Magneto-quantum oscillations of the Korringa relaxation rate of manganese ion near a two-dimensional electron gas

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Abstract

The Mn ion-spin relaxation rate (Korringa relaxation) in the vicinity of the two-dimensional electrons (2DEG) in a Mn-based semiconductor nanostructures in a quantizing magnetic field was calculated. The Korringa relaxation is an energy-consuming process due to the difference in magnetic moments of localized and electrons spins involved. The mechanism of energy transfer between the Mn spin and the 2DEG is exchange scattering of the Landau electrons with a transition from the $e \uparrow$ sub-band to the $e \downarrow$ sub-band accompanied by a change in the Mn spin. It was found that due to the presence of Landau levels and the spin-split mobility gap, the Korringa relaxation rate oscillates with the magnetic field resembling the oscillations of the resistivity in such systems. Our calculation offers a method of investigating the dynamics of a magnetic ion such as Mn in a 2DEG and provide new information on the exchange parameter as well as information about the 2D electrons themselves.

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1. Introduction

Dilute magnetic semiconductor [1] (DMS) provide unique opportunities for manipulation of carrier spins due to the strong sp-d exchange interaction between the carrier spins and localized spins of magnetic ions. A variety of magnetic, magneto-optical, and magnetotransport phenomena [2–4] makes DMS’s a very reliable model system for spintronic devices [5,6]. It was demonstrated very recently that a layer of II–VI DMS, deposited on top of a GaAs-based n-i-p light emitting diode, serves as an effective injector of spin-polarized carriers [7]. A prerequisite for such structures relying on electrical spin injections is the presence of free carriers achieved by doping of the DMS layer. However, the experimental information on the influence of free carriers on the dynamical magnetic properties of wide-gap DMS’s [e.g. (Cd,Mn)Te or (Zn,Mn)Se] is very limited. It is known for metals and bulk semiconductors with magnetic impurities that the free carriers play an important role for energy transfer away from magnetic ions [8,9]. The Korringa relaxation, which is an energy-consuming process due to the difference in magnetic moments of localized and electrons spins involved, is an example of the effect caused by such interaction. Similar effect has been reported for bulk narrow-gap DMS’s with a high concentration of free carriers [10]. However, a study of this phenomenon has not been extended, to our knowledge, to wide-gap DMS’s. One can expect that the energy and spin transfer between the systems of the carriers and magnetic ions can be modified significantly by the presence of a background of free carriers. One should also expect strong modifications in the case of low-dimensional electron systems realized in quantum wells, wires and dots, as the density of states in the nanostructures differs qualitatively from that of the three-dimensional (3D) system. It was reported recently that in (Cd, Mn)Te-based quantum wells(QW’s) the magnetic-ion system can be substantially overheated by means of the interaction with photocarriers that is enhanced by the presence of a two-dimensional electron gas(2DEG) [11]. This result shows that the interaction of magnetic ions with 2DEG is quite strong and that the modification of their spin dynamics might be very significant.

On the other hand, it has been reported [12] that in a strong magnetic field of say 147, separate Landau levels are resolved in the photoluminescence spectra, which indicates...
that the energy distribution of electrons has notable oscillations. As it is well known, for the ideal Landau-quantized 2DEG in a magnetic field, the energy spectrum is purely discrete. The energy-conservation law, therefore, cannot be satisfied in the spin-flip process under discussion and no Korringa relaxation is expected. In a real 2DEG, however, the electrons move in a fluctuating external potential created by randomly distributed impurities and or crystal imperfections. This potential lifts the degeneracy of the energy levels, and spin-flip processes then becomes possible since the energy can be conserved by a suitable change of the electronic orbital state. That is, the Landau levels acquire a homogeneous broadening and the DOS is described by a Gaussian with broadening parameter $\Gamma$. It is therefore, important to explore here in this paper the question of how dynamic magnetic properties of DMS’s QW’s, namely the Korringa relaxation (KR) rate of magnetic ions are modified by the presence of a 2DEG in a high (quantizing) magnetic field, and how this rate depends on magnetic field strength.

2. Korringa relaxation rate

We shall perform the Mn-spin relaxation rate due to interaction with the Q2DEG in the QW in the presence of a high (quantizing) magnetic field. In order to make the problem very simple and understandable we initially assume that the external magnetic field is sufficiently low that the Landau quantization of electron orbital motion can be ignored. We then find the important equations leading to the Mn spin-flip transition probability and then specialize the result to the quantizing magnetic field.

Suppose that there is thermal equilibrium between spin subbands. Let the weak magnetic field be aligned with the $z$-axis, perpendicular to the QW plane. The well occupies the space between $z = 0$ and $L$, and a Mn atom is at distance $d$ from the heterojunction. Suppose for simplicity that the walls of the QW are infinitely high. Such an approximation is supported by a recent band offsets calculation [13] in CdTe/(Cd,Mn)Te quantum wells in which a large conduction-band offset of 400 meV (the band gap of Cd$_1-x$Mn$_x$Te is linear with $x$) [14] was deduced for a Mn molar fraction $y = 0.24$ in the barrier layer with an accuracy better than 10 meV. A finite height of the walls should lead to inessential corrections that can be ignored in our case.

We shall begin by writing down the Hamiltonian of the s–d exchange model has the form [3]

$$H_{s-d} = \frac{J}{N_0} \sigma \cdot \hat{S}(\vec{r}) = \frac{1}{2N_0} \int \hat{j}(\vec{r})[2\sigma^+ S^+ + (\sigma^+ S^- + \sigma^- S^+)].$$

(1)

where the band electrons have a magnetic interaction with the well-localized $S = \frac{5}{2}$ half filled-3D shell. Here $J$ is the exchange integral for conduction electrons, $N_0$ is the number of Mn atoms in the quantum well, and the spin operators $\hat{S}$ and $\sigma$ act on the Mn and free electrons spins, respectively. The spin wave function of Mn varies as a result of an electron transition from one spin subband to another simultaneously with a change in the spin projection of the magnetic impurity by a unit. Consider a situation when the magnetic field is sufficiently strong not to cause orbital quantization, so that Mn has in the initial state the spin $S_z$ and in the final state the spin $S_z'$. Hamiltonian (1) is responsible for transition from state 5/2 to state 3/2, or in general from $M$ to $M \pm 1$. The transition probability of such a transition is given by the ‘Fermi golden rule’ [15]

$$w = \frac{2\pi}{\hbar}|\langle f|H_{s-d}|i\rangle|^2 \delta(E_f - E_i).$$

(2)

We consider now electrons from the lowest size-quantized subbands so that the wave functions of the initial and final states have the form

$$|i\rangle = \phi(d)|\vec{k}\rangle, \quad \left| \frac{5}{2} - \frac{1}{2} \right\rangle$$

(3)

$$|f\rangle = \phi(d)|\vec{k}'\rangle, \quad \left| \frac{3}{2} + \frac{1}{2} \right\rangle,$$

(4)

where

$$\phi(z) = \frac{\sin(\pi z/L)}{\sqrt{L}}.$$

(5)

is the envelope part of the wave function defining the width of the 2DEG perpendicular to the interface. Here $\vec{k}$ and $\vec{k}'$ are wave vectors of Bloch electron functions in the initial and final states in the QW plane

$$|\vec{k}\rangle = \hat{a}\vec{k}(\vec{r}) \frac{e^{\vec{k} \cdot \vec{r}}}{\sqrt{A}},$$

(6)

where $A$ is the normalizing area of the QW plane and $\hat{a}\vec{k}(\vec{r})$ is the Bloch function. The Q2D electron energy is given by

$$e(\vec{k}, \sigma) = \hbar^2 \vec{k}^2/2m^* + E_0 - \sigma \frac{JN_0}{2} \langle S^\sigma \rangle - \sigma \left( \frac{1}{2} \right) g\mu_B B,$$

where $E_0 = \hbar^2 \pi^2/2m^*L^2$, $\sigma = \pm 1, +(-)$ for spin up(down), $m^*$ is the electron effective mass which is assumed to be constant along the growth direction, $x$ is the molar fraction of Mn ions and $\langle S^\sigma \rangle$ is the thermal average of the $z$-component of Mn spin taken over all Mn ions given by $\langle S^\sigma \rangle = -5/2B_{\sigma/2}(\zeta)$, $B_{\sigma/2}(\zeta)$ being the standard Brillouin function with $\zeta = g_{\text{Mn}}\mu_B B/k_BT$. The difference in energy between the states in Eqs. (3) and (4) is given by

$$e_i - e_f = e(\vec{k}) - e(\vec{k}') + g^* \mu_B B + g_{\text{Mn}}\mu_B B.$$

(7)

The energy term $g^* \mu_B B = g\mu_B B + JN_0\langle S^\sigma \rangle$, where $g^*$ is the effective $g$-factor of Q2DEG taking into account the exchange interaction between the conduction electron and Mn ions and $g_{\text{Mn}}$ is the $g$-factor of the Mn ion. Neglecting the Zeeman energy of the Mn spin in applied magnetic field, and calculating the spin matrix element
with the help of the relation $|\sin \pm 1/2| S^-|S/2, -1/2\rangle^2$, the transition probability for the Mn spin-flip can be obtained:

$$w = \frac{5\pi}{2h} J^2 x^2 |\phi(d)\rangle^4 |\langle k| |\delta(\bar{r})| \rangle_k^2 |\delta(\varepsilon_k - \varepsilon_{F}' + g^* \mu_B B). \quad (8)$$

We proceed now to specialize Eq. (8) to the case of high magnetic field. In order to see the effects on the Korringa relaxation rate of a strong quantizing magnetic field, we have to rewrite the matrix element in Eq. (2), to account for the strong field. The electronic states are now described by the quantum numbers $(j, k_z)$, and all $k_z$ states in the $j$-th Landau level are degenerate. As before, in Eq. (2), for simplicity, a state of the system is represented as the product of ket vectors for magnetic ion and electrons. We separate these products by semicolons and denote the initial and final states, respectively, as follows:

$$|i\rangle = \phi(d)|j, k_z\rangle,$$

$$|f\rangle = \phi(d)|j', k'_z\rangle. \quad (9)$$

The difference in energy between the initial and final states is given by

$$\varepsilon_i - \varepsilon_f = \hbar [\omega_c (j' - j) + \omega_Z],$$

where $\omega_c$ and $\omega_Z = g^* \mu_B B = g \mu_B B + J N_{0}(S)^2$ denote the cyclotron frequency and the effective Zeeman energy of 2D electrons, respectively. In terms of the matrix element $\langle \hat{H}_{e-d}|\rangle$, the transition probability is written as

$$w_{fj} = \frac{5\pi}{2h} J^2 x^2 |\phi(d)\rangle^4 |\langle j', k'_z| |\delta(\bar{r})| \rangle_j k_z |^2 |\delta(\hbar [\omega_c (j' - j) + \omega_Z]). \quad (11)$$

The total transition probability for Mn-spin-flip transition is obtained summing over all occupied initial Landau and all vacant final Landau states. One gets:

$$W = \frac{5\pi}{2h} J^2 x^2 |\phi(d)\rangle^4 \sum_{j, j'} \sum_{k_z} |\langle j, k_z| |\delta(\bar{r})| \rangle_j k_z |^2 |\delta(\hbar [\omega_c (j' - j) + \omega_Z]). \quad (12)$$

The expression for $\langle j, k'_z| |\delta(\bar{r})| \rangle_j k_z |$ is evaluated at $\bar{r} = 0$, because $H_{e-d}$ is independent of the position of the magnetic impurity, and represent a complex matrix element for the electron. The sum over the quantum number $k_z$ is simplified by assuming the relation $\sum_z |\langle j, k'_z| |\delta(\bar{r})| \rangle_j k_z |^2 = 1/(2 \pi a_B^2)$, where

$$a_B = (\hbar c/eB)^{1/2}. \quad (13)$$

At finite temperatures, the relaxation rate is computed by including the occupation probability for the 2D electron states. The expression for the relaxation rate $1/T_1 = 2W$ is given by

$$\frac{1}{T_1} = \frac{2}{5\pi} \sum_{j, j'} f_{ij} (1 - f_{ij}) W, \quad (14)$$

where $f_{ij} = 1/(e^{e_F - e_F}/kT + 1)$ is the Fermi function, $e_F$ is the Fermi energy, and $e_F^* = \hbar \omega_c (j + 1/2) - \sigma/(1/2) \hbar \omega_Z + E_0$ where, as before, $\sigma = +$ for up spin and $\sigma = -$ for down spin. By performing the computation of the factor $f_{ij} (1 - f_{ij})$ in Eq. (14), the relaxation rate turns to be

$$\frac{1}{T_1} = \frac{20\pi}{h} \frac{(JL)^2 x^2 |\phi(d)|^4}{(2\pi a_B^2)^2 \sinh \beta \hbar \omega_Z} \times \sum_{j, j'} \frac{(e^{\beta \hbar \omega_Z} - 1)}{[\sinh [\beta \hbar \omega_c (j + 1/2) - \varepsilon_F] + \sinh \beta \hbar \omega_Z]} \times \frac{1}{1/T_1}. \quad (15)$$

The effect of scattering is characterized by a broadening of the Lorentzian smearing of the delta function characterizing the ideal energy-conserving. The effect of scattering on the electron distribution is neglected. This is a reasonable assumption if the peak width is smaller than $k_BT_1$. In Eq. (15), $\phi(d) = \sin(\pi z/L)$ and $\beta \equiv \hbar/k_BT$, $k_B$ being the Boltzmann constant. As before, the relaxation rate $1/T_1$ measured experimentally is averaged in the $z$ direction. Thus, using Eq. (15), and the fact that $\langle |\phi(d)|^4 \rangle = 3/8$, we obtain for the relaxation rate averaged over the QW profile

$$\langle \frac{1}{T_1} \rangle_z = \frac{15\pi}{2h} \frac{(JL)^2 x^2}{(2\pi a_B^2)^2 \sinh \beta \hbar \omega_Z} \times \sum_{j} \frac{(e^{\beta \hbar \omega_Z} - 1)}{[\sinh [\beta \hbar \omega_c (j + 1/2) - \varepsilon_F] + \sinh \beta \hbar \omega_Z]} \times \frac{1}{1/T_1}. \quad (16)$$

3. Discussion of results

Eq. (16) is the expression for the Mn ion relaxation rate near a 2D electron gas in the high magnetic field limit. It follows from Eq. (16) that at low temperatures ($k_BT \ll \hbar \omega_c$), the quantum Mn-spin relaxation rate in a high magnetic field is seen to vary exponentially with the inverse temperature(i.e.$T_1^{-1} \propto \exp[-\hbar \omega_Z/(k_BT)]$). At high temperatures ($\hbar \omega_Z \ll k_BT \ll \hbar \omega_c$), on the other hand the relaxation rate increases linearly with temperature (i.e.$T_1^{-1} \propto k_BT/\hbar \omega_Z$), resulting in the Korringa relation $1/T_1 \propto T$. In order to provide a plot of Eq. (16) we have rewritten it in a more adequate manner, namely

$$\frac{1}{T_1} = \frac{15}{8h} \frac{(JL)^2 x^2}{\pi(\beta^2 m^*)^2} T Y(y) \quad (17)$$
where \( Y(y) \) is the dimensionless function given by

\[
Y(y) = \sum_j \left( \frac{\sinh^2 g^* y}{g^*} \right)^{1/2} \frac{(e^{g^*/2} - 1)}{[y^2 + (g^*/2)^2]^2]}.
\]

Here, \( y \equiv \beta \hbar \omega_c \) and \( \lambda \equiv \beta \Gamma \) and in Eq.(16) the Fermi energy was taken as \( \epsilon_F(B) = (j + 1/2) \hbar \omega_c + E_0 - \hbar \omega_Z \).

In Fig. 1 we plot \( Y(y) \) as a function of \( y \), when summation over two Landau level is done. We have used the following parameters for a (Cd, Mn)Te heterostructure: \( T^* = 100 \, K \) \((K_B T^* = 8.6 \, meV)\), \( n^* = 0.096 \, m \), \( E_0 = 35 \, meV \), \( L = 10 \, nm \), \( \Gamma = 0.2 \, meV \), deduced from the mobility of our sample[17], and a \( g^- \) enhancement factor [18] \( g^* = 10 \) (the effective \( g \) factor of the carriers can be as high as 1000). From this plot we see that as the field is increased, for a fixed temperature, \( T^{-1} \) approaches either a local maximum or local minimum due the function \( Y(y) \). Physically, this is due to the fact that the Fermi energy \( \epsilon_F(B) = (j + 1/2) \hbar \omega_c + E_0 - \hbar \omega_Z \) crosses either a spin-split level or the center of a spin-split gap, respectively. In this way, \( T^{-1} \) oscillates with the field, but the amplitude of this oscillation decreases as the field is decreased. At low magnetic field values, the decrease in the amplitude of oscillation of \( T^{-1} \) is caused by the \( s-d \) exchange interaction between the electron in the conduction band and the magnetic ions (the Zeeman energy is \( \hbar \omega_Z = g^* \mu_B B = g \mu_B B + J N g(S) \) and it can be made large provide the effective \( g^* \) is made large). In high magnetic-fields, \( T^{-1} \) exhibits increased oscillation with increasing magnetic field. The oscillating behavior is attributed to the interplay between the discrete Landau levels and the Fermi surface. This type of field dependence is similar to SdH oscillations in the diagonal Hall resistivity \( \rho_{xx} \) [18,19].

In summary, we have determined the Korringa relaxation rate \( 1/T_1 \) of Mn ion near a 2DEG in (Cd,Mn)Te heterostructures in high magnetic fields. The result shows that \( 1/T_1 \) scales with \( L^{-2} \), \( L \) being the quantum well width and it can be enhanced as \( L \) is decreased. Due to the presence of Landau levels and the spin-split mobility gap, \( T^{-1} \) oscillates with the field resembling the oscillations of \( \rho_{xx} \). Our calculation offers a method of investigating the dynamics of a magnetic ion such as Mn in a 2DEG and provide new information on the exchange parameter as well as information about the 2D electrons themselves.

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