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2008 J. Phys.: Condens. Matter 20 465207

(http://iopscience.iop.org/0953-8984/20/46/465207)

IOP PUBLISHING JOURNAL OF PHYSICS: CONDENSED MATTER

J. Phys.: Condens. Matter **20** (2008) 465207 (5pp) doi:10.1088/0953-8984/20/46/465207

Phonon-induced decoherence of spin–orbit-driven coherent oscillations in a single InGaAs quantum dot

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Received 27 June 2008, in final form 18 September 2008
Published 21 October 2008
Online at stacks.iop.org/JPhysCM/20/465207

Abstract

The effect of direct spin–phonon interactions on spin–orbit-driven coherent oscillations in a single quantum dot proposed by Debald and Emary (2005 *Phys. Rev. Lett.* **94** 226803) is investigated theoretically in terms of the perturbation treatment based on a unitary transformation. It is shown that the decoherence rate induced by acoustic phonons strongly depends on the spin–orbit coupling strength, the magnetic field strength and the dot size.

1. Introduction

The past few years have seen great progress in both theory and experiment of quantum computation [1, 2]. Quantum computation based on semiconductor materials has become one of the most hopeful and realizable technologies for computation [3, 4]. One critically important technology for quantum computation and information processing is manipulating electron spins [5–8]. Loss and DiVincenzo [9] have proposed a scheme to use the spin state of coupled single-electron quantum dots. Another method for realizing quantum computation is using excitonic Rabi oscillation in quantum dot (QD) systems. Zrenner *et al* [10] have demonstrated that such excitonic coherent oscillations can be converted into deterministic photocurrents in a quantum dot system. On tuning an external gate voltage, the electron generated in the QD can tunnel out into nearby contacts. This process generates a photocurrent signal that is a weakly disturbed probe of the coherent state of the system. Kato *et al* [11] successfully manipulated 2D electron spins by a gigahertz electric field. Recently, the spin–orbit interaction known as the Rashba effect has been considered to be a possible control of electron spin states via gate voltages [12–14]. More recently, Debald and Emary [15] have proposed an experimental scheme to observe a spin–orbit-driven Rabi oscillation in quantum dot systems. They considered the effects of spin–orbit interaction on electrons in a small, few-electron lateral quantum dot. However, they did not consider the influence of the environment on the coherent oscillations.

Decoherence of a quantum system due to coupling to the environment is a great obstacle in experiment and in practical usage. In QD systems, phonons play an important role in the decoherence process. In this work, we will consider the phonon effect, which is considered to play a significant role in such a spin–orbit-driven coherent oscillation. Zhu *et al* [16] have shown that the damping rate of excitonic oscillation in a nanocavity depends on the exciton–phonon coupling constant and the vacuum Rabi frequency in terms of the perturbation method based on a unitary transformation. On this basis, we investigate the effect of spin–phonon interaction on a single InGaAs quantum dot system with Rashba spin–orbit interaction, which can be manipulated by an external gate voltage [17–19]. We will study how spin–phonon interaction and magnetic field influence the oscillation and demonstrate that such a coherent oscillation is very sensitive to the magnetic field.

2. Theory

We consider a simple two-dimensional quantum dot with parabolic lateral confinement potential in a perpendicular magnetic field **B** which points along the z direction. Then the electron system can be described by the Hamiltonian [15],

$$H_s = \frac{(\mathbf{p} + \frac{e}{c}\mathbf{A})^2}{2m^*} + \frac{m^*}{2}\omega_0^2(x^2 + y^2) + \frac{1}{2}g\mu_B B\sigma_z, \qquad (1)$$

where $\mathbf{p}$ is the linear momentum operator of the electron, $\mathbf{A}(\mathbf{r}) = \frac{\mathbf{B}}{2}(-y, x, 0)$ is the vector potential in the symmetric gauge, ω_0 is the characteristic confinement frequency, and $\sigma = (\sigma_x, \sigma_y, \sigma_z)$ is the vector Pauli matrices. m^* is the effective mass of the electron and g its gyromagnetic factor. μ_B is the Bohr magneton. In the second quantized notation, equation (1) becomes

$$H_s = (a_x^+ a_x + a_y^+ a_y + 1)\hbar\widetilde{\omega} + \frac{\hbar\omega_c}{2\mathrm{i}}(a_x^+ a_y - a_x a_y^+) + \tfrac{1}{2} g\mu_B B\sigma_z, \qquad (2)$$

where $\omega_c = \frac{eB}{m^* c}$ and $\widetilde{\omega}^2 = \omega_0^2 + \frac{\omega_c^2}{4}$. If we set

$$a_+ = \frac{1}{\sqrt{2}}(a_x - \mathrm{i}a_y), \qquad a_- = \frac{1}{\sqrt{2}}(a_x + \mathrm{i}a_y), \qquad (3)$$

then the Hamiltonian (2) can be written as

$$H_s = n_+\hbar\omega_+ + n_-\hbar\omega_- + \tfrac{1}{2} g\mu_B B\sigma_z, \qquad (4)$$

where $\hat{n}_+ = a_+^+ a_+$, $\hat{n}_- = a_-^+ a_-$ and $|n_+ n_-\rangle \frac{1}{\sqrt{n_+! n_-!}}(a_+^+)^{n^+} (a_-^+)^{n^-}|0\rangle$. In what follows we include the spin–orbit interaction, which is described as the Rashba Hamiltonian in this system [20],

$$H_{\mathrm{so}} = -\frac{\alpha_R}{\hbar}\left[\left(\mathbf{p} + \frac{e}{c}\mathbf{A}\right) \times \sigma\right]_z, \qquad (5)$$

where α_R is the spin–orbit coupling constant, which can be controlled by the gate voltage in experiment. On substituting equation (3) into (5), then

$$H_{\mathrm{so}} = \frac{\alpha_R}{\widetilde{l}}[\gamma_+(\sigma_+ a_+ + \sigma_- a_+^+) - \gamma_-(\sigma_- a_- + \sigma_+ a_-^+)], \qquad (6)$$

where $\gamma_\pm = 1 \pm \frac{1}{2}(\widetilde{l}/l_B)^2$, $\widetilde{l} = (\hbar/m^*\widetilde{\omega})^{\frac{1}{2}}$ and $l_B = (\hbar/m^*\omega_c)^{\frac{1}{2}}$.

The Hamiltonians of the phonon bath and its coupling to the electron spin can be written as follows [9, 21]:

$$H_{\mathrm{ph}} = \sum_{\mathbf{q}} \hbar\omega_{\mathbf{q}} b_{\mathbf{q}}^\dagger b_{\mathbf{q}}, \qquad (7)$$

$$H_{\mathrm{ph-s}} = \tfrac{1}{2}\sigma_z \sum_{\mathbf{q}} \hbar M_{\mathbf{q}}(b_{\mathbf{q}}^\dagger + b_{\mathbf{q}}), \qquad (8)$$

where $b_{\mathbf{q}}^\dagger$ ($b_{\mathbf{q}}$) and $\omega_{\mathbf{q}}$ are the creation (annihilation) operator and energy of the phonons with wavevector $\mathbf{q}$, respectively. $M_{\mathbf{q}}$ is the spin–phonon coupling constant. The effects of the phonon bath are fully described by a spectral density

$$J(\omega) = \sum_{\mathbf{q}} |M_{\mathbf{q}}|^2 \delta(\omega - \omega_{\mathbf{q}}). \qquad (9)$$

Hence we obtain the total Hamiltonian of the electron and phonon bath:

$$\begin{aligned} H &= H_s + H_{\mathrm{so}} + H_{\mathrm{ph}} + H_{\mathrm{ph-s}} \\ &= \hbar\omega_+ a_+^+ a_+ + \hbar\omega_- a_-^+ a_- + \tfrac{1}{2} g\mu_B B\sigma_z \\ &\quad + \frac{\alpha_R}{\widetilde{l}}[\gamma_+(\sigma_+ a_+ + \sigma_- a_+^+) - \gamma_-(\sigma_- a_- + \sigma_+ a_-^+)] \\ &\quad + \sum_{\mathbf{q}} \hbar\omega_{\mathbf{q}} b_{\mathbf{q}}^\dagger b_{\mathbf{q}} + \tfrac{1}{2}\sigma_z \sum_{\mathbf{q}} \hbar M_{\mathbf{q}}(b_{\mathbf{q}}^\dagger + b_{\mathbf{q}}). \end{aligned} \qquad (10)$$

Performing a unitary rotation of the spin such that $\sigma_z \to -\sigma_z$ and $\sigma_\pm \to -\sigma_\mp$, we arrive at the Hamiltonian

$$\begin{aligned} H &= \hbar\omega_+ a_+^+ a_+ + \hbar\omega_- a_-^+ a_- - \tfrac{1}{2} g\mu_B \sigma_z \\ &\quad + \frac{\alpha_R}{\widetilde{l}}[\gamma_-(\sigma_+ a_- + \sigma_- a_-^+) - \gamma_+(\sigma_- a_+ + \sigma_+ a_+^+)] \\ &\quad + \sum_{\mathbf{q}} \hbar\omega_{\mathbf{q}} b_{\mathbf{q}}^\dagger b_{\mathbf{q}} - \tfrac{1}{2}\sigma_z \sum_{\mathbf{q}} \hbar M_{\mathbf{q}}(b_{\mathbf{q}}^\dagger + b_{\mathbf{q}}). \end{aligned} \qquad (11)$$

We now derive an approximation form of this Hamiltonian by borrowing the observation from quantum optics that the terms preceded by γ_+ in equation (11) are counter-rotating, and thus negligible under the rotating-wave approximation when the spin–orbit coupling is small compared to the confinement [15, 22].

$$\begin{aligned} H &= \hbar\omega_+ a_+^+ a_+ + \hbar\omega_- a_-^+ a_- + \tfrac{1}{2}|g|\mu_B\sigma_z + \lambda(\sigma_+ a_- \\ &\quad + \sigma_- a_-^+) + \sum_{\mathbf{q}} \hbar\omega_{\mathbf{q}} b_{\mathbf{q}}^\dagger b_{\mathbf{q}} - \tfrac{1}{2}\sigma_z \sum_{\mathbf{q}} \hbar M_{\mathbf{q}}(b_{\mathbf{q}}^\dagger + b_{\mathbf{q}}), \end{aligned} \qquad (12)$$

where $\lambda = \alpha_R\gamma_-/\widetilde{l}$. Since g is negative in InGaAs, we choose the absolute value $|g|$ of g.

In order to solve the above Hamiltonian, we rewrite it as follows:

$$H = H_n + H_I, \qquad (13)$$

where

$$H_n = \hbar\omega_+ a_+^+ a_+ + \hbar\omega_- a_-^+ a_- + \tfrac{1}{2}|g|\mu_B\sigma_z, \qquad (14)$$

$$H_I = \sum_{\mathbf{q}} \hbar\omega_{\mathbf{q}} b_{\mathbf{q}}^\dagger b_{\mathbf{q}} - \tfrac{1}{2}\sigma_z \sum_{\mathbf{q}} \hbar M_{\mathbf{q}}(b_{\mathbf{q}}^\dagger + b_{\mathbf{q}}) + \lambda(\sigma_+ a_- + \sigma_- a_-^+). \qquad (15)$$

In what follows we will use the method proposed by Zhu *et al* [16] to treat these Hamiltonians. Then we apply two canonical transformations to the Hamiltonian (15),

$$H' = \mathrm{e}^{\mathrm{i}H_n t} H_I\, \mathrm{e}^{-\mathrm{i}H_n t}, \qquad (16)$$

$$H'' = \mathrm{e}^{A} H'\, \mathrm{e}^{-A}, \qquad (17)$$

with the generator

$$A = \sum_{\mathbf{q}} \frac{M_{\mathbf{q}}}{2\omega_{\mathbf{q}}} \xi_{\mathbf{q}}(b_{\mathbf{q}}^\dagger - b_{\mathbf{q}})\sigma_z, \qquad (18)$$

where $\xi_{\mathbf{q}}$ is a variational parameter, which can be determined later. Then the transformed Hamiltonian is decomposed into three parts [16]:

$$H'' = H_0'' + H_1'' + H_2'', \qquad (19)$$

where

$$H_0'' = \eta\lambda(\sigma_+ a_- + \sigma_- a_-^+) + \sum_{\mathbf{q}} \omega_{\mathbf{q}} b_{\mathbf{q}}^\dagger b_{\mathbf{q}} - \sum_{\mathbf{q}} \frac{M_{\mathbf{q}}}{2\omega_{\mathbf{q}}} \xi_{\mathbf{q}}(2 - \xi_{\mathbf{q}}), \qquad (20)$$

$$H_1'' = \tfrac{1}{2}\sum_{\mathbf{q}} M_{\mathbf{q}}(1 - \xi_{\mathbf{q}})(b_{\mathbf{q}}^\dagger + b_{\mathbf{q}})\sigma_z - \eta\lambda(\sigma_+ a_- - \sigma_- a_-^+) \times \sum_{\mathbf{q}} \frac{M_{\mathbf{q}}}{\omega_{\mathbf{q}}} \xi_{\mathbf{q}}(b_{\mathbf{q}}^\dagger - b_{\mathbf{q}}), \qquad (21)$$

$$H_2'' = \lambda\sigma_+ a_- \left\{ \exp\left[\sum_{\mathbf{q}} \frac{M_{\mathbf{q}}}{\omega_{\mathbf{q}}} \xi_{\mathbf{q}} (b_{\mathbf{q}}^\dagger - b_{\mathbf{q}}) \right] - \eta \right\} + \lambda\sigma_- a_-^+ \left\{ \exp\left[-\sum_{\mathbf{q}} \frac{M_{\mathbf{q}}}{\omega_{\mathbf{q}}} \xi_{\mathbf{q}} (b_{\mathbf{q}}^\dagger - b_{\mathbf{q}}) \right] - \eta \right\} + \eta\lambda(\sigma_+ a_- - \sigma_- a_-^+) \sum_{\mathbf{q}} \frac{M_{\mathbf{q}}}{\omega_{\mathbf{q}}} \xi_{\mathbf{q}} (b_{\mathbf{q}}^\dagger - b_{\mathbf{q}}) \tag{22}$$

where η is a variational parameter which can be adjusted to minimize H_1'' and H_2''. It is obvious that H_0'' can be solved exactly because the electron spin and the phonons are decoupled. Following the method [16], we can obtain the variational parameters $\xi_{\mathbf{q}}$ and η respectively as follows:

$$\xi_q = \frac{\omega_{\mathbf{q}}}{\omega_{\mathbf{q}} + 2\eta\Omega_m}, \tag{23}$$

$$\eta = \exp\left(-\sum_{\mathbf{q}} \frac{M_{\mathbf{q}}^2}{2\omega_{\mathbf{q}}^2} \xi_{\mathbf{q}}^2 \right), \tag{24}$$

where $\Omega_m = \lambda\sqrt{m+1}$ ($m = 0, 1, 2, \ldots$). The population inversion of electron spin is given by

$$W(t) = \langle\{0_{\mathbf{q}}\}|\langle\uparrow|e^{iH''t}\sigma_z e^{-iH''t}|\uparrow\rangle|\{0_{\mathbf{q}}\}\rangle = \frac{1}{4\pi i}\oint \frac{e^{-iEt}\,dE}{E - 2\eta\Omega_m - \sum_{\mathbf{q}} \frac{V_{\mathbf{q}}^2}{E-\omega_{\mathbf{q}}+i0^+}} + \frac{1}{4\pi i}\oint \frac{e^{iEt}\,dE}{E - 2\eta\Omega_m - \sum_{\mathbf{q}} \frac{V_{\mathbf{q}}^2}{E-\omega_{\mathbf{q}}-i0^+}}, \tag{25}$$

where $|\{0_{\mathbf{q}}\}\rangle$ stands for the vacuum state of the phonon, $|\uparrow\rangle$ is the spin-up state and $V_{\mathbf{q}} = 2\eta\Omega_m M_{\mathbf{q}}\xi_{\mathbf{q}}/\omega_{\mathbf{q}}$. The real and imaginary parts of $\sum_{\mathbf{q}} \frac{V_{\mathbf{q}}^2}{E-\omega_{\mathbf{q}}\pm i0^+}$ are denoted as $R(E)$ and $\mp\gamma(E)$,

$$R(\omega) = -(2\eta\Omega_m)^2 \int_0^\infty \frac{J(\omega')\,d\omega'}{(\omega'-\omega)(\omega'+2\eta\Omega_m)^2}, \tag{26}$$

$$\gamma(\omega) = \pi(2\eta\Omega_m)^2 J(\omega)/(\omega + 2\eta\Omega_m)^2. \tag{27}$$

The spectral density of the phonon bath is described by [23]

$$J(\omega) = 2\alpha\omega^3 e^{-(\omega/\omega_l)^2}, \tag{28}$$

where α is the spin–phonon coupling constant and $\omega_l = c/L_0$ is the cut-off frequency (c is the sound speed and L_0 is the size of the quantum dot).

The integral in equation (25) can be evaluated and in the second order approximation the result is

$$W(t) = \cos(\omega_m t)\, e^{-\gamma_m t}, \tag{29}$$

where

$$\gamma_m = 4\pi\alpha(\eta\lambda)^3 (m+1)^{\frac{3}{2}} e^{-4(m+1)(\frac{\eta\lambda}{\omega_l})^2}, \tag{30}$$

and ω_m is the solution of the equation

$$\omega - 2\eta\Omega_m + (2\eta\Omega_m)^2 \int_0^\infty \frac{J(\omega')\,d\omega'}{(\omega'-\omega)(\omega'+2\eta\Omega_m)^2} = 0. \tag{31}$$

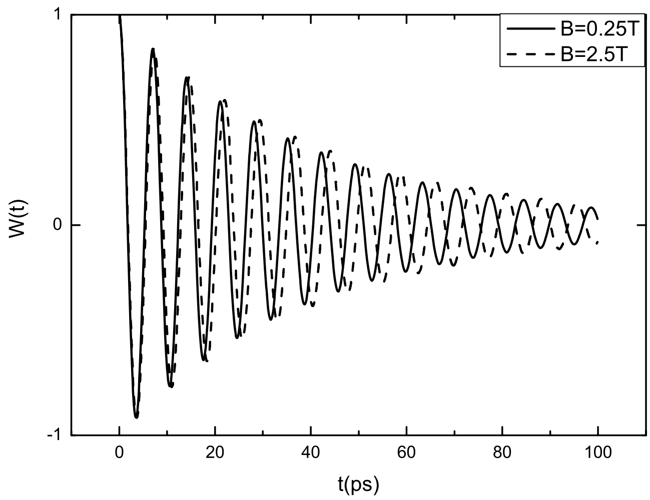

Figure 1. The population inversion as a function of time with two magnetic fields. The parameters used are $L_0 = 150$ nm, $\alpha = 0.05$ ps^2, $g = -4$ and $\alpha_R = 0.8 \times 10^{-12}$ eV m.

Equation (29) obviously stands for a damped coherent oscillation with frequency ω_m and damping rate γ_m. If there is no spin–phonon interaction ($\alpha = 0$), then $\gamma_m = 0$, $W(t) = \cos(2\Omega_m t) = \cos(2\lambda\sqrt{m+1}t)$, so the coherent oscillation in this system would not damp out as shown by Debald and Emary [15]. Since $\lambda = \alpha_R\gamma_-/\widetilde{l}$, we can alter the magnetic field strength to control the frequency of the oscillation and the damping rate. On the other hand, it is interesting in our analytical solutions that if we turn off the magnetic field the 'Rabi frequency λ' will not be equal to zero, then the oscillations still exist due to spin–orbit coupling.

For illustration of numerical results, we only consider the condition of orbit quantum number $m = 0$, then equations (29) and (30) become

$$W(t) = \cos(2\eta\lambda t)\, e^{-\gamma t}, \tag{32}$$

and

$$\gamma = 4\pi\alpha(\eta\lambda)^3 e^{-(2\frac{\eta\lambda}{\omega_l})^2}. \tag{33}$$

3. Results and discussion

In what follows we choose an InGaAs quantum dot with a dot size of 150 nm (approximate size for the dot in [15]), and other parameters are typical of InGaAs: $g = -4, \alpha = 0.05$ ps^2. Figure 1 presents the population inversion of electron spin as a function of time with two magnetic fields. Generally, the oscillation will damp out quickly when magnetic field strength B becomes smaller. The magnetic field strength in the broken curve is ten times larger than that of the solid line. In the first 10 ps, these two curves are overlapped, which means the damped times in this period are almost the same. After about 20 ps, the amplitude of each curve is not changed, but the delay of the broken curve gradually increases. Therefore, if a strong magnetic field such as superconducting magnetic field is applied, the damped time will be longer.

Figure 2 depicts the same behavior of population inversion of electron spin as in figure 1, while the only difference is that B is fixed and α_R is changed for three values. Grundler [18]

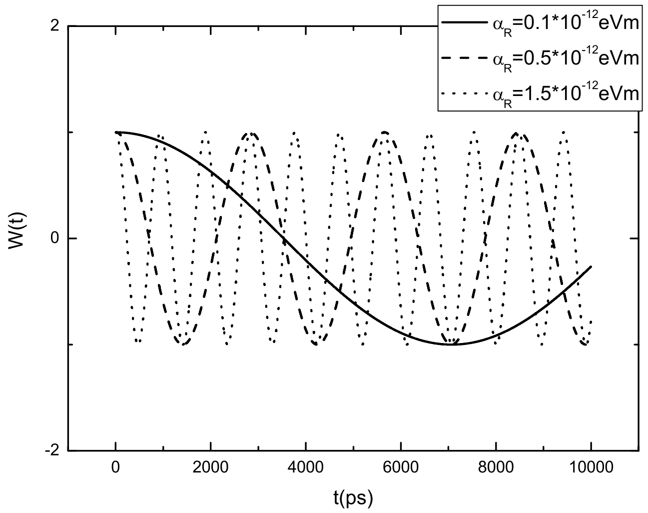

Figure 2. The population inversion as a function of time with three spin–orbit coupling constants. The parameters used are $L_0 = 150$ nm, $B = 90$ mT, $g = -4$ and $\alpha = 0.05$ ps^2.

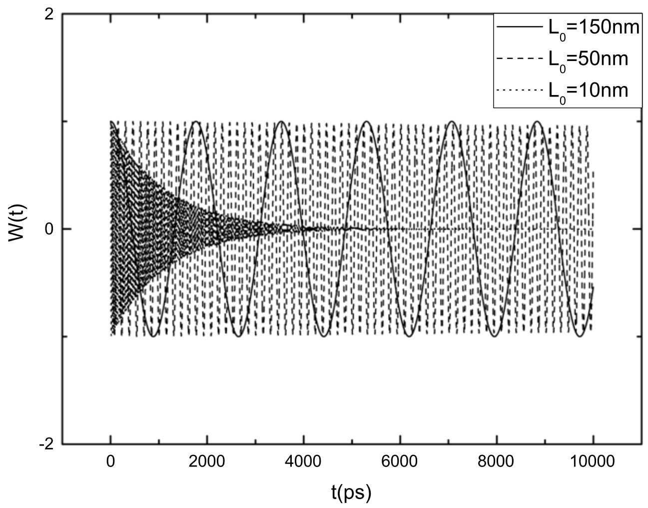

Figure 3. The population inversion as a function of time with three quantum dots. The parameters used are $B = 90$ mT, $\alpha = 0.05$ ps^2, $g = -4$ and $\alpha_R = 0.8 \times 10^{-12}$ eV m.

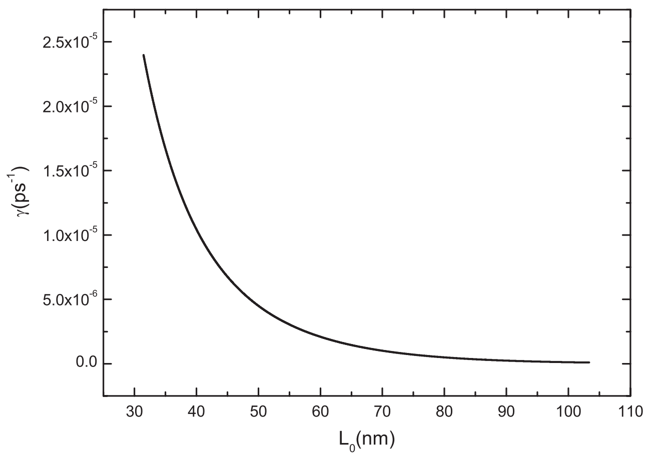

Figure 4. The damping rate γ as a function of L_0. The parameters used are $B = 90$ mT, $\alpha = 0.05$ ps^2, $\alpha_R = 0.8 \times 10^{-12}$ eV m and $g = -4$.

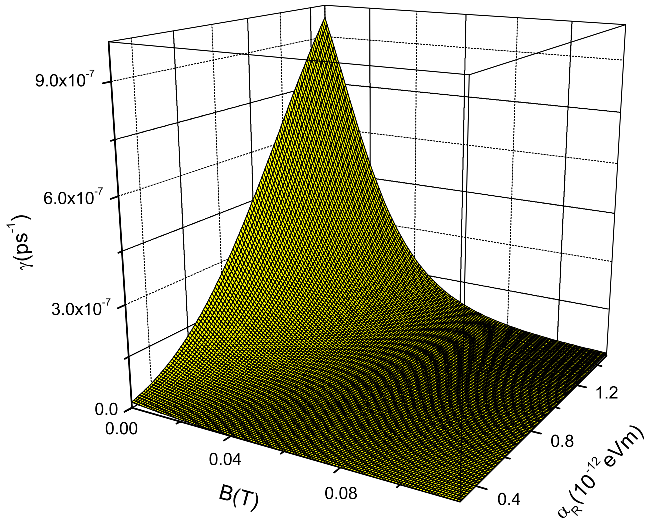

Figure 5. The damping rate γ as a function of magnetic field B and spin–orbit coupling interaction α_R. The dot size is chosen to be $L_0 = 150$ nm, $\alpha = 0.05$ ps^2 and $g = -4$.

(This figure is in colour only in the electronic version)

has shown that spin–orbit coupling constant α_R changes with electron density, which can be simply controlled by gate voltage. In a recent paper, Koga *et al* [19] have reported that α_R can be varied in the range 0.3×10^{-12}–1.5×10^{-12} eV m. The oscillation behaviors of these three curves are very similar, as shown in figure 2. A small increase of spin–orbit coupling gives rise to a relatively large change of the damped time. When $t = 10^4$ ps, the oscillation for $\alpha_R = 1.5 \times 10^{-12}$ eV m starts to damp; however, the solid curve ($\alpha_R = 0.1 \times 10^{-12}$ eV m) is not changed. This behavior stems from the parametric dependence of α_R in equations (32) and (33).

From figure 3, we can see that population inversion of electron spin also changes with the size of the QD. In this case, the magnetic field is set to $B = 90$ mT. The damped oscillations are greatly affected by the size of the QD. It is observed that for the case of a small quantum dot with $L_0 = 10$ nm the population inversion damps quickly, within about 5000 ps. But when the dot size is large as $L_0 = 150$ nm (just like the case in [15]) the damped time increases significantly.

Figures 4 and 5 demonstrate how damping rate changes with three parameters: dot size, magnetic field and spin–orbit interaction. The damping rate as a function of quantum dot size L_0 is shown in figure 4. When L_0 is smaller than 60 nm, the damping rate decreases quickly and approaches zero with increasing L_0. When L_0 becomes larger than 100 nm the decrease of γ slows down, and for $L_0 = 150$ nm γ is almost zero. This is because the cut-off frequency is inversely proportional to the dot size and equation (33) shows an exponential relation between the damping rate and the cut-off frequency. As a result, the dot size has a significant dependence on the damping rate. Equations (32) and (33) also indicate that the damping rate γ is greatly dependent on the spin–phonon coupling constant, spin–orbit coupling constant and magnetic field strength as the dot size is fixed.

From figure 5, it is obvious that by increasing the magnetic field or spin–orbit coupling the damping rate is reduced and the oscillations will persist for a long time. From these two figures, it is found that the characteristic decoherence time $T_2 = 1/\gamma$ is about 1 μs, which is long enough for quantum computation. Consequently, the damping rate of the spin–orbit-driven coherent oscillation can be effectively controlled by the magnetic field, dot size and spin–orbit interaction.

4. Conclusions

In summary, we have investigated the influence of spin–phonon interaction on the spin–orbit-driven coherent oscillation proposed by Debald and Emary [15]. The analytical results of coherent dynamics and the damping rate are obtained as functions of magnetic field strength, spin–phonon coupling constant, spin–orbit coupling constant and dot size. It is shown that the damping rate is reduced significantly for large QDs and strong magnetic fields. As the dot size and the spin–orbit coupling are fixed, the damping rate only depends on external magnetic fields, so in a realistic experiment we expect that a strong magnetic field will reduce damping rate and prolong such an oscillation lifetime. Finally, it should be noted in the present paper that a magnetic field component in the x–y plane or a tilt magnetic field will influence the decoherence rate due to the phonon bath as discussed by Golovach *et al* [24]. Here, for the sake of analytical simplicity, we only consider a magnetic field along the z direction. The effects of a magnetic field component in the x–y plane or a tilt magnetic field will certainly have an impact on the decoherence rate of electron spin. However, we cannot obtain the analytical results in such a general case by using the present treatment. The numerical methods for this case are underway.

Acknowledgments

Part of this work was supported by the National Natural Science Foundation of China (No 10774101) and the National Ministry of Education Program for Training PhD.

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