\mathbf{r}/2) - V(\mathbf{R} - \Delta\mathbf{r}/2))\rho \\ &+ \frac{\mathrm{i}}{2}\sum_{j,\alpha} A_{j\alpha}\left\{\sigma_\alpha, \frac{\partial\rho}{\partial R_j}\right\} + \mathrm{i}\sum_{j,\alpha} A_{j\alpha}\left[\sigma_\alpha, \frac{\partial\rho}{\partial \Delta r_j}\right]. \end{aligned} \tag{3}$$

The effect of spin–orbit interaction is introduced by the last two terms in equation (3), where $[\boldsymbol{\sigma}, \ldots]$ and $\{\boldsymbol{\sigma}, \ldots\}$ are the commutator and anticommutator with the Pauli spin matrices

respectively. Following the standard transformation to the Wigner function [38],

$$W_{ss'}(\mathbf{R}, \mathbf{k}, t) = \int \rho(\mathbf{R}, \Delta\mathbf{r}, s, s', t)\mathrm{e}^{-\mathrm{i}\mathbf{k}\Delta\mathbf{r}}\,\mathrm{d}^2\Delta r, \tag{4}$$

and assuming that the potential $V(\mathbf{r})$ varies slowly and smoothly with the position $\mathbf{r}$, we obtain the transport equation for a single electron with spin:

$$\frac{\partial W}{\partial t} + \frac{1}{2}\left\{v_j, \frac{\partial W}{\partial x_j}\right\} - \frac{1}{\hbar}\frac{\partial V}{\partial x_j}\frac{\partial W}{\partial k_j} + \mathrm{i}k_j[v_j, W] = \mathrm{St}W. \tag{5}$$

On the right-hand side of equation (5) we have included the phenomenological scattering term, $\mathrm{St}W$, responsible for the interactions of an electron with phonons and non-magnetic impurities. Unlike [35, 43], we are interested in a transport regime where the electron–electron interaction produces small effects on the spin dynamics in comparison with the effects of the phonon and impurity scattering [20, 22].

In the spin space the velocity operator,

$$v_j = \frac{\partial H}{\partial p_j}, \tag{6}$$

and the Wigner function, W, are (2×2) matrices, while the potential, $V(\mathbf{r})$, and the electron wavevector, $\mathbf{k}$, are scalar variables. The last term on the left-hand side of equation (5) expresses the spin rotation. Matrix equation (5) can be projected onto the set of Pauli matrices, σ_α, and the unit matrix, I, using the following relations for the Wigner function [40]:

$$W = \tfrac{1}{2}(W_n I + W_{\sigma_\alpha}\sigma_\alpha), \tag{7}$$

and the velocity operator:

$$v_j = (v_n^j I + v_{\sigma_\alpha}^j \sigma_\alpha), \qquad v_n^j = \hbar k_j/m^*, \quad v_{\sigma_\alpha}^j = A_{j\alpha}/\hbar. \tag{8}$$

In the zeroth order of approximation of the spin–orbit coupling constant, $A_{j\alpha}$, scattering events do not couple different spin components of the Wigner function. The collision term possesses a semiclassical form:

$$\mathrm{St}W(\mathbf{R}, \mathbf{k}, t) = \int S(\mathbf{k}, \mathbf{k}')(W(\mathbf{R}, \mathbf{k}', t) - W(\mathbf{R}, \mathbf{k}, t))\,\mathrm{d}^2k', \tag{9}$$

where $S(\mathbf{k}, \mathbf{k}')$ is the transition rate for electrons without spin. We use the relaxation time approximation for equation (9) with the same set of assumptions as are usually applied for transport in III–V semiconductors [21]. Corrections to the collision term, $\mathrm{St}W$, linear in the spin–orbit interaction [44], mix spin-polarized components of the Wigner function, W_σ, with the non-polarized function, W_n, and produce the effect of the electron spin polarization by the in-plane electric field [36, 44–46]. We do not consider this effect owing to the assumption that the spin–orbit coupling is small in comparison with the electron kinetic energy.

To get drift-diffusion transport equations in terms of macroscopic variables we apply the moment expansion procedure [47] to equation (5). The (2×2) matrices of the particle density and current density at the position $\mathbf{R}$ are defined as

$$n(\mathbf{R}) = \int W\,\mathrm{d}^2k \tag{10}$$

and

$$J_j(\mathbf{R}) = \tfrac{1}{2}\int \{v_j, W\}\,\mathrm{d}^2k, \tag{11}$$

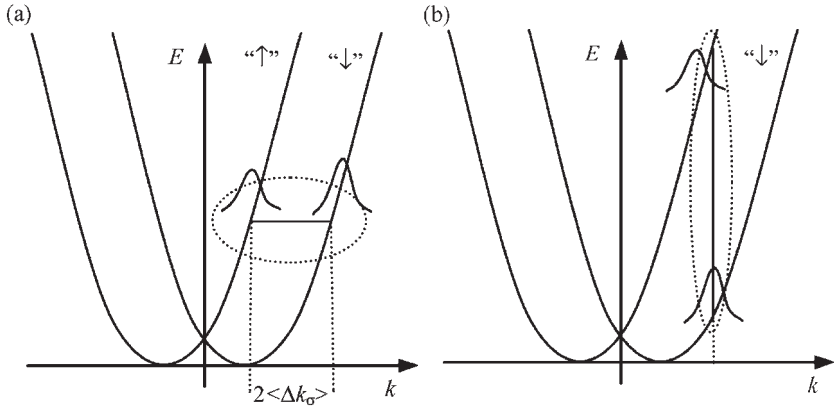

Figure 1. Different electron spin-polarized states in a 2DEG: (a) an electron state created with energy conservation; (b) an electron state created with wavevector conservation.

respectively. Similarly to equations (7) and (8), these matrix variables are projected onto the set of the basis matrices (σ_α, $\alpha = x, y, z$, and I) to get relations for the particle density, spin density, particle current density, and spin current density [35]:

$$\begin{aligned} n_n &= \int W_n \, \mathrm{d}^2k, \\ n_{\sigma_\alpha} &= \int W_{\sigma_\alpha} \, \mathrm{d}^2k, \end{aligned} \tag{12}$$

$$\begin{aligned} J_n^j &= \int (v_n^j W_n + v_{\sigma_\alpha}^j W_{\sigma_\alpha}) \, \mathrm{d}^2k, \\ J_{\sigma_\alpha}^j &= \int (v_n^j W_{\sigma_\alpha} + v_{\sigma_\alpha}^j W_n) \, \mathrm{d}^2k. \end{aligned} \tag{13}$$

The vector of the spin density can be expressed as $n_{\boldsymbol{\sigma}}(\mathbf{R}) = n_n(\mathbf{R})\mathbf{P}(\mathbf{R})$, where $\mathbf{P}(\mathbf{R})$ is, normalized to one, the spin polarization vector [48] of a small area, d^2r, of 2DEG at the position $\mathbf{R}$. The spin density, $n_{\boldsymbol{\sigma}}(\mathbf{R})$, corresponds to the density of the magnetic moment as $\boldsymbol{\mu}(\mathbf{R}) = -g\mu_{\mathrm{B}} n_{\boldsymbol{\sigma}}(\mathbf{R})$. We assume that the vectors of the particle and spin currents, equation (13), can be written using the average flow velocity, $\langle v_j \rangle$, as

$$\begin{aligned} J_n^j &= \langle v_j \rangle n_n + (v_{\sigma_\alpha}^j - \hbar \langle \Delta k_{\sigma_\alpha}^j \rangle / m^*) n_{\sigma_\alpha}, \\ J_{\sigma_\alpha}^j &= \langle v_j \rangle n_{\sigma_\alpha} + (v_{\sigma_\alpha}^j - \hbar \langle \Delta k_{\sigma_\alpha}^j \rangle / m^*) n_n. \end{aligned} \tag{14}$$

An additional parameter, $\langle \Delta k_{\sigma_\alpha}^j \rangle$, is introduced to define an electron spin-polarized state. For electrons injected from a ferromagnetic contact to the QW it is simple to show that $\langle \Delta k_{\sigma_\alpha}^j \rangle = m^* v_{\sigma_\alpha}^j / \hbar$ (see figure 1(a)), owing to the particle density and current density conservation at the interface [49]. Another possible electron spin-polarized state has been considered in the work by Mishschenko and Halperin [42], where $\langle \Delta k_{\sigma_\alpha}^j \rangle = 0$; see figure 1(b).

In this work we study the first case, where the electron spin-polarized state is created with constant energy rather than a constant electron wavevector. Moreover, we assume that the electron state is inhomogeneously broadened in momentum space due to temperature effects, and that the electron velocity can be expanded about the macroscopic flux velocity as

$$v_j = \langle v_j \rangle + \delta v_j. \tag{15}$$

The particle and spin conservation equations are obtained by integration of equation (5) over an electron wavevector:

$$\begin{aligned} &\frac{\partial n_n}{\partial t} + \frac{\partial J_n^j}{\partial x_j} = 0, \\ &\frac{\partial n_\sigma}{\partial t} + \frac{\partial J_\sigma^j}{\partial x_j} - \frac{2m^*}{\hbar}[v_\sigma^j \times J_\sigma^j] = 0. \end{aligned} \tag{16}$$

The effect of spin–orbit interaction appears in the second equation as a rotational term, where $[a_\sigma \times b_\sigma]$ is used for vector multiplication in the spin space. This term is proportional to the average flow velocity, unlike in the case of spin-polarized transport in an external magnetic field [37]. The drift-diffusion equations for the particle and spin currents are derived by applying the operator

$$\tfrac{1}{2}\int \{v^j, \ldots\}\, \mathrm{d}^2k \tag{17}$$

to equation (5). Assuming the conventional relation $\langle \delta v_n^j \delta v_n^l \rangle = \delta_{jl} kT/m^*$, we obtain the particle current density and spin current density:

$$\begin{aligned} J_n^j &= -\frac{\tau}{m^*}\left(kT\frac{\partial n_n}{\partial x_j} + \frac{\partial V}{\partial x_j} n_n\right), \\ J_\sigma^j &= -\frac{\tau}{m^*}\left(kT\frac{\partial n_\sigma}{\partial x_j} + \frac{\partial V}{\partial x_j} n_\sigma - \frac{2m^* kT}{\hbar}[v_\sigma^j \times n_\sigma]\right). \end{aligned} \tag{18}$$

In equation (18) we neglected terms quadratic in v_σ^j mixing polarized and non-polarized components of the current. These corrections should be considered in the model, accounting for non-conservation of the spin current density (equation (11)) in systems with a spin–orbit interaction [50]. The term proportional to $[v_\sigma^j \times n_\sigma]$ is responsible for the Dyakonov–Perel spin relaxation [20, 22]. Within the approximations applied, equations (16) and (18) do not mix electron transport in different spatial directions. Therefore, the spin-polarized transport in the 2DEG can be considered as a 1D problem.

The set of the drift-diffusion transport equations (16) and (18) for spin-polarized electrons in the presence of the spin–orbit interaction term linear in momentum is the main result of this work. To include the effects of an electron–electron interaction in the effective field approximation, the transport equations (16) and (18) should be supplemented with the Poisson equation.

3. Examples and discussion

We apply the derived equations to study transport in an asymmetric single QW grown in the (0, 0, 1) direction in terms of the crystallographic axes. The electric field in the plane of the QW is assumed homogeneous and equal to the external field, E. The x axis of the spatial coordinate system is oriented along the electric field and forms an angle ξ with the (1, 0, 0) direction in the xy plane. The spin–orbit interaction term, linear in the electron wavevector, is

$$\begin{aligned} H_{\mathrm{SO}} = {} & k_x(\sigma_y(\beta\langle k_z^2\rangle \sin 2\xi - \eta) - \sigma_x \beta \langle k_z^2\rangle \cos 2\xi) \\ & + k_y(\sigma_x(\beta\langle k_z^2\rangle \sin 2\xi + \eta) + \sigma_y \beta \langle k_z^2\rangle \cos 2\xi), \end{aligned} \tag{19}$$

where the spin coordinate system is oriented parallel to the spatial one. The parameters η and β are the Rashba [27] and Dresselhaus [28] spin–orbit coupling constants respectively. For the following derivation we specify a new spin coordinate system. The z spin axis is parallel to

the effective magnetic field produced by spin–orbit interaction (equation (19)) for an electron propagating along the external electric field direction. The y spin axis is orthogonal to the QW plane. In this coordinate system the spin–orbit term, H_{SO} (equation (1)), is

$$\begin{aligned} H_{SO} &= A_{xz}k_x\sigma_z + (A_{yx}\sigma_x + A_{yz}\sigma_z)k_y, \\ A_{xz} &= \sqrt{\eta^2 + (\beta\langle k_z^2\rangle)^2 - 2\eta\beta\langle k_z^2\rangle \sin 2\xi}, \\ A_{yx} &= ((\beta\langle k_z^2\rangle)^2 - \eta^2)/A_{xz}, \\ A_{yz} &= -2\eta\beta\langle k_z^2\rangle \cos 2\xi / A_{xz}. \end{aligned} \tag{20}$$

With the notation utilized, the spin polarization of electrons propagating parallel the x axis will precess about the z spin axis. In the case of Rashba spin–orbit interaction only, $\eta \neq 0$, $\beta = 0$, the z spin axis is oriented along the y spatial axis. For a non-zero Dresselhaus term, $\eta = 0$, $\beta \neq 0$, it is parallel to the x axis.

The drift-diffusion spin transport equation is

$$\frac{\partial n_\sigma}{\partial t} - \mathbf{D}\frac{\partial^2 n_\sigma}{\partial x^2} - \boldsymbol{\alpha}\frac{\partial n_\sigma}{\partial x} + \boldsymbol{\beta} n_\sigma = 0, \tag{21}$$

where

$$\begin{aligned} \mathbf{D} &= \begin{pmatrix} D & 0 & 0 \\ 0 & D & 0 \\ 0 & 0 & D \end{pmatrix}, \qquad \boldsymbol{\alpha} = \begin{pmatrix} \mu E & 2B_{xz}D & 0 \\ -2B_{xz}D & \mu E & 0 \\ 0 & 0 & \mu E \end{pmatrix}, \\ \boldsymbol{\beta} &= \begin{pmatrix} D(B_{xz}^2 + B_{yz}^2) & -\mu E B_{xz} & -B_{yx}B_{yz}D \\ \mu E B_{xz} & D(B_{xz}^2 + B_{yx}^2 + B_{yz}^2) & 0 \\ -B_{yx}B_{yz}D & 0 & DB_{yx}^2 \end{pmatrix}, \\ D &= \frac{kT\tau}{m^*}, \qquad \mu = -\frac{e\tau}{m^*}, \qquad B_{j\alpha} = \frac{2m^* A_{j\alpha}}{\hbar^2}. \end{aligned} \tag{22}$$

We consider a few examples of spin dynamics in a 2DEG using equation (21):

(1) At time $t = 0$ the 2DEG is homogeneously polarized, $E = 0$, and only one spin–orbit interaction mechanism (Rashba or Dresselhaus) is responsible for the spin evolution. Equation (21) is transformed to the spin relaxation equation

$$\frac{\partial n_\sigma}{\partial t} = -\boldsymbol{\beta} n_\sigma, \tag{23}$$

where the coupling coefficients $\boldsymbol{\beta}$ are equal to that derived in the work [22].

(2) Stationary injection of spin-polarized electrons at the position $x = 0$ into an infinite QW. The spin–orbit constants are coupled by the relation $\eta = \beta\langle k_z^2\rangle$. Equation (21) can be diagonalized for any orientation of electron transport with respect to the crystallographic directions. The solution is

$$\begin{aligned} n_{\sigma_x} &= n_{\sigma_x}^0 \exp\left\{-\left(\frac{\mu E}{2D} + \sqrt{\left(\frac{\mu E}{2D}\right)^2 + B_{yz}^2}\right)x\right\}\cos(B_{xz}x), \\ n_{\sigma_y} &= n_{\sigma_y}^0 \exp\left\{-\left(\frac{\mu E}{2D} + \sqrt{\left(\frac{\mu E}{2D}\right)^2 + B_{yz}^2}\right)x\right\}\sin(B_{xz}x), \\ n_{\sigma_z} &= n_{\sigma_z}^0 \exp\left\{-\left(\frac{\mu E}{2D} + \left|\frac{\mu E}{2D}\right|\right)x\right\}. \end{aligned} \tag{24}$$

Analogously to [16], for the z component of the spin polarization the Dyakonov–Perel spin relaxation mechanism [20, 22] is suppressed. The transverse component of the spin polarization evolves about the effective magnetic field and decays with the characteristic spin dephasing length $L_\perp = \left(\frac{\mu E}{2D} + \sqrt{\left(\frac{\mu E}{2D}\right)^2 + B_{yz}^2}\right)^{-1}$. The effect of the electric field, E, on the transverse spin dephasing length is similar to that obtained in [51] and later considered in [8, 37]. The remarkable property of the solution (24) is that the temperature affects coefficients $B_{j\alpha}$ through an effective mass only. Usually, this effect is weak. Therefore, the distributions of the spin polarization should be nearly the same at different temperatures once the relation $\mu E/D$ is conserved.

We estimate the length of the coherent spin precession, $L_\mathrm{p} = 2\pi/B_{xz}$, and the transverse spin dephasing length, $L_\perp$, for a 10 nm width GaAs/AlGaAs QW. The Dresselhaus spin–orbit constant $\beta = 25.5$ eV Å^3 is taken from [52]. This corresponds to $\beta\langle k_z^2\rangle = 0.025$ eV Å. The same order of the value for the Rashba spin–orbit coupling constant can be achieved by an appropriate doping of a heterostructure. The calculated values of the spin dephasing and spin precession lengths for different orientations of the electron transport are shown in figure 2.

The essential requirement for the realization of the spin-FET proposed by Datta and Das [10] is $L_\mathrm{p} \ll L_\perp$. As follows from figure 2, this relation is valid for transport within the small angle about the $(1, -1, 0)$ direction in terms of the crystallographic axes. The applied electric field increases the spin dephasing length while it does not affect the spin precession length. Moreover, the range of transport directions usable for the spin-FET [10] varies with the in-plane electric field.

(3) Stationary injection of spin-polarized electrons into a 10 nm width GaAs/AlGaAs QW along the $(1, -1, 0)$ crystallographic direction. At the injection boundary $n_{\sigma_x} = 0$, $n_{\sigma_y} = 0$, $n_{\sigma_z} = n_{\sigma_z}^0$. The spin–orbit coupling constants are not equal. This configuration can be utilized for the spin-FET proposed by Schliemann *et al* [16]. The longitudinal spin density component, n_{σ_z}, decays as

$$n_{\sigma_z} = n_{\sigma_z}^0 \exp\left\{-\left(\frac{\mu E}{2D} + \sqrt{\left(\frac{\mu E}{2D}\right)^2 + B_{yx}^2}\right)x\right\} \tag{25}$$

where $B_{yx} = 2m^*\Delta\eta/\hbar^2$. The calculated spin dephasing length, $L_\parallel$, as a function of the relative difference of the spin–orbit constants, $\Delta = |\eta/(\beta\langle k_z^2\rangle) - 1|$, is shown in figure 3. Design of a non-ballistic spin-FET [16] requires efficient modulation of the spin scattering length from $L_\parallel \gg L_\mathrm{D}$ to $L_\parallel \ll L_\mathrm{D}$ (L_D is the device length), varying the difference between the Rashba and Dresselhaus spin–orbit constants. As follows from our calculations (see figure 3), the device operation can be optimized for different relations between η and $\beta\langle k_z^2\rangle$ by varying the in-plane electric field. For example, if the allowable range of spin dephasing lengths is limited to $L_\mathrm{D}/3 < L_\parallel < 3L_\mathrm{D}$, then for the device length $L_\mathrm{D} = 3$ μm, Δ should be varied in the ranges $0.03 < \Delta < 0.22$ ($E = -0.01$ V cm^{-1}), $0.08 < \Delta < 0.3$ ($E = -100$ V cm^{-1}), $0.13 < \Delta < 0.42$ ($E = -300$ V cm^{-1}), and $0.22 < \Delta < 0.69$ ($E = -1$ kV cm^{-1}). Recent experiments demonstrate that 25% variations of the parameter Δ can be easily achieved for GaAs/AlGaAs QWs [23].

Although within the model developed for $\Delta = 0$ the Dyakonov–Perel spin relaxation is completely suppressed for spatial transport along the $(1, -1, 0)$ direction, other spin scattering mechanisms will determine the spin dephasing. For example, it was shown in [53] that spin–orbit terms cubic in an electron momentum can appreciably modify the spin dynamics

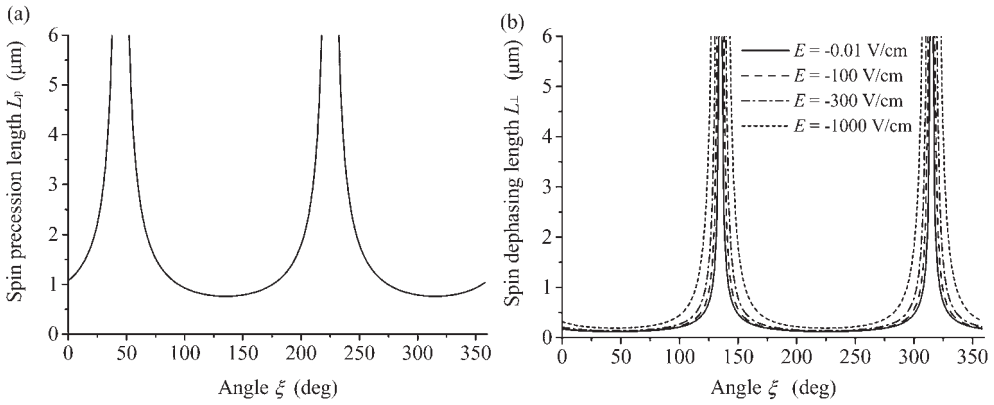

Figure 2. The spin precession length (a) and transverse spin dephasing length (b) for different transport orientations with respect to the (0, 0, 1) direction in terms of the crystallographic axes at room temperature.

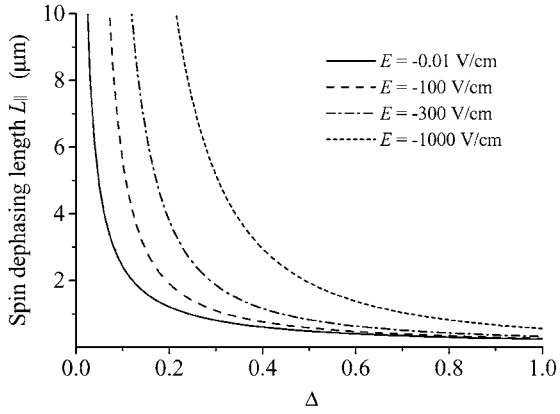

Figure 3. The longitudinal spin dephasing length as a function of the relative difference of the spin–orbit constants $\Delta = |\eta/(\beta\langle k_z^2\rangle) - 1|$ for different values of the applied electric field at $T = 300$ K. The electric field is along the $(1, -1, 0)$ direction.

during the ballistic transport in a quasi-1D structure. The upper limit for the spin dephasing length will possibly be defined by the spin scattering at nuclear spins [54]. For high electron concentrations electron–electron collisions will affect the spin dynamics even without the momentum dissipation [35, 43, 55]. Moreover, the relaxation time approximation is a rough model for transport in polar semiconductors [56]. However, we assume that equations (21) and (22) can be valid for a more general, field-dependent form of the diffusion coefficient and mobility $D(E)$, $\mu(E)$ [21]. These transport parameters can be obtained from a Monte Carlo modelling [57, 58].

All-electric measurements of the spin polarization in semiconductor heterostructures [59–61] are complicated by additional spin-independent effects [60]. Existing results do not have a single theoretical explanation [32, 33, 60, 62]. However, our model is consistent with results obtained in experiments on optical manipulation of spin coherence in strained semiconductor layers [63]. For spin transport along the $(1, -1, 0)$ direction, the spin polarization rotates by

the angle $\varphi = \pi$ in a distance $l \approx 23\ \mu$m. This value is independent of the electric field in the range studied. Moreover, in stronger electric fields the spin coherence is conserved for a longer distance. Both of these features agree with our results (equation (24)).

A drift-diffusion transport equation similar to equation (21) has been obtained by Pershin [64] within a stochastic approach.

4. Conclusions

We have developed a semiclassical drift-diffusion model for spin-polarized transport in a non-degenerate 2DEG controlled by the spin–orbit interaction linear in the electron momentum. The effects of an in-plane electric field and the transport orientation with respect to the crystallographic directions are discussed for a single quantum well grown in the (0, 0, 1) direction in terms of the crystallographic axes. The model derived agrees with results of optical measurements of the coherent spin dynamics in semiconductor layers. It could be useful for spintronic device modelling.

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