The effect of *k*-cubic Dresselhaus spin–orbit coupling on the decay time of persistent spin helix states in semiconductor two-dimensional electron gases*

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We study the theoretical effect of *k*-cubic (i.e. cubic-in-momentum) Dresselhaus spin–orbit coupling on the decay time of persistent spin helix states in semiconductor two-dimensional electron gases. We show that the decay time of persistent spin helix states may be suppressed substantially by *k*-cubic Dresselhaus spin–orbit coupling, and after taking the effect of *k*-cubic Dresselhaus spin–orbit interaction into account, the theoretical results obtained accord both qualitatively and quantitatively with other recent experimental results.

Keywords: persistent spin helix state, Rashba spin–orbit coupling, *k*-cubic Dresselhaus spin–orbit coupling, decay time

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1. Introduction

In recent years, the understanding of the physics of semiconductor electronic systems with spin–orbit interaction (SOI) has attracted a great deal of theoretical and experimental interest due to their potential applications in the emerging field of semiconductor spintronics.[1–22] For example, the presence of SOI may provide a natural way of manipulating electron spins via purely electric means, without use of magnetic fields. The most prominent examples are the spin-Hall effect and current-induced spin polarization, both of which have received substantial theoretical and experimental interest in the past few years.[5–7] However, SOI also causes the undesirable effect of decay (relaxation) on the nonequilibrium spin polarizations of electrons. Since the functionality of most of the proposed semiconductor spintronic devices rely on the effective generation of long-lived nonequilibrium spin polarizations,[3,4] the decay time of nonequilibrium spin polarizations of electrons must be lengthened if we want to use them as a practical source for information processing. Thus the suppression of the relaxation of nonequilibrium spin polarizations in semiconductor electronic systems with SOI is a significant challenge in semiconductor spintronics.[9]

In the study of semiconductor electronic systems with SOI, a great deal of attention has been focused on semiconductor two-dimensional electron gases (2DEGs) formed in III–V type semiconductor quantum wells or heterostructures.[10–22] The breaking of inversion symmetry produces the Rashba SOI, and the strength of the Rashba SOI can be tuned by purely electric means.[23] In addition to the Rashba SOI, the structural inversion asymmetry of such systems also produces the Dresselhaus SOI.[24] For a semiconductor 2DEG with both Rashba and *k*-linear (i.e. linear-in-momentum) Dresselhaus SOI, the physics may exhibit some novelties when the strengths of the two types of SOI are equal. One particular interesting example is the prediction of the so-called persistent spin helix (PSH) states.[10] For 2DEGs formed in III–V type semiconductor quantum wells grown along the [001] direction, a PSH state is a helical spin density wave with the wavevector along the [110] direction. For such a helical spin density wave, it was first predicted in Ref. [10] that, if only the Rashba and *k*-linear Dresselhaus SOI is considered, a new type of exact $SU(2)$ symmetry will emerge when the Rashba and *k*-linear Dresselhaus spin–orbit coupling constants are equal and the magnitude of the wavevector is set to a certain characteristic value. This $SU(2)$ symmetry will significantly enhance the decay time of a helical spin density wave with this characteristic wavevector. Such a helical spin density wave is thus termed a persistent spin helix (PSH) state. The existence of such PSH states has been demonstrated by means of transient spin-gating spectroscopy in GaAs/AlGaAs quantum wells.[11,12] The tunable suppression of the relaxation of such PSH states may find some practical applications in semiconductor spintronics.

In most of the recent theoretical and experimental studies on PSH states in semiconductor 2DEGs, only the effects of Rashba and *k*-linear Dresselhaus SOI have been investigated in some detail.[10–15] But in a real semiconductor 2DEG system, the structural inversion asymmetry allows not only *k*-linear

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Dresselhaus SOI, but also k-cubic (i.e. cubic-in-momentum) Dresselhaus SOI.[24] An interesting question about PSH states is whether they are stable in the presence of k-cubic Dresselhaus SOI or whether their basic properties will be affected significantly by k-cubic Dresselhaus SOI. The understanding of what influences k-cubic Dresselhaus SOI may have on such PSH states is important, not only for their practical applications, but also for the clarification of some important discrepancies between the theoretical predictions and the relevant experimental results. In this paper we will present a detailed theoretical investigation, based on the spin-diffusion approach, on the effect of k-cubic Dresselhaus SOI on PSH states in semiconductor 2DEG systems formed in III–V type semiconductor quantum wells. We will show that the decay time of PSH states in such systems may be suppressed substantially by k-cubic Dresselhaus SOI. After taking into account the effect of k-cubic Dresselhaus SOI, the theoretical results obtained accord both qualitatively and quantitatively with some recent experimental results.

The paper is organized as follows: In Section 2 we give a brief description on the theoretical model and the theoretical formalism used in this paper. The results obtained will be presented and discussed in some detail in Section 3. Finally, a summary is given in the last section.

2. Model and formalism

We consider a semiconductor 2DEG system of noninteracting electrons in a III–V type semiconductor quantum well grown along the [001] axis. We choose the [100] axis as the x axis, the [010] axis as the y axis, and the [001] axis as the z axis. In the presence of both Rashba, k-linear, and k-cubic Dresselhaus SOIs, the single-particle effective-mass Hamiltonian of such a semiconductor 2DEG system reads

$$\begin{aligned}\hat{H} &= \frac{\hbar^2 k^2}{2m} + \alpha_{\rm R}(k_y\hat{\sigma}_x - k_x\hat{\sigma}_y) + \alpha_{\rm D_1}(k_x\hat{\sigma}_x - k_y\hat{\sigma}_y) \\ &\quad + \alpha_{\rm D_3}(k_x k_y^2\hat{\sigma}_x - k_y k_x^2\hat{\sigma}_y) + V_i(r), \end{aligned} \tag{1}$$

where m is the effective mass of electrons, $\boldsymbol{k} = (k_x, k_y)$ is the two-dimensional wavevector, $\hat{\boldsymbol{\sigma}} = (\hat{\sigma}_x, \hat{\sigma}_y, \hat{\sigma}_z)$ are the Pauli matrices, $\alpha_{\rm R}$, $\alpha_{\rm D_1}$, and $\alpha_{\rm D_3}$ denote the coupling constants of Rashba, k-linear, and k-cubic Dresselhaus SOIs, respectively, and $V_i(\boldsymbol{r})$ denotes a random spin-independent impurity potential. We assume that the random impurity potential is short-ranged, implying that the fluctuations of the impurity potential satisfy $\overline{\delta V(\boldsymbol{r})\delta V(\boldsymbol{r}')} = n_i v_0^2 \delta(\boldsymbol{r} - \boldsymbol{r}')$, with n_i being the density of impurities and v_0 the scattering amplitude.

For a disordered and noninteracting spin–orbit coupled 2DEG system, the spin dynamics, at much longer length scales than the mean-free path of electrons, can be described by a set of coupled spin-diffusion equations. These equations can be obtained by the standard density matrix formalism.[16–19] Assuming that a nonequilibrium spin density distribution is created at the initial time $t = 0$, then at length scales much longer than the mean-free path of electrons, the nonequilibrium spin density distribution at time $t > 0$ can be expressed as

$$S_\alpha(\boldsymbol{r}, t) = \int {\rm d}\boldsymbol{r}' P_{\alpha\beta}(t, \boldsymbol{r}, \boldsymbol{r}') S_\beta(0, \boldsymbol{r}'), \tag{2}$$

where $S_\beta(0, \boldsymbol{r}')$ is the initial spin density created at $\boldsymbol{r}'$ and $P_{\alpha\beta}(t, \boldsymbol{r}, \boldsymbol{r}')$ $(\alpha, \beta = x, y, z)$ is the diffusion Green's function. In the momentum space and for a homogeneous system, equation (2) can also be expressed as

$$[\hat{I} - \hat{P}(\omega, \boldsymbol{k})] \cdot \boldsymbol{S} = 0, \tag{3}$$

where $\hat{I}$ is the unit matrix and $\hat{P}(\omega, \boldsymbol{k})$ the Fourier transformation of the diffusion Green's function.

The diffusion Green's function can be calculated by using the disorder-averaged retarded (advanced) single-particle Green's function.[16] For short-ranged spin-independent random impurity scattering, the first Born approximation can be made in the calculation of the disorder-averaged retarded (advanced) single-particle Green's function. In the first Born approximation and in the momentum space, the 2×2 retarded (advanced) single-particle Green's function $\hat{G}^{\rm R(A)}$ satisfies the following equations of motion,[25]

$$\left[\omega - \hat{H}_0 \pm \frac{\rm i}{2\tau}\right] \cdot \hat{G}^{\rm R(A)}(\omega, \boldsymbol{k}) = \hat{I}, \tag{4}$$

where $\hat{H}_0$ $(= \hat{H} - V_i(r))$ is the disorder-independent part of the Hamiltonian (1) and $\tau = 1/n_i v_0^2 m$ is the momentum scattering time of electrons. From Eq. (4), the 2×2 retarded (advanced) single-particle Green's function $\hat{G}^{\rm R(A)}$ can be obtained. In the linear-response approximation, the diffusion Green's function can be obtained by using the 2×2 retarded (advanced) single-particle Green's function $\hat{G}^{\rm R(A)}$ through the following formula,[16]

$$\begin{aligned} P_{\alpha\beta}(\omega, \boldsymbol{k}) &= \frac{1}{4\pi\rho_{\rm F}\tau} \int \frac{{\rm d}^2\boldsymbol{q}}{(2\pi)^2} {\rm Tr}\Big[\hat{\sigma}_\alpha \hat{G}^{\rm R}(\omega/2, \boldsymbol{k} + \boldsymbol{q}/2) \\ &\quad \times \hat{\sigma}_\beta \hat{G}^{\rm A}(\omega/2, \boldsymbol{k} + \boldsymbol{q}/2)\Big], \end{aligned} \tag{5}$$

where ${\rm Tr}[\cdots]$ denotes the trace over spin indices and $\rho_{\rm F}$ the density of states at the Fermi level. Substituting the retarded (advanced) single-particle Green's function $\hat{G}^{\rm R(A)}$ obtained from Eq. (4) into Eq. (5), the Fourier transformations of the diffusion Green's function can be obtained. In this paper, we will concentrate on the limit of weak spin–orbit interactions, i.e. the momentum scattering time of electrons is much shorter than the spin-precession period due to spin–orbit interactions. This diffusive limit corresponds to the low frequency and long wavelength limit of Eq. (3). Taking the low frequency

and long wavelength limit of Eq. (3) and then transforming the equations back to real space, we get the following coupled spin-diffusion equations for the time-dependent and position-dependent spin densities:

$$\frac{\partial S_x}{\partial t} = D\nabla^2 S_x + AS_y + \left(B\frac{\partial}{\partial x} + C\frac{\partial}{\partial y}\right)S_z - \frac{S_x}{\tau_s}, \quad (6)$$

$$\frac{\partial S_y}{\partial t} = D\nabla^2 S_y + AS_x + \left(C\frac{\partial}{\partial x} + B\frac{\partial}{\partial y}\right)S_z - \frac{S_y}{\tau_s}, \quad (7)$$

$$\frac{\partial S_z}{\partial t} = D\nabla^2 S_z - \left(B\frac{\partial}{\partial x} + C\frac{\partial}{\partial y}\right)S_x - \left(C\frac{\partial}{\partial x} + B\frac{\partial}{\partial y}\right)S_y - \frac{2S_z}{\tau_s}. \quad (8)$$

Here $D = v_F^2\tau/2$ is the diffusion constant (with v_F the Fermi velocity), τ_s is a spin relaxation time defined by

$$\tau_s = \frac{\tau}{8(mD/\hbar)^2[\alpha^2 + (\beta_{D_1} + \beta_{D_3})^2 + \beta_{D_3}^2]}, \quad (9)$$

A, B, and C are three coefficients defined by

$$A = -\frac{8(mD/\hbar)^2\alpha(\beta_{D_1} + \beta_{D_3})}{k_F\tau[\alpha^2 + (\beta_{D_1} + \beta_{D_3})^2 + \beta_{D_3}^2]^{1/2}}, \quad (10)$$

$$B = 8(mD/\hbar)^2\alpha/k_F\tau, \quad (11)$$

$$C = 8(mD/\hbar)^2(\beta_{D_1} + \beta_{D_3})/k_F\tau, \quad (12)$$

where $\alpha = \hbar^2 k_F/m\alpha_R$, $\beta_{D_1} = \hbar^2 k_F/m\alpha_{D_1}$, $\beta_{D_3} = 4\hbar^2/k_F m\alpha_{D_3}$, with k_F denoting the Fermi wave vector.

The relaxation kinetics of a helical spin density wave can be described theoretically, based on the spin diffusion equations (6)–(8). For a helical spin density wave with a given wavevector $\boldsymbol{q}$, the time development of the position-dependent spin densities can be expressed as $\boldsymbol{S}(\boldsymbol{r},t) = \boldsymbol{s}(\boldsymbol{q},t)\,e^{i\boldsymbol{q}\cdot\boldsymbol{r}}$, where $\boldsymbol{s}(\boldsymbol{q},t)$ is the time-dependent factor. Then from Eqs. (6)–(8), the time-dependent factor $\boldsymbol{s}(\boldsymbol{q},t)$ will satisfy the following equation:

$$\frac{\partial \boldsymbol{s}(\boldsymbol{q},t)}{\partial t} = \hat{L}(\boldsymbol{q})\boldsymbol{s}(\boldsymbol{q},t), \quad (13)$$

where $\hat{L}(\boldsymbol{q})$ is a 3×3 matrix whose elements can be written down directly from Eqs. (6)–(8). In general, the 3×3 matrix $\hat{L}(\boldsymbol{q})$ has three eigenmodes. We denote the set of eigenvalues and the normalized eigenvectors of the three eigenmodes as $-\lambda_i(\boldsymbol{q})$ and $\boldsymbol{n}_i(\boldsymbol{q})$ $(i = 1,2,3)$, respectively. With these definitions, the solution of the time-dependent factor with a given initial condition $\boldsymbol{s}(\boldsymbol{q},t=0) = \boldsymbol{s}_0$ can be found by expanding $\boldsymbol{s}_0$ in terms of the set of the eigenvectors of $\hat{L}(\boldsymbol{q})$. The solution of $\boldsymbol{s}(\boldsymbol{q},t)$ is then of the form

$$\boldsymbol{s}(\boldsymbol{q},t) = \sum_{i=1,2,3} a_i\boldsymbol{n}_i(\boldsymbol{q})\,e^{-\lambda_i(\boldsymbol{q})t}, \quad (14)$$

where a_i $(i = 1,2,3)$ are the expansion coefficients. From Eq. (14) one can see that the decay rate of a helical spin density wave with a given wavevector $\boldsymbol{q}$ will be determined by the set of the eigenvalues of the 3×3 matrix $\hat{L}(\boldsymbol{q})$, and the decay time of the i-th eigenmode is defined by the inverse of the absolute value of the corresponding eigenvalue.

3. Results and discussions

Now we use the formalism introduced in Section 2 to study the decay time of PSH states in 2DEGs with both Rashba, k-linear, and k-cubic Dresselhaus SOIs. In accordance with the experimental setup of Refs. [11] and [12], we consider the propagation of helical spin density waves along the [110] direction in 2DEGs formed in GaAs/AlGaAs quantum wells grown along the [001] axis. The input material parameters used in the calculations are as follows: the electron effective mass m, the electron density n, the electron momentum scattering time τ, the Rashba and k-linear and k-cubic Dresselhaus spin–orbit coupling constants α_R, α_{D_1}, and α_{D_3}. The values of these input parameters will be set by making reference to the material parameters presented in Refs. [11] and [12]. The direction of the wavevector $\boldsymbol{q}$ will be fixed along the [110] direction and the Rashba and k-linear Dresselhaus spin–orbit coupling constants will be set to equal.

For clarity in making comparisons, we first present the results obtained by the formalism derived in Section 2, but neglecting k-cubic Dresselhaus SOI (i.e. setting $\alpha_{D_3} = 0$). In Fig. 1 we show the dependence of the decay times of the three eigenmodes of the spin diffusion equations on the magnitude q of the wavevector in the absence of k-cubic Dresselhaus SOI. From Fig. 1 we can see that the decay times of two eigenmodes decrease monotonously as the magnitude q of the wavevector increases. These are two ordinary modes. In addition to the two ordinary modes, from the spin diffusion equations we also obtain an anomalous mode whose features contrast significantly with those of the two ordinary modes. For this eigenmode, the decay time has a maximum value when the magnitude of the wavevector is equal to a special value q_{max}, i.e. the dispersion curve of the decay time in q-space shows a sharp peak at $q = q_{max}$. The characteristic value q_{max} of the magnitude of the wavevector is determined by the equal Rashba and k-linear Dresselhaus spin–orbit coupling constants, with $q_{max} \approx 4m\alpha_R/\hbar^2 = 4m\alpha_{D_1}/\hbar^2$. In the region of $q < q_{max}$, the decay time of the anomalous mode increases rapidly as q is increased towards q_{max}, and in the region of $q > q_{max}$, the decay time decreases rapidly as q is increased further. Since the decay time of this anomalous mode at $q \approx q_{max}$ is significantly larger than that of the two ordinary modes, in the vicinity of $q \approx q_{max}$ the decay rate of a helical spin density wave propagating along the [110] direction will be determined principally

by the decay time of this anomalous mode. The nonequilibrium spin polarization state described by this anomalous mode is just the PSH state proposed in Ref. [10]. Since, in the vicinity of $q \approx q_{\max}$, the decay rate of a helical spin density wave propagating along the [110] direction is determined principally by the decay time of the PSH mode, our discussion below will be focused on the properties of this eigenmode of the spin diffusion equations.

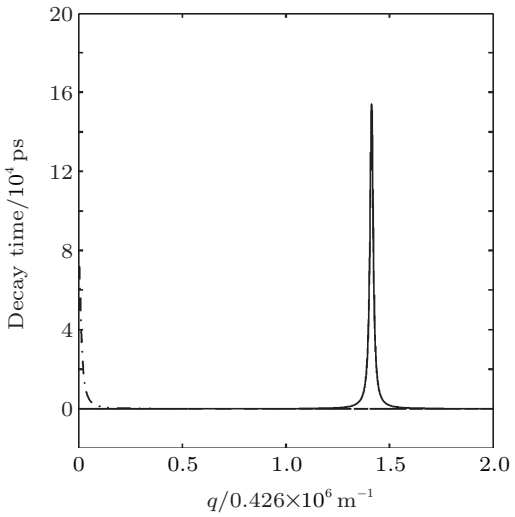


Fig. 1. The dependence of the decay time of the three eigenmodes of the spin diffusion equations on the magnitude q of the wavevector in the absence of the k-cubic Dresselhaus SOI. The direction of the wavevector $\boldsymbol{q}$ is fixed along the [110] direction, and by making reference to the material parameters presented in Refs. [11] and [12], the values of other input parameters used in the calculations are set to the following: $m = 0.067m_e$, $n = 7.8\times10^{11}$ cm^{-2}, $\tau = 5.7\times10^{-12}$ s, $\alpha_R = \alpha_{D_1} = 0.16\times10^{-10}$ eV·cm.

Though the theoretical results shown in Fig. 1 are in qualitative agreement with the recent relevant experimental observations,[11,12] we note that there are two important discrepancies between them. First, the maximum values of the decay time of the PSH mode obtained in the absence of k-cubic Dresselhaus SOI are much larger (about two orders of magnitude larger) than the corresponding experimental results. Secondly, the experimental results indicate that the variations of the decay time of a PSH state, with the variations of the magnitude of the wavevector, are much smoother but not so sharp as the theoretical results indicates (see, e.g. the experimental results shown in Fig. 2 of Ref. [12]). We find that these two important discrepancies can be accounted for by taking the effect of the k-cubic Dresselhaus SOI into account. From the spin diffusion equations derived in Section 2, we find that, in addition to two ordinary modes (whose decay time decreases monotonously as the magnitude of the wavevector increases), an anomalous mode can still be obtained in the presence of k-cubic Dresselhaus SOI, which shares much similarities with the PSH mode obtained in the absence of k-cubic Dresselhaus SOI. This is illustrated clearly in Figs. 2(a) and 2(b), where we show the dependence of the decay time of the anomalous mode on the magnitude q of the wavevector obtained in several different cases with $\alpha_{D_3} \neq 0$. From Figs. 2(a) and 2(b) one can see that the main features of the anomalous mode remain the same as that of the PSH mode shown in Fig. 1. After the effect of the k-cubic Dresselhaus SOI is taken into account, there is still a characteristic value $q_{\max}$ of the magnitude of the wavevector at which the decay time of the anomalous mode has a maximum value. In the region of $q < q_{\max}$, the decay time of the anomalous mode increases monotonously as q increases, and in the region of $q > q_{\max}$, the decay time decreases monotonously as

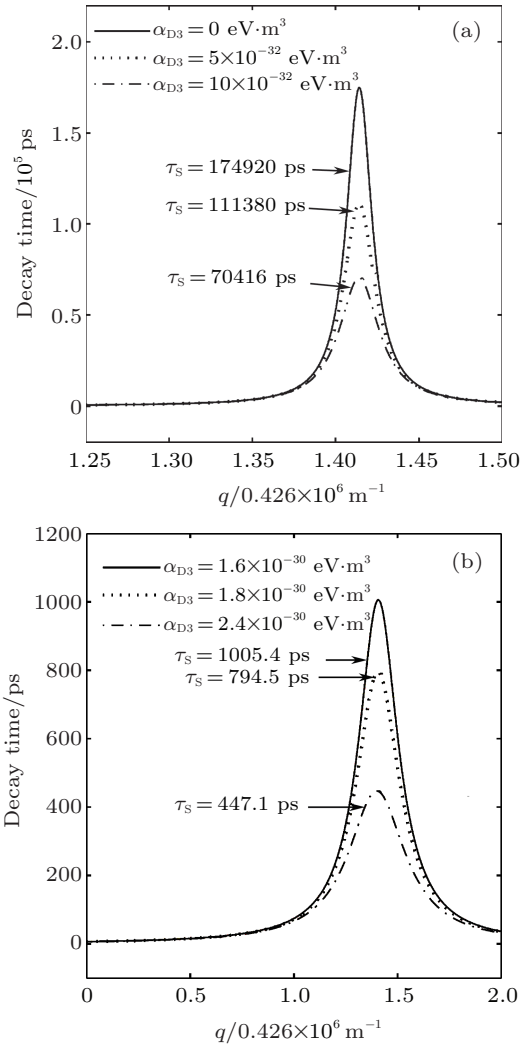


Fig. 2. (a) The q-dependence of the decay time of the PSH mode obtained in three different cases with $\alpha_{D_3} = 0\times10^{-32}$ eV·m^3, 5×10^{-32} eV·m^3, and 10×10^{-32} eV·m^3, respectively. (b) The q-dependence of the decay time of the PSH mode obtained in three different cases with $\alpha_{D_3} = 1.6\times10^{-30}$ eV·m^3, 1.8×10^{-30} eV·m^3, and 2.4×10^{-30} eV·m^3, respectively. The direction of the wavevector $\boldsymbol{q}$ is fixed along the [110] direction and the values of other input parameters used in the calculations are set identically to those given in the caption of Fig. 1. The peak values of the decay time of the PSH mode shown in Fig. 2(a) are about two orders of magnitude larger than that shown in Fig. 2(b).

q increases. But by making detailed comparison between the results shown in Fig. 1 and Figs. 2(a) and 2(b), one can see that there are two important discrepancies between them. First, after the effect of the k-cubic Dresselhaus SOI is taken into account, the variations of the decay time of the anomalous mode with the change of the magnitude of the wavevector will become smoother and smoother as the strength of the k-cubic Dresselhaus SOI increases, i.e. the peak of the dispersion curve of the decay time in q-space will be broadened significantly by the k-cubic Dresselhaus SOI. Second and more importantly, from Figs. 2(a) and 2(b) one can see that the maximum value of the decay time of the anomalous mode at $q = q_{\max}$ may be reduced significantly as the strength of the k-cubic Dresselhaus SOI increases, suggesting that the stability of a PSH state in a real semiconductor 2DEG system may be suppressed substantially by the k-cubic Dresselhaus SOI. By making a detailed comparison with some recent relevant experimental results (see, e.g. the experimental results shown in Fig. 2 of Refs. [11] and [12]), we find that, after taking into account the effect of the k-cubic Dresselhaus SOI, the theoretical results obtained by our formalism accord not only qualitatively but also quantitatively with the relevant experimental results. For example, from Fig. 2(b) one can see that, as the coupling constant of the k-cubic Dresselhaus SOI is increased to around $(1.5 \sim 2.5) \times 10^{-30}$ eV·m^3, typical for 2DEGs formed in III–V type semiconductor quantum wells,[13] the maximum value of the decay time of the PSH mode at $q = q_{\max}$ will be decreased to be around $(4 \sim 10) \times 10^2$ ps, in quantitative accordance with the relevant experimental results reported in Refs. [11] and [12]. From Figs. 2(a) and 2(b) we also note that, while the maximum value of the decay time of the anomalous mode will be reduced significantly as the strength of the k-cubic Dresselhaus SOI increases, the characteristic value $q_{\max}$ of the magnitude of the wavevector (at which the decay time of the anomalous mode has the maximum value) is insensitive to the variation of the strength of the k-cubic Dresselhaus SOI but is still determined mainly by the equal Rashba and k-linear Dresselhaus spin–orbit coupling constants. The characteristic value $q_{\max}$ of the magnitude of the wavevector obtained by our formalism, which is around the order of $(0.5 \sim 1.0) \times 10^6$ m^{-1} for the cases shown in Fig. 2(b), is also in accordance with the relevant experimental data.[11,12]

From a strict theoretical point of view, for a 2DEG formed in a III–V type semiconductor quantum well grown along the [001] axis, the exact $SU(2)$ symmetry shown in Ref. [10] can occur only if the direction of the wavevector of a helical spin density wave is strictly along the [110] direction. But it is not usually easy to arrange the direction of the wavevector strictly along the [110] direction in an actual experimental setup. Thus it is interesting to investigate whether the PSH mode can still survive if the direction of the wavevector is slightly deviated from the [110] direction. From the spin diffusion equations derived in Section 2, we find that if the direction of the wavevector is only slightly deviated from the [110] direction, an anomalous mode similar to the PSH mode discussed above can still be obtained. This is illustrated in Fig. 3, where we show the q-dependence of the decay time of the anomalous mode when the direction of the wavevector is slightly deviated from the [110] direction. From Fig. 3 one can see that the main features of the anomalous mode are similar to those of the PSH mode discussed above, but the peak value of the decay time of the anomalous mode decreases rapidly as the direction of the wavevector is turned away from the [110] direction.

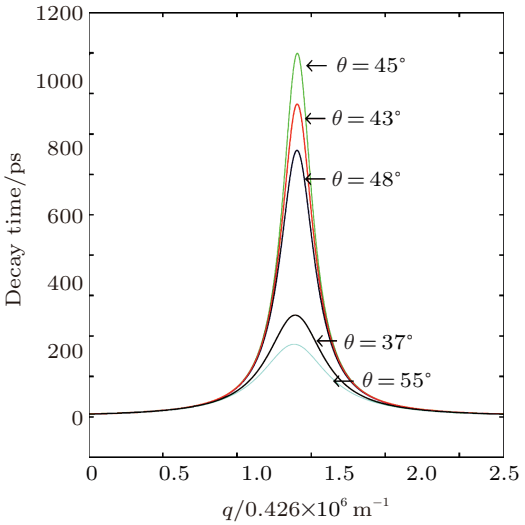


Fig. 3. (color online) The q-dependence of the decay time of the PSH mode as the direction of the wavevector is slightly deviated from the [110] direction. In the plot θ denotes the angle between the direction of the wavevector and the [100] axis. The values of other input parameters used in the calculations are given by: $m = 0.067m_e$, $n = 7.8 \times 10^{11}$ cm^{-2}, $\tau = 5.7 \times 10^{-12}$ s, $\alpha_R = \alpha_{D_1} = 0.16 \times 10^{-10}$ eV·cm, $\alpha_{D_3} = 1.7 \times 10^{-30}$ eV·m^3.

4. Summary

In summary, we have presented a detailed theoretical investigation on the effect of k-cubic Dresselhaus SOI on the decay time of PSH states in 2DEGs with both Rashba, k-linear and k-cubic Dresselhaus SOIs. The results indicate that the stability of a PSH state may be suppressed substantially by k-cubic Dresselhaus SOI, and after taking the effect of the k-cubic Dresselhaus SOI into account, the theoretical results obtained accord but qualitatively and quantitatively with some recent experimental results.

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