Magnetic–Field–Induced Exciton Tunneling in Shallow Quantum Wells

A. Getter and I. E. Perakis

Department of Physics and Astronomy, Vanderbilt University, Nashville, Tennessee 37235

Abstract

We study the effect of the magnetic field orientation on the electroabsorption spectra of excitons confined in extremely shallow quantum wells. When the applied electric field is parallel to the quantum well plane, we demonstrate that, for in–plane magnetic field orientation, the discrete confined exciton peak undergoes a transition into a continuum resonance. In contrast, for perpendicular magnetic fields, the exciton peak exhibits the usual Stark red–shift. We show that such a dramatic dependence on the magnetic field orientation originates from a resonant coupling between the confined and the bulk–like excitons. Such coupling is caused by the interplay between the quantum–well potential and a velocity–dependent two–body interaction between the exciton center–of–mass and relative motion degrees of freedom induced by the in–plane magnetic field. As a result, the exciton tunnels out of the quantum well as a whole without being ionized. We discuss possible experimental applications of this effect.

Pacs numbers: 71.35.Ji, 71.35.-y, 78.66.-w, 78.66.Fd
The response of semiconductor excitons and hydrogenic atoms to strong magnetic and electric fields has been the focus of much attention in condensed matter, atomic, and molecular physics. When the field–induced forces become comparable to the internal Coulomb interactions, the response of such electron–hole (e–h) systems becomes non–perturbative and provides valuable insight into many–body and confinement effects. In addition to confining the e–h motion (Landau quantization), a magnetic field induces a momentum–dependent interaction between the center–of–mass (CM) and relative (RM) e–h degrees of freedom \[1\]. Some experimental manifestations of such two–body correlations were discussed in Ref. \[2\] for atomic ions, and in Ref. \[3\] for shallow quantum well excitons. In bulk semiconductors, the effect of an electric field is to ionize the RM Coulomb–bound state and thus to broaden the band–edge absorption spectrum (Franz–Keldysh effect) \[4\]. In quantum wells (QW’s), the confining potential inhibits such an ionization for electric fields perpendicular to the QW plane, which results in an exciton red–shift (Quantum Confined Stark Effect) \[5\]. The latter effect has become the basis of self–electrooptic–effect switching devices \[6\]. For efficient device operation, it is important to identify physical systems where an electric field causes large changes in the absorption spectrum (large contrast ratio). In addition, in order to avoid a carrier build–up within the QW region that would result into exciton broadening, it is important that the optically–excited e–h pairs rapidly escape from the QW. In shallow QW’s with depth comparable to the bulk exciton binding energy, it was shown that the combination of strong room–temperature excitons and very short carrier escape times \[7\] improves the device switching speed \[8\]. Such considerations spurred recent efforts to study the special properties of shallow QW excitons in the crossover between 3D– and 2D–like behavior. This confinement regime can be realized both in III–V \[9–12\] and in II–VI \[13\] semiconductor QW’s.

In this Letter, we study the electroabsorption spectrum of magneto–excitons in extremely shallow QW’s, whose depth is smaller than the Coulomb binding energy. Unlike in typical electroabsorption experiments \[3\]–\[5\], we consider weak electric fields \(E \parallel \) to the QW plane and compare the exciton absorption lineshape for magnetic fields parallel \((B_\parallel)\)
or perpendicular ($B_\perp$) to the QW plane. We show that, with the same in–plane electric field, the discrete QW exciton peak undergoes a dramatic transformation if we turn on $B_\parallel$ perpendicular to the electric field and gives way to a continuum resonance with diminished absorption strength. Such a strong dependence on the magnetic field orientation should be observable in electroabsorption experiments for in–plane applied electric fields. We attribute this transition to a strong two–body interaction between the CM and RM degrees of freedom induced by $B_\parallel$. We show that, with $E \neq 0$, this leads to a resonant tunneling of the exciton as a whole out of the QW, without ionization of the e–h Coulomb–bound state. We also discuss some possible applications of this effect.

We start with the exciton Hamiltonian. We choose the z–axis along the QW growth direction and the x–axis parallel to the in–plane electric field $E$. The magnetic field is always chosen perpendicular to the electric field, pointing either along the z–axis ($B_\perp$) or along the y–axis ($B_\parallel$). We work in the Landau gauge and denote by $(R, P)$ and $(r = r_e - r_h, p)$ the CM and RM position and momentum operators respectively. A unitary transformation of the Hamiltonian defined by the operator $U = \exp(-ie r \times B/2\hbar c)$ allows a partial separation of the CM and RM degrees of freedom [1]:

$$H = H_0 + V_e \left( Z + \frac{m_h}{M} z \right) + V_h \left( Z - \frac{m_e}{M} z \right),$$  \hspace{1cm} (1)

where $M = m_e + m_h$ is the total exciton mass, $m_e$ and $m_h$ are the electron and hole masses, $V_e$ and $V_h$ are the electron and hole QW potentials respectively, and the Hamiltonian $H_0$ is

$$H_0 = \frac{P^2}{2M} + H_{RM} + H_{int}. \hspace{1cm} (2)$$

Here, the Hamiltonian $H_{RM}$ describes a RM quasiparticle in the presence of one–body potentials due to the Coulomb interaction and the electric and magnetic fields, 

$$H_{RM} = \frac{D^2}{2\mu} - \frac{e^2}{\epsilon r} + \frac{e^2}{8\mu c^2} (B \times r)^2 + \frac{e}{2\mu c} \frac{m_h - m_e}{m_e + m_h} B \cdot L + e E \cdot r, \hspace{1cm} (3)$$

where $\mu = m_e m_h / M$ is the reduced e–h mass, $\epsilon$ the dielectric constant, and $L$ the angular momentum operator, while
describes a two–body interaction between the CM and RM degrees of freedom. Since we are concerned with extremely shallow QW’s, we assume the bulk exciton band structure and consider a single valence band with \( m_h = 0.15 m_0 \), \( m_0 \) being the free electron mass, and a conduction band with mass \( m_e = 0.067 m_0 \).

In the absence of the QW potential, the CM and RM subsystems are separable and the two–body exciton problem can be reduced to two one–body problems. Since in this case the CM and RM motions are uncorrelated, the exciton wave function is a product of a CM and a RM quasiparticle contribution. One might therefore expect weak CM–RM correlations in extremely shallow QW’s (with depth smaller than the Coulomb binding energy), in which case a factorized (adiabatic) exciton wave function \[3,13]\ would provide a good approximation. However, as discussed in Ref. \[3]\, such an adiabatic approximation fails in the absence of translational invariance due to a strong CM–RM correlation induced by \( H_{\text{int}} \). This is due to the existence (for \( H_{\text{int}} \neq 0 \)) of small–energy excitations between the ground states of \( H_{\text{RM}} + H_{\text{int}} \) corresponding to different CM–momentum values. In particular, in the absence of translational invariance, \( P_Z \) is no longer a constant of motion, so that the CM motion can excite the above low–energy RM degrees of freedom, which in turn affect strongly the CM motion. For \( E=0 \), such CM–RM correlation effects on the discrete exciton ground state can be described by using the general variational wave function of Ref. \[3]\. However, for \( E, B \parallel \neq 0 \), we find that such a calculation does not converge, indicating that the QW–confined exciton state is no longer the ground state of the system.

We proceed by expanding the two–body exciton wave function \( \Phi(Z,r) \) in the basis of eigenstates of the bulk Hamiltonian \( H_0 \) for different values of \( P_Z \):

\[
\Phi(Z,r) = \sum_{P_Z} \Psi(P_Z) e^{iP_Z Z / \hbar} \phi_{P_Z}(r),
\]

where \( \phi_{P_Z}(r) \) is the ground state of \( H_0 \) for a given value of \( P_Z \). The exciton wave function Eq. (3) takes into account excitations between the low–lying RM ground states of \( H_0 \) corresponding to different values of \( P_Z \). These are induced by the CM motion and are neglected...
in the adiabatic approximation. It is assumed that, due to the fact that the finite exciton binding energy is larger than the QW depth, the mixing of the excited magneto–exciton states for a given \( P_Z \) is weak. For \( E = 0 \) or \( B_\parallel = 0 \), the above wave function compares very well with the variational solution of Ref. [3] in extremely shallow QW’s.

Using the wave function Eq. (5), the Schrödinger equation with the total Hamiltonian \( H \) leads to the following equation for the CM–momentum wave function \( \Psi(P_Z) \):

\[
[\varepsilon(P_Z) - \varepsilon] \Psi(P_Z) = -\sum_{P'_Z} V_{\text{eff}}(P_Z, P'_Z) \Psi(P'_Z)
\]

(6)

where \( \varepsilon \) is the exciton energy, \( \varepsilon(P_Z) \) is the ground state energy of the Hamiltonian \( H_0 \) for a given value of \( P_Z \), and

\[
V_{\text{eff}}(P_Z, P'_Z) = \int dZ e^{i(P'_Z - P_Z)Z/\hbar} \langle \phi_{P_Z} | V_e \left( Z + \frac{m_e}{M} z \right) + V_h \left( Z - \frac{m_e}{M} z \right) | \phi_{P'_Z} \rangle
\]

(7)

is an effective non–local QW potential that depends on both the RM wave functions and the CM momentum. To calculate \( \varepsilon(P_Z) \) and \( V_{\text{eff}} \), we diagonalized the RM Hamiltonian \( H_0 \) for different values of \( P_Z \) using a real–space Gaussian basis set similar to Ref. [3] and including basis elements with all the prefactors allowed by the reduced symmetry when \( E \neq 0 \). Although for \( B_\parallel = 0 \) Eq. (3) has a discrete ground state, with finite \( B_\parallel \) and \( E \) we found many eigenstates closely–spaced in energy, suggesting that the breakdown of adiabaticity changes the qualitative behavior of the CM motion.

With the above–obtained exciton wave function \( \Phi(Z, r) \), we calculated the absorption spectrum using Fermi’s golden rule. Our results in the frequency range of the confined exciton are presented in Fig. 1(a), for magnetic field perpendicular to the QW plane \( (B_\perp) \), and in Fig. 1(b), for magnetic field parallel to the QW plane \( (B_\parallel) \). Both magnetic field orientations are perpendicular to that of the in–plane electric field. Despite the fact that the QW depths \( V_e = 1.2 \text{meV} \) and \( V_h = 0.8 \text{meV} \) are much smaller than the bulk exciton binding energy, \( E_B \sim 4 \text{meV} \), our calculations show a sharp contrast for the two different magnetic field orientations between the changes of the magneto–exciton lineshape induced by the electric field. With \( B_\parallel = 0 \), we obtain a small Stark red–shift of the exciton peak [see Fig. 5]
With finite $B\parallel$ however, the same weak in–plane electric field induces a large exciton broadening and decrease in absorption strength [see Fig. 1(b)]. It is important to note here that such an effect is not due to the ionization of the RM exciton state. In fact, since both $B\parallel$ and $B\perp$ are perpendicular to the electric field pointing along the x–axis, the same diamagnetic potential, quadratic in $x$, inhibits the RM ionization by opposing the electric field potential, which is linear in $x$. Unlike in the $B=0$ case, the electric field reduces the Coulomb binding energy without ionizing the magneto–exciton.

To interpret the above dramatic difference in the electroabsorption spectrum, let us compare the exciton Hamiltonians for the two different magnetic field orientations. The differences in $H_{\text{RM}}$ cannot explain our effect, which is caused by the two–body interaction $H_{\text{int}}$. In the bulk semiconductor, the CM momentum $P$ is a constant of motion. For optically–active excitons, $P = 0$ in the dipole approximation and therefore $H_{\text{int}} = 0$ does not affect the linear optical spectra. In a QW however, $P_Z$ becomes a dynamical variable due to the breakdown of the translational invariance. For $B\parallel = 0$, $H_{\text{int}}$ only depends on the in–plane components of the CM momentum, which vanish for optically–active excitons. Therefore, in this case, $H_{\text{int}}$ does not affect the absorption spectrum. On the other hand, with finite $B\parallel$, one has $H_{\text{int}} = -eB\parallel xP_Z/cM$. This two–body Hamiltonian may be thought of as a fluctuating electric field potential, which intimately couples the CM and RM degrees of freedom. By changing the magnetic field orientation, we therefore tune the strength of $H_{\text{int}}$, which affects the exciton CM motion via the renormalization of $\varepsilon(P_Z)$ and $V_{\text{eff}}$ [see Eq. (1)].

Let us first consider the dispersion relation $\varepsilon(P_Z)$ of the CM degree of freedom. With magnetic field perpendicular to the QW plane, the CM–RM interaction $H_{\text{int}}$ vanishes and the momentum–dependence of $\varepsilon(P_Z)$ is quadratic. This is no longer the case when both $B\parallel$ and $E$ are finite. In Fig. 2(a) we show the dispersion relation for in–plane magnetic field $B\parallel$ and different values of the electric field. Our numerical calculation is consistent with the analytic asymptotic expressions derived in Ref. [1] for a bulk semiconductor and high magnetic fields. The total effective electric field $\mathcal{E}(P_Z) = E - B\parallel P_Z/cM$ acting on the
RM depends on the CM momentum, which in QW’s is no longer a constant of motion. For values of $P_Z$ corresponding to large $\mathcal{E}(P_Z)$, such an electric field dominates over the Coulomb interaction and leads to a linear dependence of $\varepsilon(P_Z)$ on the CM momentum. On the other hand, for values of $P_Z$ corresponding to small $\mathcal{E}(P_Z)$, the Coulomb interaction leads to a local minimum in the CM dispersion relation. The most important feature of the spectrum of Fig. 2(a) is the Coulomb–induced degeneracy between high momentum eigenstates and those corresponding to $P_Z$ close to the dispersion minimum.

Let us now turn to the effective QW potential $V_{\text{eff}}$. As can be seen from Eq. (6), $V_{\text{eff}}$ causes an appreciable mixing of all the low–energy ground states of the RM Hamiltonian $H_{R\text{M}} + H_{\text{int}}$ corresponding to the different values of $P_Z$. Such an excitation of the RM is absent within the adiabatic approximation and changes the qualitative behavior of the CM momentum wave function $\Psi(P_Z)$. Indeed, for finite $B_\parallel$ and $E$, we find many exciton eigenstates with energies closely–spaced around the energy of the confined exciton. The probability densities for some of these optically–active states are shown in Fig. 2(b). As can be seen, the CM state is a superposition of two states, one with wave function with finite momentum distribution centered at small positive $P_Z$, and another with wave function sharply peaked at large $P_Z$. The first peak in $\Psi(P_Z)$ corresponds to the QW–confined exciton state, whose energy lies below the local minimum of $\varepsilon(P_Z)$, while the second sharp peak comes from the coupling of bulk excitons whose energies are degenerate with that of the discrete confined state. The above states are resonantly coupled by $V_{\text{eff}}$, meaning that the confined exciton can tunnel into the continuum of bulk excitons. Such an effect for $B_\parallel \neq 0$ manifests itself via the transformation of the discrete QW exciton into a continuum resonance [see Fig. 1(b)]. On the other hand, for $B_\parallel = 0$, there is no degeneracy and we obtain a sharp QW–confined exciton peak [see Fig. 1(a)].

Let us now discuss some possible applications of our results. First, a magnetic field of a few Tesla allows one to strongly modify the exciton absorption by using very weak electric fields. For example, as shown in Fig. 1(b), an electric field of only $E \sim 3 \text{kV/cm}$ leads to a $\sim 80\%$ decrease in the exciton strength. Such a sharp contrast ratio is desirable for the
efficient operation of switching devices [8,6]. Unlike in typical electroabsorption experiments [8,6,5], sharp changes in the absorption spectrum are achieved by using a weak in–plane electric field that does not ionize the RM exciton state. This suppresses the undesirable Franz–Keldysh broadening that limits the contrast ratio. In our case, the confined exciton as a whole tunnels out of the QW and into degenerate dark bulk–like exciton states with high CM momentum. This as well as the fact that the electric field shifts the minimum of $\varepsilon(P_Z)$ to nonzero CM momentum values lead to a very strong decrease in the exciton absorption strength.

In conclusion, we showed that, by changing the magnetic field orientation from perpendicular to parallel to the plane of an extremely shallow quantum well, the discrete confined exciton state transforms into a continuum resonance in the presence of a weak in–plane electric field. This transition is due to a resonant tunneling of the confined exciton as a whole out of the QW region, without ionization of the RM bound state. This is caused by the interplay between the quantum well potential and a two–body velocity–dependent interaction between the CM and RM degrees of freedom induced with an in–plane magnetic field. The above dramatic effect can be observed with electroabsorption experiments that apply in–plane electric fields to extremely shallow GaAs/AlGaAs QW excitons. The role of shallow QW potentials and the above interactions on the transition from regularity to chaos [2] and the dynamical exciton–exciton interactions in ultrafast four–wave–mixing spectroscopy [14] remain to be studied.

This work was supported by the NSF CAREER award ECS-9703453, and, in part, by ONR Grant N00014-96-1-1042 and by Hitachi Ltd. Part of this work was performed while A.G. was at Rutgers University. We thank W. Knox, K. W. Goossen, J. E. Cunningham, T. V. Shahbazyan, and especially S. A. Jackson for valuable discussions.
REFERENCES

[1] L. P. Gorkov and I. E. Dzyaloshinskii, Sov. Phys. JETP 26, 449 (1968).

[2] P. Schmelcher and L. S. Cederbaum, Phys. Rev. Lett. 74, 662 (1995).

[3] M. Fritze et. al. Phys. Rev. Lett. 76, 106 (1996); A. Getter et. al., Solid State Comm. 98, 379 (1996); A. Getter, PhD Thesis, Rutgers University (1998).

[4] D. Dow and D. Redfield, Phys. Rev. B 1, 3358 (1970).

[5] D. A. B Miller et. al., Phys. Rev. Lett. 53, 2713 (1984); Phys. Rev. B 32, 1043 (1985).

[6] See e.g D. A. B. Miller, Opt. Quantum Electron. 22, 61 (1990).

[7] G. von Plessen et. al., Appl. Phys. Lett. 63, 2372 (1993); Phys. Rev. B 53, 13688 (1996).

[8] K. W. Goossen et. al., Phys. Rev. B 45, 13773 (1992); Appl. Phys. Lett. 57, 2582 (1990).

[9] J. Tignon et. al., Phys. Rev. B 58, 7076 (1998); Appl. Phys. Lett. 72, 1217 (1998).

[10] R. C. Iotti and L. C. Andreani, Phys. Rev. B 56, 3922 (1997).

[11] P. E. Simmonds et. al., Phys. Rev. B 50, 11251 (1994).

[12] I. Brener et. al., Phys. Rev. Lett. 70, 319 (1993).

[13] See e.g., J. Kossut et. al., Phys. Rev. B 56, 9775 (1997); J. Warnock et. al., Phys. Rev. B 48, 17321 (1993); A. Alexandrou et. al., Phys. Rev. B 50, 2727 (1994); N. Dai et. al., Phys. Rev. B 50, 18153 (1994).

[14] P. Kner et. al., Phys. Rev. Lett. 78, 1319 (1997).
FIGURES

FIG. 1. Change in the QW–confined magneto–exciton absorption lineshape induced by an in–plane electric field of $E=2.9$ kV/cm for (a) perpendicular magnetic field $B_\perp=10$T, and (b) in–plane magnetic field $B_\parallel=10$T.

FIG. 2. (a) CM dispersion relation $\varepsilon(P_Z)$ for $B_\parallel=10$ T and $E=0$ (solid curve), 1.16KV/cm (dashed curve), 1.74 KV/cm (long–dashed curve), 2.32 KV/cm (dotted–dashed curve), and 2.9 KV/cm (dotted curve). (b) Momentum probability density $|\Psi(P_Z)|^2$ of several low–lying QW exciton states for $B_\parallel=10$ T and $E=2.9$KV/cm
Fig. 1
Absorption Spectrum (arb. units)

Energy (Rydbergs)

E=2.9 kV/cm
E=0

(b)

Fig. 1
Fig. 2
Fig. 2