Vol 18 No 8, August 2009
1674-1056/2009/18(08)/3523-07

Chinese Physics B

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The kinetic magnetoelectric effect in laterally boundary-confined ballistic two-dimensional hole gases*

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(Received 26 August 2008; revised manuscript received 23 March 2009)

A theoretical investigation is presented on the characteristics of the kinetic magnetoelectric effect in laterally boundary-confined ballistic two-dimensional hole gases. It was shown that, though the momentum-dependent effective magnetic fields felt by charge carriers due to the spin-orbit interaction are in-plane orientated in such systems, both in-plane polarized and normal polarized nonequilibrium spin polarization densities could be electrically induced by the kinetic magnetoelectric effect, and the induced nonequilibrium spin polarizations exhibit some interesting characteristics. The characteristics we found indicate that there may be some possible relation between this effect and some recent experimental findings.

Keywords: spin-orbit coupling, electrically induced spin polarization, two-dimensional hole gas
PACC: 7225, 7335A, 7215G

1. Introduction

Electrically-induced nonequilibrium spin polarizations and/or spin currents in spin-orbit coupled nonmagnetic semiconducting electronic systems have recently attracted considerable theoretical and experimental interest because of their possible applications in semiconductor spintronics.[1] Many of these recent studies have been focused on the so-called spin Hall effect,[2−12] which consists of the generation of a spin current in the direction perpendicular to a charge current circulating through a spin-orbit coupled system. In the spin Hall effect in a laterally boundary-confined spin-orbit coupled system, the flow of the transverse spin Hall current will cause nonequilibrium accumulations of spins in the vicinity of the boundaries, leading to the generation of electrically-induced antisymmetric boundary spin polarizations. From an experimental point of view, since the spin Hall current is not a quantity that can be measured directly, electrically-induced antisymmetric boundary spin polarizations in laterally boundary-confined spin-orbit coupled systems should be the principal observable signature of the spin Hall effect. Recently, electrically-induced antisymmetric boundary spin polarizations have indeed been detected experimentally in some laterally boundary-confined spin-orbit coupled systems and interpreted theoretically as the direct manifestation of the spin Hall effect.[7−12] But it should be noted that, though intensive theoretical works have been devoted to the understanding of the spin Hall effect, a number of controversial issues still exist concerning some fundamental aspects of this effect or the theoretical interpretations of some recent relevant experimental data.[13−31] For example, as was argued in Refs.[25]–[28], even how to define the spin Hall current in a spin-orbit coupled system is also a much more subtle issue. Due to the existence of such controversial issues, in addition to the spin Hall effect, other possible effects and/or contributions should also be examined in order to obtain a clearer understanding of some recent experimental findings.

A different effect that can also lead to the generation of electrically-induced nonequilibrium spin polarizations in spin-orbit coupled semiconducting electronic systems is the so-called kinetic magnetoelectric effect.[32−39] Like the spin Hall effect, this effect also arises from the spin-orbit interaction of charge carriers. But unlike the spin Hall effect, in this ef-

*Project supported by the National Natural Science Foundation of China (Grant No 10874049), the State Key Program for Basic Research of China (Grant No 2007CB925204) and the Natural Science Foundation of Guangdong Province of China (Grant No 07005834).

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fect the generation of electrically-induced nonequilibrium spin polarization densities is not due to the flow of a spin current but originates from the combined action of the spin-orbit coupling and the time-reversal symmetry-breaking caused by the external electric field.[32−35] Since the mechanisms involved are significantly different, some significant differences exist between electrically-induced nonequilibrium spin polarizations due to the kinetic magnetoelectric effect and those due to the spin Hall effect. While many theoretical efforts have been devoted recently to the understanding of the spin Hall effect in spin-orbit coupled semiconducting electronic systems, the kinetic magnetoelectric effect in such systems has not attractted much theoretical attention, and, especially, whether there may be some possible relation between this effect and some recent experimental findings remains an interesting and largely open problem. Moreover, in most of the previous studies about this effect, only the models with k-linear spin-orbit coupling (i.e., the spin-orbit interactions depend linearly on the momenta of charge carriers) were investigated in some detail, and the possible influences of boundary confinements were still neglected. But in many actual spin-orbit coupled semiconducting electronic systems, the spin-orbit interactions are usually not linearly momentum-dependent, and the properties of a spin-orbit coupled semiconducting electronic system with nonlinearly k-dependent spin-orbit interaction may be significantly different from those of a system with k-linear spin-orbit interaction.[40] A typical example of nonlinearly k-dependent spin-orbit interaction is the so-called k-cubic Rashba spin-orbit coupling, a model that captures most of the properties of two-dimensional hole gases. In the present paper we shall present a theoretical investigation on the characteristics of the kinetic magnetoelectric effect in laterally boundary-confined ballistic two-dimensional hole gases with k-cubic Rashba spin-orbit interaction. We will show that the electrically-induced nonequilibrium spin densities, due to the kinetic magnetoelectric effect in such systems, will exhibit some interesting characteristics. The characteristics found indicate that there may be some possible relation between this effect and some recent experimental findings observed in laterally boundary-confined two-dimensional hole gases.[8]

This paper is organized as follows: In Section 2 the model considered and the theoretical formulation employed will be briefly outlined. In Section 3 the main results obtained and the essential physical mechanism behind them will be discussed in some detail.

2. Model and theoretical formulation

The system considered is a laterally boundary-confined two-dimensional hole gas described by the following single-particle Hamiltonian,

$$\begin{aligned}\hat{H} = & -\frac{\hbar^2}{2m^*}\left(\frac{\partial^2}{\partial x^2}+\frac{\partial^2}{\partial y^2}\right) \\ & -\frac{\alpha}{\hbar}\left[\left(\frac{\partial}{\partial x}-\mathrm{i}\,\frac{\partial}{\partial y}\right)^3\hat{S}_+ - \left(\frac{\partial}{\partial x}+\mathrm{i}\,\frac{\partial}{\partial y}\right)^3\hat{S}_-\right] \\ & + V(x,y), \end{aligned} \tag{1}$$

where m^* is the effective mass of holes and $\hat{S}_\pm \equiv \hat{S}_x \pm \mathrm{i}\,\hat{S}_y$, with $\hat{S}_\gamma$ ($\gamma = x, y, z$) denoting the usual spin operators of spin 1/2. We have chosen the x-axis along the longitudinal direction and the y-axis along the transverse direction of the system and the z-axis along the normal of the two-dimensional plane. The second-term in the Hamiltonian (1) describes the usual k-cubic Rashba spin-orbit interaction, with the parameter α denoting the spin-orbit coupling strength.[40] The last term in the Hamiltonian (1) denotes the boundary-confinement potential. For simplicity and clarity, in this paper we consider only the ballistic transport limit where the effects of disorder scattering can be neglected. We also assume that the length of the system is much larger than the width so that the system can be considered as a translational invariant along the longitudinal direction. Under such a condition the confining potential $V(x, y)$ is independent of the longitudinal coordinates x. By making such simplifications, the essential physics contained in the mechanism can be more easily and transparently shown. As for the form of the confining potential, the usual hard-wall boundary condition will be used, i.e., the confining potential will be set to $V(x, y) = 0$ for $0 < y < L$ and $V(x, y) = \infty$ for $y < 0$ or $y > L$, where L is the width of the system and the two lateral boundaries of the system are assumed to be located at $y = 0$ and at $y = L$, respectively.

Since the system under study has translational invariance along the longitudinal direction, the longitudinal wave vector of charge carriers is a good quantum number. Hence all eigenfunctions of the Hamiltonian can be labeled by the longitudinal wave vector k_x, which will be denoted explicitly as $\Psi(k_x; x, y)$ below. In the absence of the spin-orbit coupling (i.e., $\alpha = 0$)

but in the presence of the hard-wall boundary condition, the eigenfunction of the n-th transverse mode is given by $\Psi^{(0)}_{n,\sigma}(k_x;x,y) = \sqrt{\frac{2}{L}}|\sigma\rangle \mathrm{e}^{\mathrm{i}\,k_x x}\sin\frac{n\pi y}{L}$, where $n\ (=1,\,2,\,\ldots\,,\,\infty)$ is the label of the transverse mode and $|\sigma\rangle$ denotes the usual spin-up ($\sigma = +$) or spin-down ($\sigma = -$) spinor states with respect to the z-axis. In the presence of spin-orbit coupling (i.e., $\alpha \neq 0$), quantum mixing between different transverse modes and between the spin-up and spin-down spinor states will be induced by the spin-orbit coupling, then the general form of the eigenfunctions of the Hamiltonian (1) can be expressed as

$$\Psi(k_x;x,y) = \sum_{n=1}^{\infty}\sum_{\sigma=\pm} C_{n,\sigma}\,\Psi^{(0)}_{n,\sigma}(k_x;x,y), \qquad (2)$$

where the coefficients $\{C_{n,\sigma}\}$ need to be determined by solving the Schrödinger equation. Inserting the wave function given by Eq.(2) into the Schrödinger equation: $\hat{H}\,\Psi(k_x;x,y) = E\,\Psi(k_x;x,y)$, where E is the eigenenergy, and making use of the orthogonality of the basis functions $\{\,\Psi^{(0)}_{n,\sigma}(k_x;x,y)\}$, one can obtain an infinite set of coupled linear equations for the coefficients $\{C_{n,\sigma}\}$ as follows:

$$a_m C_{m,+} + b_m C_{m,-} - \sum_{n(\neq m)}^{\infty} d_{n,m} C_{n,-} = E C_{m,+}, \quad (3)$$

$$a_m C_{m,-} - b_m C_{m,+} + \sum_{n(\neq m)}^{\infty} d_{n,m} C_{n,+} = E C_{m,-}, \quad (4)$$

where $n,\,m = 1,\,2,\,\ldots,\,\infty$, and $\{a_n\}$, $\{b_n\}$ and $\{d_{n,m}\}$ are defined by

$$a_n = \frac{\hbar^2}{2m^*}\left[k_x^2 + \left(\frac{n\pi}{L}\right)^2\right], \qquad (5)$$

$$b_n = \mathrm{i}\,\alpha k_x^3 - 3\,\mathrm{i}\,\alpha k_x\left(\frac{n\pi}{L}\right)^2, \qquad (6)$$

$$d_{n,m} = \left(\frac{\mathrm{i}\,\alpha n^3\pi^2}{L^3} - \frac{3\,\mathrm{i}\,\alpha k_x^2 n}{L}\right)\times\left[\frac{\cos(m+n)\pi}{m+n} + \frac{\cos(m-n)\pi}{m-n} - \frac{2m}{m^2-n^2}\right]. \qquad (7)$$

One can note that the coupled linear equations (3) and (4) can be rewritten in an explicit matrix form as follows:

$$\hat{\boldsymbol{A}}(k_x,\alpha,L)\cdot\hat{\phi} = E\hat{\phi}, \qquad (8)$$

where $\hat{\phi} \equiv [C_{1,+}, C_{1,-}, \ldots, C_{n,+}, C_{n,-}, \ldots]^T$ is a column vector composed of all the coefficients $\{C_{n,\sigma}\}$, and $\hat{\boldsymbol{A}}(k_x,\alpha,L)$ is a Hermitian matrix whose elements depend on the longitudinal wave vector k_x, the spin-orbit coupling strength α, and the system width L. Equation (8) means that the column vector $\hat{\phi}$ is an eigenvector of the matrix $\hat{\boldsymbol{A}}$, and E is the corresponding eigenvalue. Though the dimension of the matrix $\hat{\boldsymbol{A}}$ is infinite in principle, in actual calculations one can truncate the summation over the label n in Eq.(2) at a certain step so that the dimension of $\hat{\boldsymbol{A}}$ becomes finite. Such a truncation is appropriate in principle since the contributions of those transverse modes whose energies are much larger than the Fermi level of the system are negligibly small in the low temperature limit and in the linear response regime. After such a truncation is made, the matrix equation (8) can be solved conveniently by numerical procedures. For clarity, in the following the eigenvalues and the normalized eigenvectors of the matrix $\hat{\boldsymbol{A}}$ are denoted explicitly as $\epsilon_j(k_x)$ and $\hat{\phi}_j(k_x)$ ($j = 1,\,2,\,\ldots$), respectively, which are arranged in order of $\epsilon_1(k_x) \leq \epsilon_2(k_x) \leq \epsilon_3(k_x) \leq \ldots$. After a series of normalized eigenvectors $\{\hat{\phi}_j(k_x)\}$ are obtained, from Eq.(2) one can obtain correspondingly a series of normalized eigenfunctions of the Hamiltonian (1), which is denoted explicitly as $\Psi^{(j)}(k_x;x,y)$ ($j = 1,\,2,\,\ldots$) below. The position-dependent spin-expectation values corresponding to the eigenfunctions $\Psi^{(j)}(k_x;x,y)$ are given by

$$S^{(j)}_{\gamma}(k_x;y) = \Psi^{(j)}(k_x;x,y)^{\dagger}\hat{S}_{\gamma}\,\Psi^{(j)}(k_x;x,y), \qquad (9)$$

where $S^{(j)}_{\gamma}(k_x;y)$ ($\gamma = x,y,z$) denotes the x, y, or z component of the position-dependent spin-expectation values. (Since the system under study is translationally invariant along the longitudinal direction, the spin-expectation values are independent of the longitudinal coordinate.) The position-dependent net spin density due to all charge carriers (denoted as $\langle S_\gamma(y)\rangle$ below) is given by the ensemble average of the contributions from all charge carriers, which can be divided explicitly into two parts as follows:

$$\langle S_\gamma(y)\rangle = \langle S_\gamma(y)\rangle_+ + \langle S_\gamma(y)\rangle_-, \qquad (10)$$

where $\langle S_\gamma(y)\rangle_+$ denotes the contributions of charge carriers moving to the $+x$ direction and $\langle S_\gamma(y)\rangle_-$ denotes the contributions of charge carriers moving to the $-x$ direction, which are defined by, respectively,

$$\langle S_\gamma(y)\rangle_+ = \sum_j \int_0^{\infty}\frac{\mathrm{d}k_x}{2\pi} S^{(j)}_{\gamma}(k_x;y) f_+[\epsilon_j(k_x)], \quad (11)$$

$$\langle S_\gamma(y)\rangle_- = \sum_j \int_{-\infty}^{0}\frac{\mathrm{d}k_x}{2\pi} S^{(j)}_{\gamma}(k_x;y) f_-[\epsilon_j(k_x)]. \quad (12)$$

Here $f_+[\epsilon_j(k_x)]$ denotes the filling of charge carriers in the j-th sub-band and moving to the $+x$ direction, and $f_-[\epsilon_j(k_x)]$ the filling of charge carriers in the j-th sub-band but moving to the $-x$ direction. Due to the time-reversal symmetry of the spin-orbit interaction, one has $\epsilon_j(k_x) = \epsilon_j(-k_x)$ and $S_\gamma^{(j)}(-k_x; y) = -S_\gamma^{(j)}(k_x; y)$ in general. By use of these relations, the net spin densities given by Eqs.(10)–(12) can also be expressed as

$$\langle S_\gamma(y)\rangle = \sum_j \int_0^\infty \frac{\mathrm{d}k_x}{2\pi} S_\gamma^{(j)}(k_x; y) \times \{f_+[\epsilon_j(k_x)] - f_-[\epsilon_j(k_x)]\}. \tag{13}$$

From Eq.(13) one can see that the net spin density due to all charge carriers will vanish exactly in the equilibrium state (since $f_+[\epsilon_j(k_x)] \equiv f_-[\epsilon_j(k_x)]$ is in the equilibrium state). However, when a longitudinal charge current circulates through the system, the particle number density of charge carriers moving to the $+x$ direction is different from that of moving to the $-x$ direction (i.e., $f_+[\epsilon_j(k_x)] \neq f_-[\epsilon_j(k_x)]$), and due to this time-reversal symmetry-breaking, nonequilibrium spin polarizations may be induced, provided that the spin-expectation values given by Eq.(9) are nonvanishing. This is just the essential mechanism responsible for the generation of electrically-induced nonequilibrium spin densities in the kinetic magnetoelectric effect.[32−35] This mechanism is quite different from that of the spin Hall effect, where the electrically-induced nonequilibrium spin polarizations are caused by the accumulations of spins near the boundaries of a sample, which are driven to the boundaries by the flow of the transverse spin Hall current. Because the mechanisms involved are quite different, there will be some significant differences between electrically-induced nonequilibrium spin polarizations due to the kinetic magnetoelectric effect, and those due to the spin Hall effect. For example, according to the picture of the spin Hall effect, no nonequilibrium accumulations of spins could be electrically induced in an infinite and homogeneous spin-orbit coupled system without boundaries, i.e., in the spin Hall effect electrically-induced nonequilibrium spin densities can occur only near the boundaries of a boundary-confined spin-orbit coupled system.[3−4] In contrast to the spin Hall effect, according to the mechanism of the kinetic magnetoelectric effect, nonequilibrium spin densities can be electrically induced in an infinite and homogeneous spin-orbit coupled system without boundaries.[32−35]

3. Results and discussion

To calculate the electrically-induced nonequilibrium spin densities through the use of the formulation introduced in Section 2, one needs knowledge about the fillings of charge carriers in each energy sub-band in the presence of a longitudinal charge current. If we restrict ourselves to qualitative discussions, in the presence of a longitudinal charge current, the fillings of charge carriers in each energy sub-band can be characterized simply by introducing a pseudo-chemical potential shift in the Fermi–Dirac distribution function as follows: $f_\pm[\epsilon_j(k_x)] = f_0[\epsilon_j(k_x) - \epsilon_\mathrm{F} \pm \frac{1}{2}eV]$, where f_0 denotes the usual Fermi–Dirac distribution function, ϵ_F denotes the Fermi energy in the equilibrium state, and eV denotes the pseudo-chemical potential shift between the charge carriers moving to the $+x$ direction and those moving to the $-x$ direction. The longitudinal charge current circulating through the system can then be given by

$$J_x = e\sum_j \int_0^\infty \frac{\mathrm{d}k_x}{2\pi} v_x^{(j)}(k_x) \times \left\{ f_0\left[\epsilon_j(k_x) - \epsilon_\mathrm{F} + \frac{1}{2}eV\right] - f_0\left[\epsilon_j(k_x) - \epsilon_\mathrm{F} - \frac{1}{2}eV\right]\right\}, \tag{14}$$

where $v_x^{(j)}(k_x)$ denote the longitudinal velocities of charge carriers filling into the j-th sub-band with the longitudinal wave vector k_x. According to Eq.(13), in the linear response regime (i.e., $eV \ll \epsilon_\mathrm{F}$) the electrically-induced nonequilibrium spin densities are given by

$$\langle S_\gamma(y)\rangle = eV\sum_j \int_0^\infty \frac{\mathrm{d}k_x}{2\pi} S_\gamma^{(j)}(k_x; y)\frac{\partial f_0}{\partial \epsilon_\mathrm{F}}, \tag{15}$$

which is linearly proportional to V. Through this formula, the electrically-induced nonequilibrium spin densities can be calculated conveniently by numerical procedures. The numerical procedures are as follows: Firstly, we solve numerically the matrix equation (8) to obtain a series of eigenenergies $\{\epsilon_j(k_x)\}$ and a series of normalized eigenvectors $\{\hat{\phi}_j(k_x)\}$. Secondly, we insert the obtained eigenvectors $\{\hat{\phi}_j(k_x)\}$ into Eq.(2) to obtain a series of eigenfunctions $\{\Psi^{(j)}(k_x; x, y)\}$ of the Hamiltonian (1) and then calculate the spin-expectation values $\{S_\gamma^{(j)}(k_x; y)\}$ by use of Eq.(9). After $\{S_\gamma^{(j)}(k_x; y)\}$ are obtained, the spatial distribution of the electrically-induced nonequilibrium spin densities can be calculated through Eq.(15).

In Figs.1 and 2 the typical profiles of the transversal position dependence of the electrically-induced nonequilibrium spin densities obtained by using the above formulae are shown schematically. From Figs.1 and 2 one can see that, for the model considered, both in-plane polarized and normal-polarized nonequilibrium spin densities could be electrically induced by the kinetic magnetoelectric effect in the presence of boundary confinement, and more interestingly, from Fig.1 one can see that the normal-polarized nonequilibrium spin densities will manifest in a similar form to that due to the spin Hall effect: Firstly, the normal-polarized nonequilibrium spin densities occur near both lateral boundaries of the system; Secondly, the transversal spatial distributions of the normal polarized nonequilibrium spin densities are antisymmetric on both sides of the system, i.e., the spin polarizations are along opposite directions on both sides of the system. Such characteristics are very similar to that as conceived for electrically-induced nonequilibrium spin densities due to the spin Hall effect in a laterally boundary-confined spin-orbit coupled system, but the essential physical mechanisms behind them are much different, which will be discussed in more detail below. It should be noted that in some previous studies such characteristics are believed to be unique for the spin Hall effect in laterally boundary-confined spin-orbit coupled systems, thus they can serve as the principal observable signature of the spin Hall effect from the experimental point of view.[8] The results

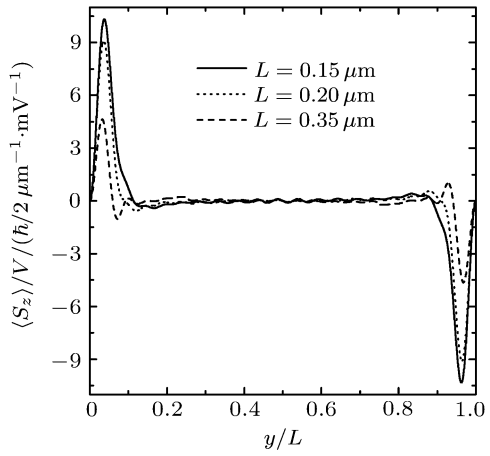


Fig.1. The transversal position dependence of the electrically-induced normal-polarized nonequilibrium spin densities (divided by V) due to the mechanism considered. The following parameters are used in the calculations: $m^* = 0.6m_e$, $\epsilon_F = 10$ meV, $\alpha = 10$ meV·nm^3. (The two boundaries of the system are assumed to be located at $y = 0$ and at $y = L$, respectively. The values of L are shown in the figure.)

obtained above indicate that such characteristics can also be manifested by the kinetic magnetoelectric effect in laterally boundary-confined spin-orbit coupled systems, i.e., they may not be unique for the spin Hall effect. Due to this fact and considering that a number of controversial issues about the spin Hall effect are still not resolved, in addition to the spin Hall effect, other possible effects and/or contributions should also be examined in order to gain a clearer understanding of some recent experimental findings.

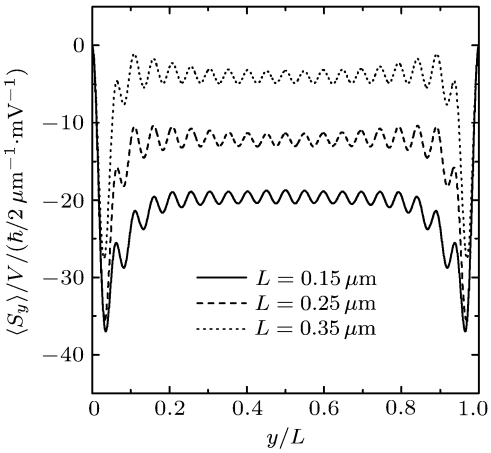


Fig.2. The transversal position dependence of the electrically-induced in-plane polarized nonequilibrium spin densities (divided by V) due to the mechanism considered. The parameters used in the calculations are the same as those in Fig.1.

It should be pointed out that, according to the usual physical picture of the kinetic magnetoelectric effect, if the effects of boundary confinement are not considered, only in-plane polarized nonequilibrium spin densities could be electrically induced by the kinetic magnetoelectric effect in two-dimensional spin-orbit coupled systems.[32–35] The physical mechanism behind this is that the effective magnetic fields felt by the charge carriers due to the spin-orbit coupling are usually in-plane orientated in two-dimensional spin-orbit coupled systems, and due to this fact, in the absence of boundary confinement, the z component of the spin-expectation values will vanish identically for all eigenstates of charge carriers. For example, for the infinite k-cubic Rashba model without the presence of boundary-confinement, the spinor eigenfunctions take the following simple form: $|\boldsymbol{k}\lambda\rangle = \frac{1}{\sqrt{2}}\begin{bmatrix} \mathrm{i}\,\lambda k_-^3/k^3 \\ 1 \end{bmatrix}$ (here $\lambda = \pm$ and $k_- \equiv k_x - \mathrm{i}\,k_y$), from which one can see that the z component of the spin-expectation values vanishes identically. But for a real system with boundaries, the effects of boundary confinement may

be important. After the effects of boundary confinement are taken into account, an important consequence is that the spinor eigenstates of the total Hamiltonian of the system under study may no longer be in-plane spin-orientated, i.e., the spinor wave functions of charge carriers may be distorted by the substantial spatial variation of the confining potential near the boundaries, and the normal components of the spin-expectation values given by Eq.(9) may hence become nonvanishing. Indeed, one may check that the spinor wave functions $|\boldsymbol{k}\lambda\rangle$ just mentioned are no longer the spinor eigenfunctions of the total Hamiltonian of the system under study in the presence of a spatially non-uniform confining potential, and due to the distortions of the spinor wave functions, not only the in-plane component but also the normal component of the spin-expectation values given by Eq.(9) will become nonvanishing. Due to this fact, from Eq.(13) one can see that both in-plane polarized and normal-polarized nonequilibrium spin densities could be electrically induced when a longitudinal charge current circulates through the system. Since the distortions of the spinor wave functions of charge carriers are caused by the substantial spatial variations of the confining potential near the boundaries, if the width of the system is sufficiently large, such distortions will become negligibly small at positions far away from both lateral boundaries. This means that if the width of the system is sufficiently large (e.g., for the cases shown in Fig.1), the normal-polarized nonequilibrium spin densities will be mainly localized near both lateral boundaries of the system, similar to that discerned in the spin Hall effect. According to the above physical picture, the characteristic localization length for the normal-polarized boundary spin polarizations is determined roughly by the width of the boundary regions in which the confining potential varies substantially.[1)] To obtain a quantitative estimate on the magnitude of the electrically-induced nonequilibrium boundary spin polarizations due to the mechanism discussed above so as to see whether they could be detected experimentally in an actual sample, we take two-dimensional p-doped GaAs as an example. For a strip sample of two-dimensional p-doped GaAs with the typical sample parameters $m^* \sim 0.6m_{\rm e}$, $\epsilon_{\rm F} \sim 10$ meV, $\alpha \sim 10$ meV$\cdot$nm^3, $V \sim 0.1$ mV, and $L \sim 10$ μm, the maximum of the electrically-induced nonequilibrium boundary spin polarizations due to the mechanism discussed above is around the order $0.1\% \sim 1\%$, which is within the experimentally detectable range by current technology.[7−12]

Unlike the normal-polarized nonequilibrium spin densities, from Fig.2 one can see that the in-plane polarized nonequilibrium spin densities are not mainly localized near the two lateral boundaries of the system, though their magnitudes can be enhanced substantially at some places near the two lateral boundaries. From the physical point of view, since the effective magnetic fields felt by the charge carriers due to the spin-orbit coupling are in-plane orientated, the in-plane components of the spin-expectation values given by Eq.(9) are nonvanishing in the absence of boundary confinement. This means that unlike the normal-polarized nonequilibrium spin densities, the in-plane polarized nonequilibrium spin densities are not caused by the distortions of the spinor wave functions of charge carriers near the boundaries, though their magnitudes can be enhanced substantially at some places near the boundaries. Another significant difference between the in-plane polarized and the normal-polarized nonequilibrium spin densities is that, unlike the normal-polarized nonequilibrium spin densities, the transversal spatial distributions of the in-plane polarized nonequilibrium spin densities are symmetric but not antisymmetric on both sides of the system, i.e., the spin polarizations are along the same direction on both sides of the system. This difference can be easily understood based on the intrinsic symmetry properties of the Hamiltonian of the system under study. From the form of the Hamiltonian (1), one can see that if the transversal spatial variation of the confining potential is symmetric on both sides of the system (i.e., $V(y) = V(L-y)$), the total Hamiltonian of the system will be invariant under the combined transformation of a real space coordinate transformation of $y \to L-y$ and a spin space rotation of an angle π around the S_y axis, under which the spin operators will transform as $\hat{S}_z \to -\hat{S}_z$ but $\hat{S}_y$ remains unchanged. From this transformation invariance, one can see immediately that the spin-expectation values given by Eq.(9) satisfy the following symmetry relations: $S_y^{(j)}(k_x; x, y) = S_y^{(j)}(k_x; x, L-y)$ and

[1)] It should be noted that in a real system the spatial variation of the confining potential is not so sharply localized at the boundaries as described by the hard-wall boundary condition, which is assumed in our calculations. Due to this fact, the characteristic localization length of the induced boundary spin polarizations in a real system should be substantially larger than that shown in Fig.1.

$S_z^{(j)}(k_x; x, y) = -S_z^{(j)}(k_x; x, L-y)$. Due to these intrinsic symmetry properties, from Eq.(13) one can see that the transversal spatial distributions of the in-plane polarized nonequilibrium spin densities are symmetric but the normal-polarized nonequilibrium spin densities are antisymmetric on both sides of the system, in agreement with the numerical results shown in Figs.1 and 2. Finally, it should be pointed out that the effect introduced above may also be present in other types of spin-orbit coupled semiconducting electronic systems, i.e., the mechanism discussed above is not limited to the k-cubic Rashba spin-orbit interaction in principle but may be general in other types of spin-orbital interactions (e.g., k-linear Rashba spin-orbit interaction or Dresselhaus spin-orbit interaction). But from the strictly theoretical point of view, in order to confirm this point unambiguously, we need to study the different types of spin-orbit interactions separately. This is beyond the scope of the present paper.

In summary, we have presented a theoretical investigation on the characteristics of the kinetic magnetoelectric effect in laterally boundary-confined ballistic two-dimensional hole gases with k-cubic Rashba spin-orbit coupling. We have shown that the electrically-induced nonequilibrium spin densities due to the kinetic magnetoelectric effect in such systems can exhibit some interesting characteristics. The physical mechanism behind the characteristics has been discussed in some detail. The characteristics we found indicate that there may be some possible relation between this effect and some recent experimental findings observed in laterally boundary-confined two-dimensional hole gases, which remains largely an open aspect at present and deserves more detailed theoretical investigations in the future.

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