Magnetoexciton escape from shallow quantum wells induced by in-plane electric fields

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We study the change in the magnetoexciton absorption induced by an electric field parallel to the plane of a shallow quantum well. 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 redshift. We show that such a transition originates from a resonant coupling between the discrete confined exciton state and the continuum of dark bulk exciton states. As a result, the confined exciton tunnels out of the quantum well *as a whole* without being ionized. We discuss the possible experimental applications of this effect. [S0163-1829(99)03547-X]

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 nonperturbative and provides valuable insight into many-body and confinement effects. An external magnetic field confines the e-h motion (Landau quantization) and, in addition, induces a momentum-dependent interaction between the center-of-mass (CM) and relative motion (RM) e-h degrees of freedom.[1–3] 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 redshift (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 buildup within the QW region that would result in 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 (Refs. 9–12 and 3) and in II-VI (Ref. 13) semiconductor QW's.

In this paper, we study the change in the magnetoexciton absorption line shape induced by a weak electric field *E parallel* to the plane of an extremely shallow QW and perpendicular to the magnetic field. In particular, we compare the exciton absorption 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 $B_\parallel$ and gives way to a *continuum resonance* with diminished absorption strength. This is in sharp contrast to the situation considered in Ref. 3, where it was shown that in the absence of an in-plane electric field, $B_\parallel$ *increases the confinement* of the *discrete* QW exciton. We attribute the above transition to the tunneling of the confined exciton *as a whole* out of the quantum well, *without ionization of the e-h Coulomb bound state*. We show that this effect is induced by the interplay between the in-plane electric field potential, the two-body interaction between the CM and RM degrees of freedom, and the QW potential. As a result, the discrete confined exciton state becomes resonant with the continuum of dark bulk exciton states, to which it is coupled by the shallow QW potential.

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 $\mathbf{E}$. 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 $(\mathbf{R},\mathbf{P})$ and $(\mathbf{r}=\mathbf{r}_e-\mathbf{r}_h,\mathbf{p})$ the CM and RM position and momentum operators, respectively. A unitary transformation of the Hamiltonian, defined by the operator $U=\exp[-ie(\mathbf{r}\times\mathbf{B})\cdot\mathbf{R}/2\hbar c]$, allows a partial separation of the CM and RM degrees of freedom:[1]

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

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

$$H_0=\frac{P_Z^2}{2M}+H_{\rm RM}+H_{\rm int}. \tag{2}$$

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

$$H_{\rm RM}=\frac{\mathbf{p}^2}{2\mu}-\frac{e^2}{\epsilon r}+\frac{e^2}{8\mu c^2}(\mathbf{B}\times\mathbf{r})^2+\frac{e}{2\mu c}\frac{m_h-m_e}{m_e+m_h}\mathbf{B}\cdot\mathbf{L}+e\mathbf{E}\cdot\mathbf{r}, \tag{3}$$

where $\mu=m_em_h/M$ is the reduced e-h mass, ϵ the dielectric constant, and $\mathbf{L}$ the angular momentum operator. The term

$$H_{\rm int}=\frac{e}{cM}(\mathbf{P}\times\mathbf{B})\cdot\mathbf{r} \tag{4}$$

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 $m_h=0.15m_0$, m_0 being the free electron mass, and a conduction band with mass $m_e=0.067m_0$. Below we only consider zero in-plane components of the CM momentum since we are dealing with optically active excitons.

In the absence of the QW, the CM and RM degrees of freedom are separable and the two-body exciton problem can be reduced to two one-body problems. As a result, in this case the exciton wave function can be expressed as a product of a CM and a RM quasiparticle contribution. One might therefore expect weak CM-RM correlations in extremely shallow QW's (with a depth smaller than the Coulomb binding energy), in which case a factorized (adiabatic) exciton wave function[3,13] would provide a good approximation. However, such an adiabatic approximation fails in the absence of translational invariance due to a strong CM-RM correlation induced by $H_{\rm int}$. This is due to the existence (for $H_{\rm int}\neq 0$) of small-energy excitations between the different exciton ground states of $H_{\rm RM}+H_{\rm int}$ 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 affects strongly the CM motion. For $E=0$, the ground state is a *discrete* exciton bound state, and the CM-RM correlations can be described by using the general variational wave function of Ref. 3. However, as we demonstrate below, in the presence of an in-plane electric field, $E\neq 0$, the QW-confined exciton is no longer the ground state of the system.

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

$$\Phi(Z,\mathbf{r})=\sum_{P_Z}\Psi(P_Z)e^{iP_ZZ/\hbar}\phi_{P_Z}(\mathbf{r}), \tag{5}$$

where $\phi_{P_Z}(\mathbf{r})$ is the ground state of H_0 for a given value of P_Z. The exciton wave function, Eq. (5), takes into account the excitations between the low-lying RM ground states of H_0 corresponding to different values of P_Z. Such an excitation of the RM by the CM motion is neglected in the adiabatic approximation. For $E=0$, the above wave function compares very well with the variational wave function[3] in extremely shallow QW's. This implies that the mixing of excited RM states for a given P_Z is weak when the exciton binding energy exceeds the QW depth.

Using the wave function, Eq. (5), the Schrödinger equation with the Hamiltonian H in Eq. (1) leads to the following equation for the CM-momentum wave function $\Psi(P_Z)$:

$$[\varepsilon(P_Z)-\varepsilon]\Psi(P_Z)=-\sum_{P'_Z}V_{\rm eff}(P_Z,P'_Z)\Psi(P'_Z), \tag{6}$$

where ε 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

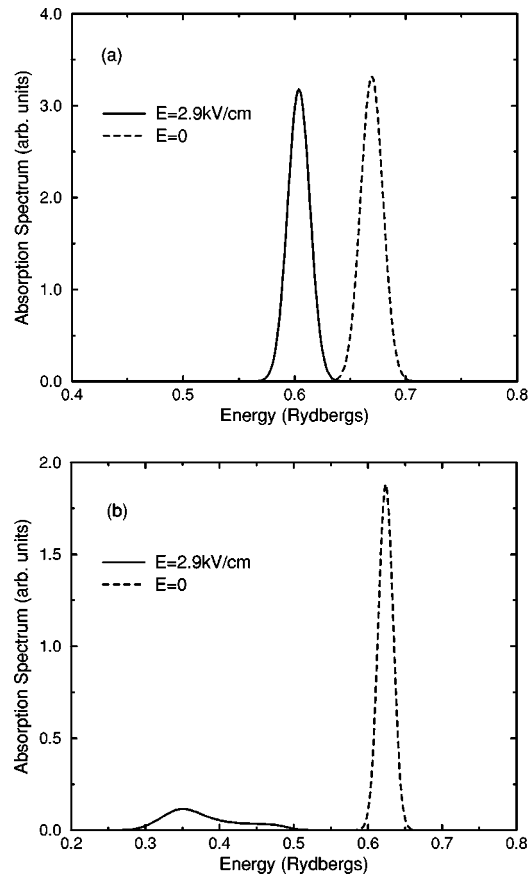

FIG. 1. Change in the QW-confined magnetoexciton absorption line shape induced by an in-plane electric field of $E=2.9$ kV/cm for (a) perpendicular magnetic field $B_\perp=10$ T, and (b) in-plane magnetic field $B_\parallel=10$ T. The value of the Rydberg is $E_B=4$ meV.

$$V_{\rm eff}(P_Z,P'_Z)=\int dZ e^{i(P'_Z-P_Z)Z/\hbar}\langle\phi_{P_Z}|V_e\left(Z+\frac{m_h}{M}z\right)+V_h\left(Z-\frac{m_e}{M}z\right)|\phi_{P'_Z}\rangle \tag{7}$$

is an effective nonlocal QW potential that depends on both the RM wave functions and the CM momentum. To calculate $\varepsilon(P_Z)$ and $V_{\rm eff}$, we diagonalized the RM Hamiltonian H_0 for different values of P_Z using a real-space Gaussian basis set[3] and including basis elements with all the prefactors allowed by the reduced symmetry for $E\neq 0$. Even though with $B_\perp$ or for $E=0$ Eq. (6) does have a discrete QW exciton ground state, with finite $B_\parallel$ *and* E we found many exciton eigenstates very closely spaced in energy, which suggests a transition in the energy spectrum.

With the above exciton wave function $\Phi(Z,\mathbf{r})$, 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 a magnetic field perpendicular to the QW plane ($B_\perp$), and in Fig. 1(b) for a magnetic field parallel to the QW plane ($B_\parallel$). Both magnetic field orientations are perpendicular to the in-plane electric field. We considered a QW with width $L=100$ Å and depths $V_e=1.2$ meV and $V_h=0.8$ meV, which correspond to the typical Al content values $x\simeq 0.2\%$ in extremely shallow

$GaAs/Al_xGa_{1-x}As$ QW's.[3] Despite the fact that the QW depths are much smaller than the bulk exciton binding energy, $E_B \sim 4$ meV, we find a sharp contrast for the two different magnetic field orientations between the changes of the exciton line shape induced by the *same* in-plane electric field. With $B_\parallel = 0$, we obtain a small Stark redshift of the exciton peak [see Fig. 1(a)]. 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 the above 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, which is *quadratic* in x, inhibits the RM ionization by opposing the electric field potential (which is *linear* in x). In other words, the in-plane electric field reduces the Coulomb binding energy without ionizing the magnetoexciton, unlike in the $B = 0$ case.

To interpret the above transition in the electroabsorption spectrum, let us compare the exciton Hamiltonians for the two different magnetic field orientations. Obviously, the differences in H_{RM} cannot explain our effect, which is caused by the interplay between the in-plane electric field potential, the two-body interaction H_{int}, and the QW potential. For $B_\parallel = 0$, H_{int} only depends on the in-plane components of the CM momentum, which vanish in the dipole approximation for optically active excitons. Therefore, in this case, H_{int} does not affect the absorption spectrum. On the other hand, with finite $B_\parallel$, one has $H_{\mathrm{int}} = -eB_\parallel x P_Z / cM$. In a QW, P_Z becomes a dynamical variable due to the breakdown of the translational invariance. The two-body Hamiltonian H_{int} may then be thought of as a *fluctuating* electric field potential. By changing the magnetic field orientation, we therefore tune the strength of H_{int}, which affects the exciton CM motion via the renormalization of $\varepsilon(P_Z)$ and V_{eff} [see Eq. (6)]. For $E = 0$, this would increase the exciton confinement.[3] Below we demonstrate that an in-plane electric field leads to quite the opposite effect by *changing the low-energy spectrum* of the Hamiltonian H *without* ionizing the RM Coulomb bound state.

Let us first consider the dispersion relation $\varepsilon(P_Z)$ of the CM degree of freedom. With the magnetic field perpendicular to the QW plane, the CM-RM interaction H_{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. This can be seen in Fig. 2(a), which shows the dispersion relation for the in-plane magnetic field $B_\parallel$ and different values of the external electric field E. Such momentum dependence is consistent with the analytic asymptotic expressions derived in Ref. 1 for high magnetic fields and can be understood as follows. 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 is no longer a constant of motion in QW's. For values of P_Z corresponding to large $\mathcal{E}(P_Z)$, such an electric field leads to large e-h separations and thus dominates over the Coulomb interaction and, for $E \neq 0$, leads to a linear dependence of $\varepsilon(P_Z)$ on the CM momentum.[1] On the other hand, for CM momenta such that $\mathcal{E}(P_Z)$ is small, the Coulomb interaction leads to a *local minimum* in the CM dispersion relation. The most important feature of the spectrum of

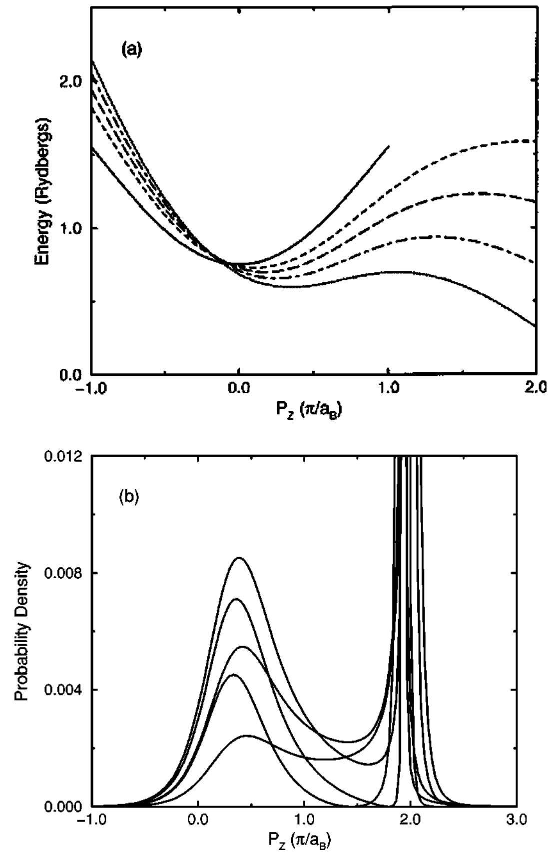

FIG. 2. (a) CM dispersion relation $\varepsilon(P_Z)$ versus P_Z (in units of the inverse Bohr radius a_B) for $B_\parallel = 10$ T and $E = 0$ (solid curve), 1.16 kV/cm (dashed curve), 1.74 kV/cm (long-dashed curve), 2.32 kV/cm (dotted-dashed curve), and 2.9 kV/cm (dotted curve). (b) Momentum probability density $|\Psi(P_Z)|^2$ of several low-lying QW exciton states for $B_\parallel = 10$ T and $E = 2.9$ kV/cm.

Fig. 2(a) is the Coulomb-induced *degeneracy* between high momentum exciton eigenstates and those corresponding to P_Z close to the dispersion minimum. Such a degeneracy is *absent* for $E = 0$, in which case $B_\parallel$ simply increases the effective mass.[3]

Let us now turn to the effective QW potential V_{eff}. As can be seen from Eq. (6), V_{eff} mixes the low-energy ground states of the RM Hamiltonian, $H_{\mathrm{RM}} + H_{\mathrm{int}}$, corresponding to the different values of P_Z. Such an excitation of the RM, absent within the adiabatic approximation, drastically changes the low-energy spectrum of the Hamiltonian H when *both* $B_\parallel$ and E are finite. Indeed, we find that the discrete confined exciton ground state gives way to many exciton eigenstates with very closely spaced energies. With increasing basis size, this set of eigenstates becomes continuous. The momentum 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 whose wave function has a finite momentum distribution centered at small positive P_Z, and another with a wave function sharply peaked at large P_Z. The first peak in $\Psi(P_Z)$ comes from the discrete QW-confined exciton state, whose energy lies below the local minimum of $\varepsilon(P_Z)$, while the second sharp peak comes from its coupling to the continuum of the bulk exciton states with high CM momenta, whose energies are brought into resonance by the electric field [see

Fig. 2(a)]. The above states are resonantly coupled via V_{eff}, meaning that the confined exciton can tunnel into the continuum of dark bulk excitons. Such an effect manifests itself in the transformation of the discrete QW exciton into a continuum resonance [see Fig. 1(b)]. On the other hand, for $B_{\parallel}=0$ or for $E=0$, 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 teslas 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 $E\sim3$ kV/cm leads to a $\sim80\%$ 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] this is 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 the degenerate dark bulklike 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 leads to a very strong decrease in the exciton absorption strength.

In conclusion, we showed that, if we change the magnetic field orientation from perpendicular to parallel to the plane of an extremely shallow quantum well, a weak in-plane electric field transforms the discrete confined exciton state into a continuum resonance. This transition is due to the resonant tunneling of the confined exciton as a whole out of the QW region, without ionization of the RM bound state, and is caused by an interplay among the in-plane electric field potential, the two-body interaction between the CM and RM degrees of freedom, and the QW potential. This effect can be observed with electroabsorption experiments in extremely shallow GaAs/$Al_xGa_{1-x}As$ QW's.

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