Controllable kinetic magnetoelectric effect in two-dimensional electron gases with both Rashba and Dresselhaus spin-orbit couplings

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A microscopic model calculation of the electric-field-induced nonequilibrium spin accumulation due to kinetic magnetoelectric effect in two-dimensional electron gases with both Rashba and Dresselhaus spin-orbit couplings is presented. The results show that by modulating the relative strength of the two different spin-orbit couplings, both the magnitude and the polarization direction of the nonequilibrium spin accumulation can be effectively controlled.

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During the last several years, there has been growing interest in spin-based semiconductor electronics, where the electron spin rather than charge is at the very center of interest.[1] Creation and effective control of spin-polarized currents and/or nonequilibrium spin densities in nonmagnetic semiconductors are necessary for the practical implementation of semiconducting spin electronics, and various ways have been proposed for solving this problem; e.g., spin injections from ferromagnetic metals or from diluted magnetic semiconductors.[2–6] In order to take advantage of traditional semiconductor technologies, it would be much more attractive if application of an electric field alone would suffice to induce a nonequilibrium spin density or spin-polarized current in a nonmagnetic semiconductor, and several novel effects, e.g., *intrinsic* or *extrinsic* spin Hall effect[7–12] and kinetic magnetoelectric effect,[13,14] have been predicted that might offer such opportunities. All these effects rely on spin-orbit interactions in semiconducting materials. A spin Hall effect can arise from both intrinsic and extrinsic spin-orbit couplings in two-dimensional or bulk semiconducting materials, in which an external electric field will induce a spin current in the direction perpendicular to the electric field, and a nonequilibrium and inhomogeneous spin accumulation can hence result near the edges of a sample.[11,12] Kinetic magnetoelectric effect is due to intrinsic spin-orbit couplings in two-dimensional electronic systems, e.g., Rashba spin-orbit coupling in two-dimensional electron gases. In the kinetic magnetoelectric effect in a Rashba two-dimensional electron gas, an external electric field will induce a homogeneous nonequilibrium spin density in the sample with the spin polarization direction perpendicular to the electric field.[13,14] In both these effects, the magnitude of the nonequilibrium spin density is proportional to the driving electric field and hence can be easily controlled by varying the electric field. For the practical uses of these effects, it would be desirable if both the magnitude and the polarization direction of the nonequilibrium spin accumulation can be effectively controlled. In this brief report we investigate how to control the polarization direction of the nonequilibrium spin accumulation due to kinetic magnetoelectric effect in a semiconducting material. We show that if there are two different intrinsic spin-orbit couplings existing in a semiconducting material and the relative strength of the two different spin-orbit interactions can be effectively modulated by applying some kinds of external forces (e.g., by a gate voltage), then both the magnitude and the polarization direction of the nonequilibrium spin accumulation due to kinetic magnetoelectric effect in such a system might be effectively controlled. Two-dimensional electron gases (2DEGs) with both Rashba and Dresselhaus spin-orbit couplings can serve as examples of such a system. As is well known, Rashba spin-orbit coupling in 2DEGs is caused by the lacking of inversion symmetry of the trapping wells and can be effectively modulated by applying a gate voltage. In contrast, Dresselhaus spin-orbit coupling is due to structural inversion asymmetry and hence is independent of the applied gate voltage. Thus, in 2DEGs with both Rashba and Dresselhaus spin-orbit couplings, the relative strength of the two different spin-orbit interactions can be effectively modulated by applying a gate voltage. In this brief report we present a microscopic model calculation of the electric-field-induced nonequilibrium spin accumulation due to kinetic magnetoelectric effect in 2DEGs with both Rashba and Dresselhaus spin-orbit couplings. (Nonequilibrium spin accumulation due to an intrinsic spin Hall effect in a thin strip of such a system was investigated by Malshukov *et al.* For details please refer to Ref. 15.) We will investigate in detail how the magnitude and the polarization direction of the nonequilibrium spin accumulation changes as the relative strength of the two different spin-orbit interactions varies.

The single-particle Hamiltonian for 2DEGs with both Rashba and Dresselhaus spin-orbit coupling reads[16,17]

$$\hat{H}_{so} = \frac{\hbar^2 k^2}{2m} + \alpha(\hat{\sigma}_x k_y - \hat{\sigma}_y k_x) + \beta(\hat{\sigma}_x k_x - \hat{\sigma}_y k_y), \qquad (1)$$

where $\hat{\sigma}_i$ $(i=x,y,z)$ are the usual Pauli matrices, $\mathbf{k}$ is the wave vector of a conduction electron, α is the Rashba spin-orbit coupling constant, and β the Dresselhaus spin-orbit coupling constant. For a given wave vector $\mathbf{k}$, the eigenfunctions of $\hat{H}_{so}$ can be expressed as

$$|\mathbf{k}\lambda\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} e^{i\phi_{\mathbf{k}}/2} \\ \lambda e^{-i\phi_{\mathbf{k}}/2} \end{pmatrix}, \qquad (2)$$

where $\lambda = \pm 1$ and the angle $\phi_{\mathbf{k}}$ is defined by

$$\tan\phi_{\mathbf{k}} = \frac{\alpha k_x + \beta k_y}{\alpha k_y + \beta k_x}. \tag{3}$$

The corresponding eigenvalues of $\hat{H}_{so}$ are given by

$$\epsilon_{\mathbf{k}\lambda} = \frac{\hbar^2\mathbf{k}^2}{2m} + \lambda\sqrt{(\alpha^2+\beta^2)k^2 + 4\alpha\beta k_x k_y}. \tag{4}$$

The expectation value of the spin of an electron in the state $|\mathbf{k}\lambda\rangle$ will given by

$$\mathbf{S}_\lambda^{(0)}(\mathbf{k}) = \frac{\hbar}{2}\langle\mathbf{k}\lambda|\hat{\boldsymbol{\sigma}}|\mathbf{k}\lambda\rangle = S_{\lambda,x}^{(0)}(\mathbf{k})\mathbf{e}_x + S_{\lambda,y}^{(0)}(\mathbf{k})\mathbf{e}_y, \tag{5}$$

where

$$S_{\lambda,x}^{(0)}(\mathbf{k}) = \frac{\hbar}{2}\lambda\cos(\phi_{\mathbf{k}}), \quad S_{\lambda,y}^{(0)}(\mathbf{k}) = -\frac{\hbar}{2}\lambda\sin(\phi_{\mathbf{k}}). \tag{6}$$

In the presence of impurity scatterings, the total Hamiltonian will be $\hat{H}=\hat{H}_{so}+V(\mathbf{r})$, where $V(\mathbf{r})$ is the impurity scattering potential. (For simplicity, we assume that only *spinless* impurity scatterings are present.) The scattering state of a conduction electron can be determined from the Lippman-Schwinger equation.[18] In the Born approximation, one has

$$|\mathbf{k}\lambda\rangle_{sc} = |\mathbf{k}\lambda\rangle + \sum_{k'\lambda'}\frac{V_{\mathbf{k}'\lambda',\mathbf{k}\lambda}}{\epsilon_{k\lambda} - \epsilon_{k'\lambda'} + i0^+}|\mathbf{k}'\lambda'\rangle, \tag{7}$$

where $|\mathbf{k}\lambda\rangle_{sc}$ is the scattering state and $V_{\mathbf{k}'\lambda',\mathbf{k}\lambda} = \langle\mathbf{k}'\lambda'|V(\mathbf{r})|\mathbf{k}\lambda\rangle$ the scattering matrix elements. The expectation value of the electron spin will be modified to be $S_{\lambda,i}(\mathbf{k})=\frac{\hbar}{2}\langle\mathbf{k}\lambda|\hat{\sigma}_i|\mathbf{k}\lambda\rangle_{sc}$ $(i=x,y,z)$. Substituting Eq. (7) into this equation, one gets

$$S_{\lambda,i}(\mathbf{k}) = S_{\lambda,i}^{(0)}(\mathbf{k}_0) + \hbar\sum_{k'\lambda'}\mathrm{Re}\left[\frac{V_{\mathbf{k}'\lambda',\mathbf{k}\lambda}\langle\hat{\sigma}_i\rangle_{\mathbf{k}\lambda,\mathbf{k}'\lambda'}}{\epsilon_{\mathbf{k}\lambda} - \epsilon_{\mathbf{k}'\lambda'} + i0^+}\right], \tag{8}$$

where $\langle\hat{\sigma}_i\rangle_{\mathbf{k}\lambda,\mathbf{k}'\lambda'} \equiv \langle\mathbf{k}\lambda|\hat{\sigma}_i|\mathbf{k}'\lambda'\rangle$. In the weak impurity scattering regime, the off-diagonal matrix elements of the density matrix can be neglected and the net spin density can be given by

$$\langle S_i\rangle = \sum_{\mathbf{k},\lambda} S_{\lambda,i}(\mathbf{k})f_\lambda(\mathbf{k}), \tag{9}$$

where $\langle S_i\rangle$ is the ith component of the net spin density and $f_\lambda(\mathbf{k})$ the distribution function of conduction electrons (i.e., the diagonal matrix elements of the density matrix). In the absence of external electric field, $f_\lambda(\mathbf{k})$ will be given by the equilibrium Fermi distribution function, $f_\lambda(\mathbf{k})=f_0(\epsilon_{\mathbf{k}\lambda}) =[1+e^{(\epsilon_{\mathbf{k}\lambda}-\epsilon_F)/k_BT}]^{-1}$. From Eqs. (2) and (8), one can show that $S_{\lambda,i}(-\mathbf{k})=-S_{\lambda,i}(\mathbf{k})$. Since $f_0(\epsilon_{\mathbf{k}\lambda})= f_0(\epsilon_{-\mathbf{k}\lambda})$, the integrals in Eq. (9) vanish exactly in the equilibrium state, i.e., no net spin density survives in the absence of external electric fields. When an external electric field is applied, the distribution function $f_\lambda(\mathbf{k})$ will become asymmetric under the transformation $\mathbf{k}\leftrightarrow-\mathbf{k}$ since the Fermi sphere will be displaced along the direction of the external electric field; hence a net spin density may be result. In the weak impurity scattering regime, the nonequilibrium distribution function can be derived by solving the Boltzmann transport equation, which reads (in a steady state and assuming a homogeneous system)

$$\dot{\mathbf{k}}\cdot\frac{\partial f_\lambda(\mathbf{k})}{\partial\mathbf{k}} = \left(\frac{\partial f_\lambda}{\partial t}\right)_{\text{coll.}}, \tag{10}$$

where $\dot{\mathbf{k}}=-e\mathbf{E}/\hbar$ and $(\partial f_\lambda/\partial t)_{\text{coll.}}$ is the collision integral due to impurity scatterings:

$$\left(\frac{\partial f_\lambda}{\partial t}\right)_{\text{coll.}} = -\sum_{k'\lambda'} w_{\mathbf{k}\lambda,\mathbf{k}'\lambda'}f_\lambda(\mathbf{k})[1-f_{\lambda'}(\mathbf{k}')]\delta(\epsilon_{\mathbf{k}\lambda}-\epsilon_{\mathbf{k}'\lambda'}) + \sum_{k'\lambda'} w_{\mathbf{k}'\lambda',\mathbf{k}\lambda}f_{\lambda'}(\mathbf{k}')[1-f_\lambda(\mathbf{k})]\delta(\epsilon_{\mathbf{k}\lambda}-\epsilon_{\mathbf{k}'\lambda'}). \tag{11}$$

where $w_{\mathbf{k}\lambda,\mathbf{k}'\lambda'}$ are the transition probabilities given by the Fermi's golden rule, $w_{\mathbf{k}\lambda,\mathbf{k}'\lambda'} \equiv (2\pi/\hbar)n_i|V_{\mathbf{k}\lambda,\mathbf{k}'\lambda'}|^2$ (n_i is the density of impurities). In the linear response regime, the general form of the solution for $f_\lambda(\mathbf{k})$ can be expressed as

$$\delta f_\lambda(\mathbf{k}) = e\frac{\partial f_0(\epsilon_{\mathbf{k}\lambda})}{\partial\epsilon_{\mathbf{k}\lambda}}[a_{\mathbf{k}\lambda}(\mathbf{E}\cdot\mathbf{v}_{\mathbf{k}\lambda}) + b_{\mathbf{k}\lambda}(\mathbf{E}\times\mathbf{e}_z)\cdot\mathbf{v}_{\mathbf{k}\lambda}], \tag{12}$$

where $\delta f_\lambda(\mathbf{k})\equiv f_\lambda(\mathbf{k})-f_0(\epsilon_{\mathbf{k}\lambda})$, $\mathbf{v}_{\mathbf{k}\lambda}=\hbar^{-1}\nabla_{\mathbf{k}}\epsilon_{\mathbf{k}\lambda}$ is the velocity of conduction electrons, $\mathbf{e}_z$ is a unit vector along the normal of the two-dimensional plane, $a_{\mathbf{k}\lambda}$ and $b_{\mathbf{k}\lambda}$ are two unknown coefficients that need to be determined self-consistently from Eq. (10). The last term in Eq. (12) characterizes the anisotropy of impurity scatterings. Because the energy spectrums of conduction electrons are anisotropic in the presence of both Rashba and Dresselhaus spin-orbit couplings, the impurity scatterings may be highly anisotropic even if the impurity scattering potential $V(\mathbf{r})$ is isotropic. Substituting Eq. (12) into Eq. (10), one can find that the coefficients $a_{\mathbf{k}\lambda}$ and $b_{\mathbf{k}\lambda}$ will satisfy the following equations:

$$a_{\mathbf{k}\lambda}/\tau_{\mathbf{k}\lambda}^{(1)} + b_{\mathbf{k}\lambda}/\tau_{\mathbf{k}\lambda}^{(2)} = 1, \tag{13}$$

$$a_{\mathbf{k}\lambda}/\tau_{\mathbf{k}\lambda}^{(2)} - b_{\mathbf{k}\lambda}/\tau_{\mathbf{k}\lambda}^{(1)} = 0, \tag{14}$$

where $\tau_{\mathbf{k}\lambda}^{(1,2)}$ are two characteristic relaxation times defined by

$$\frac{1}{\tau_{\mathbf{k}\lambda}^{(1)}} = \sum_{\mathbf{k}',\lambda'} w_{\mathbf{k}'\lambda',\mathbf{k}\lambda}\left\{1 - \frac{|\mathbf{v}_{\mathbf{k}'\lambda'}|}{|\mathbf{v}_{\mathbf{k}\lambda}|}\cos[\theta(\mathbf{v}_{\mathbf{k}\lambda}\wedge\mathbf{v}_{\mathbf{k}'\lambda'})]\right\}, \tag{15}$$

$$\frac{1}{\tau_{\mathbf{k}\lambda}^{(2)}} = \sum_{\mathbf{k}',\lambda'} w_{\mathbf{k}'\lambda',\mathbf{k}\lambda}\frac{|\mathbf{v}_{\mathbf{k}'\lambda'}|}{|\mathbf{v}_{\mathbf{k}\lambda}|}\sin[\theta(\mathbf{v}_{\mathbf{k}\lambda}\wedge\mathbf{v}_{\mathbf{k}'\lambda'})], \tag{16}$$

in which $\theta(\mathbf{v}_{\mathbf{k}\lambda}\wedge\mathbf{v}_{\mathbf{k}'\lambda'})$ denotes the angle between $\mathbf{v}_{\mathbf{k}\lambda}$ and $\mathbf{v}_{\mathbf{k}'\lambda'}$. From Eqs. (13) and (14), one gets

$$a_{\mathbf{k}\lambda} = \frac{\tau_{\mathbf{k}\lambda}^{(1)}}{1+[\tau_{\mathbf{k}\lambda}^{(1)}/\tau_{\mathbf{k}\lambda}^{(2)}]^2}, \quad b_{\mathbf{k}\lambda} = \frac{\tau_{\mathbf{k}\lambda}^{(2)}}{1+[\tau_{\mathbf{k}\lambda}^{(2)}/\tau_{\mathbf{k}\lambda}^{(1)}]^2}. \tag{17}$$

The electric-field-induced nonequilibrium spin density can be calculated by substituting Eq. (8) and the solution for

$f_\lambda(\mathbf{k})$ into Eq. (9), but no simple expressions can be obtained due to the complexities involved in the integrals in Eqs. (8) and (15) to (16); hence, numerical calculations must be resorted to.[19] For simplicity, in the following we will assume a δ-function-shaped (short-ranged) impurity scattering potential, $V(\mathbf{r})=V_0\delta(\mathbf{r})$, where V_0 characterizes the impurity scattering strength. The following parameters will be needed in the numerical calculations: the electron's effective mass m, the electron density n, the impurity density n_i, the impurity scattering strength V_0, and the spin-orbit coupling constants α and β. We take $m=0.06m_e$, $n=1.9\times10^{12}$ cm^{-2}, $n_i=1.0\times10^{10}$ cm^{-2}, and $V_0=1.4\times10^{-11}$ meV cm. These values are typical of current high quality 2DEG samples. The Rashba and Dresselhaus spin-orbit coupling strengths will be characterized by two spin-orbit coupling energy scales $\varepsilon_\alpha=m\alpha^2/\hbar^2$ and $\varepsilon_\beta\equiv m\beta^2/\hbar^2$. In current high quality 2DEG samples, the range of ε_α and ε_β is $0\sim3$ meV.[20] We assume that the external electric field $\mathbf{E}$ is applied along the x direction, $\mathbf{E}=E_x\mathbf{e}_x$. From Eqs. (9) and (12), one can see that both $\langle S_x\rangle$ and $\langle S_y\rangle$ will be proportional to E_x, hence the magnitude of the nonequilibrium spin accumulation (given by $\langle S\rangle=[\langle S_x\rangle^2+\langle S_y\rangle^2]^{1/2}$) will also be proportional to E_x but the spin polarization direction will be independent of E_x. In order to see whether the polarization direction of the nonequilibrium spin accumulation and/or its magnitude can be effectively controlled by tuning the relative strength of the Rashba and Dresselhaus spin-orbit couplings, in Figs. 1(a) and 1(b) and Figs. 2(a) and 2(b), we have plotted schematically the variations of both the magnitude and the polarization direction of the nonequilibrium spin accumulation with the changes of the Rashba and Dresselhaus spin-orbit coupling energies, respectively. (It should be noted that, unlike the intrinsic spin Hall effect in thin strips of such systems,[15] the kinetic magnetoelectric effect does not lead to the generation of a nonvanishing z component of the spin density.[13,14]) From Figs. 1 and 2 one can see that the variations of $\langle S\rangle$ with the changes of the spin-orbit coupling energies are not very substantial ($\langle S\rangle/E_x$ remains on the order of 10^5 μ_B cm^{-1}/V), but the polarization direction of the nonequilibrium spin accumulation can be modulated significantly as the relative strength of the two different spin-orbit interactions varies. From Figs. 1(b) and 2(b), one can see that, if $\alpha\gg\beta$, the polarization direction of the spin accumulation will tend to be perpendicular to the external electric field, and if $\alpha\ll\beta$, the spin polarization direction will tend to be parallel to the external electric field. This suggests that by tuning the relative strength of the two different spin-orbit interactions (which can be achieved by applying a gate voltage), the polarization direction of the nonequilibrium spin accumulation can be effectively controlled. From the physical point of view, this controllability is due to the following. As can be seen from the Hamiltonian (1), both Rashba and Dresselhaus spin-orbit couplings act like a momentum-dependent effective magnetic field, but the directions of the effective magnetic fields due to Rashba and Dresselhaus spin-orbit couplings are different from each other. In the absence of external electric fields, the averaged effective magnetic fields due to both Rashba and Dresselhaus spin-orbit couplings vanish exactly, and hence no net spin density can survive. If

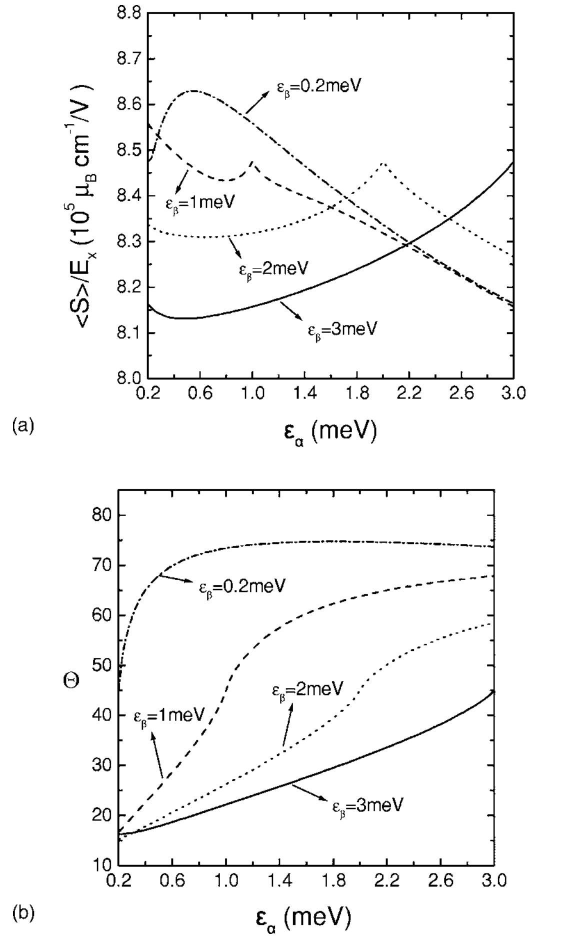

FIG. 1. Variations of (a) the magnitude (divided by the electric field E_x) and (b) the polarization direction of the nonequilibrium spin accumulation with the changes of the Rashba spin-orbit coupling energy ε_α. In both (b) and Fig. 2(b), the symbol Θ stands for the angle (in units of degrees) between the spin polarization direction and the external electric field. The parameters used were given in the text and shown in the figures.

an external electric field is applied (e.g., in the x direction), the averaged effective magnetic fields due to both Rashba and Dresselhaus spin-orbit couplings will become nonvanishing because $\langle k_x\rangle\neq0$ and also $\langle k_y\rangle\neq0$ if $\beta\neq0$; thus a nonequilibrium spin density can result. Because $\langle k_x\rangle\gg\langle k_y\rangle$, the average effective magnetic field $\langle\mathbf{H}_R\rangle$ due to Rashba spin-orbit coupling will tend to be along the y direction [i.e., $\langle\mathbf{H}_R\rangle=\langle\alpha(k_y\mathbf{e}_x-k_x\mathbf{e}_y)\rangle\simeq-\alpha\langle k_x\rangle\mathbf{e}_y$] but the effective magnetic field $\langle\mathbf{H}_D\rangle$ due to Dresselhaus spin-orbit coupling tends to be along the x direction [i.e., $\langle\mathbf{H}_D\rangle=\langle\beta(k_x\mathbf{e}_x-k_y\mathbf{e}_y)\rangle\simeq\beta\langle k_x\rangle\mathbf{e}_x$]. The total effective magnetic field felt by the conduction electrons is the sum of the two effective fields $\langle\mathbf{H}_R\rangle$ and $\langle\mathbf{H}_D\rangle$, whose direction will depend sensitively on the relative strength of the Rashba and Dresselhaus spin-orbit couplings, thus, by tuning the relative strength of the two different spin-orbit couplings, the polarization direction of the nonequilibrium spin accumulation can be effectively modulated.

An interesting question related to kinetic magnetoelectric effect is that whether the electric-field-driven charge current will become spin polarized due to the occurrence of nonequi-

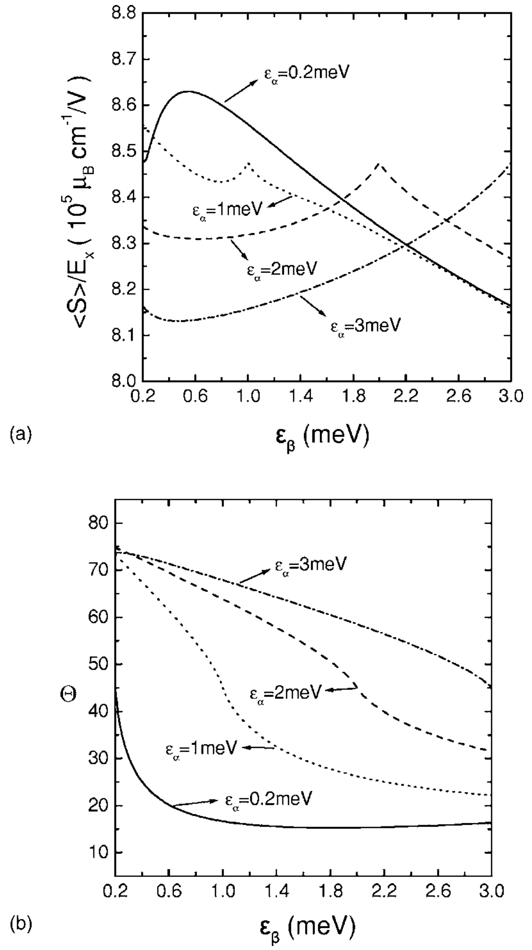

FIG. 2. Variations of (a) the magnitude (divided by the electric field E_x) and (b) the polarization direction of the nonequilibrium spin accumulation with the changes of the Dresselhaus spin-orbit coupling energy ε_β. The parameters used were given in the text and shown in the figures.

librium in-plane spin accumulation. Recently, Rashba warned that, in a spin-orbit coupled system, nonvanishing background spin currents (in the standard definitions of spin currents) may be obtained even in the absence of external electric fields, and such *equilibrium* background spin currents are not *transport* spin currents and should be removed in the calculations of transport spin currents.[21] After removing the equilibrium background spin currents, the transport spin current flowing to the x direction with spin parallel to the x or y axis will be given by

$$J_x^i = \sum_{\mathbf{k},\lambda} J_x^i(\mathbf{k},\lambda)\,\delta f_\lambda(\mathbf{k}), \quad (i=x,y), \tag{18}$$

where $\delta f_\lambda(\mathbf{k}) = f_\lambda(\mathbf{k}) - f_0(\epsilon_{\mathbf{k}\lambda})$ and $J_x^i(\mathbf{k},\lambda) \equiv \langle \mathbf{k}\lambda | \hat{J}_x^i | \mathbf{k}\lambda \rangle_{sc}$, in which $\hat{J}_x^i$ is the spin current operator defined by $\hat{J}_x^i = \frac{\hbar}{2}\{\sigma_i, \hat{v}_x\}$ and $\hat{v}_x = \hbar^{-1}\,\partial \hat{H}/\partial k_x$ is the velocity operator. The transport spin current flowing to the y direction with spin parallel to the x or y direction can be calculated similarly. By use of Eqs. (2) and (7), one can show that $J_x^i(-\mathbf{k},\lambda) = j_x^i(\mathbf{k},\lambda)$. On the other hand one, has $\delta f_\lambda(-\mathbf{k}) = -\delta f_\lambda(\mathbf{k})$ according to Eq. (12); thus, the integrals in Eq. (18) vanish exactly, suggesting that no spin-polarized transport currents accompany with the electric-field-driven nonequilibrium in-plane spin accumulation. A similar conclusion was obtained in Ref. 13 for a Rashba two-dimensional electron gas based on the Green's function approach.

In summary, a microscopic model calculation of kinetic magnetoelectric effect in two-dimensional electron gases with both Rashba and Dresselhaus spin-orbit couplings was presented. The results show that by tuning the relative strength of the two different spin-orbit couplings, both the magnitude and the polarization direction of the nonequilibrium spin accumulation due to this effect can be effectively controlled. Finally, it should be stressed that the approach used in the present paper is valid only in the weak impurity scattering regime, where the spin-orbit eigenstates are well defined objects. In the presence of strong impurity scatterings, the off-diagonal matrix elements of the density matrix become important and hence more strict theoretical approaches (e.g., the quantum kinetic equation approach) should be applied.

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