

Far-infrared $s \rightarrow p^\pm$ interexciton transitions in InGaAs/GaAs coupled double quantum wells

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We consider $s \rightarrow p^\pm$ interexciton far-infrared (FIR) magneto-optical transitions in coupled double quantum wells (DQWs). Spatially direct (intrawell) and indirect (interwell) excitons in strained $\text{In}_x\text{Ga}_{1-x}\text{As}/\text{GaAs}$ symmetric DQWs with a simple valence band are considered. The evolution of the transition energies and oscillator strengths as functions of the transverse magnetic field B in different regimes is studied: we consider the direct regime (zero and low transverse electric field \mathcal{E}), the indirect regime (high \mathcal{E}), and the indirect–direct crossover (induced by increasing B at intermediate \mathcal{E}). © 1996 American Institute of Physics. [S0021-3640(96)01515-0]

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Currently, there is considerable experimental interest^{1,2} in magnetoexcitons (MXs) in coupled double quantum wells (DQWs). In these systems, under the application of an external transverse electric field \mathcal{E} , the indirect (interwell) exciton, which has large radiative lifetimes, becomes the ground state. This opens up a possibility for observing collective effects in neutral $e-h$ systems in high B at low excitonic temperatures. Spectroscopic identification of collective effects requires a detailed knowledge of the magneto-optics of excitons. Our theory³ of the interband magneto-optics of MXs in DQWs is in good agreement with experiment.¹ Intraband FIR magnetospectroscopy can be very effective in resolving the fine structure of the ground and excited states of quasi-2D excitons in DQWs. In this paper, we study theoretically the FIR magneto-optical transitions of such excitons in fields $B > 1$ T.

We consider symmetric strained $\text{In}_x\text{Ga}_{1-x}\text{As}/\text{GaAs}$ DQWs. Due to the strain, the light-hole band is split off in such structures. As a result, the features present in the excitonic spectra in a DQW are entirely due to coupling through the barrier and are not complicated by heavy- and light-hole anticrossings. Here we present the results for the particular values of the well and barrier widths $L_1 = L_2 = L_b = 60$ Å and $x = 0.2$. For such DQWs the intermediate-barrier regime is realized at zero and low B .¹³ With increasing B , it gradually changes to a wide-barrier regime. This is because the difference in the Coulomb binding energies of the direct (D) and indirect (I) excitons, $\delta E_{DI} = E_D - E_I$, increases with B and becomes larger than the single-particle e and h symmetric–antisymmetric splittings Δ_e, Δ_h . For each, s, p^\pm, \dots state, there are four excitonic terms in

a DQW (instead of one for a single QW in the lowest electric subband).^{1,3-5} For a symmetric DQW at $\mathcal{E}=0$, each of the D and I states splits into states which are symmetric (S) and antisymmetric (A) under the simultaneous inversion $\{z_e \rightarrow -z_e, z_h \rightarrow -z_h\}$ (Refs. 3 and 4). In the wide-barrier regime, the excitonic symmetric-antisymmetric splittings $\Delta_X \approx \Delta_e \Delta_h / \delta E_{DI}$ are determined by the two-particle $e-h$ cotunneling through the barrier.³ Thus, Δ_X are suppressed by the excitonic effects ($\sim \delta E_{DI}^{-1}$), in particular, they decrease with increasing B .

In the Faraday geometry (the light propagates along \mathbf{B}), the interaction of an exciton with the ac electric field (with amplitude \mathcal{F}_0 and frequency ω) of circularly polarized FIR radiation is of the form

$$\delta \hat{V}^\pm = \frac{e \mathcal{F}_0}{\omega} \left(\frac{\pi_e^\pm}{m_e} - \frac{\pi_h^\pm}{m_h} \right) \exp(-i\omega t), \quad (1)$$

where “+” is for the left ($\Delta l_z = 1$) and “-” is for the right ($\Delta l_z = -1$) circularly polarized light; $\pi_j^\pm = \pi_{jx} \pm i \pi_{jy}$ ($\pi_j = -i \hbar \nabla_{\rho j} - (e_j/c) \mathbf{A}_j$, $j = e, h$) are the e and h inter-Landau raising (π_e^+, π_h^-) and lowering (π_e^-, π_h^+) operators, $\delta \hat{V}^\pm$ commutes with the operator of the magnetic⁶ center-of-mass momentum $\hat{\mathbf{K}}$ of the exciton in B : $[\delta \hat{V}^\pm, \hat{\mathbf{K}}] = 0$; thus \mathbf{K} is conserved. Clearly, FIR intraband spectroscopy probes all populated excitonic states having finite \mathbf{K} ; this is in contrast with interband transitions, for which only $\mathbf{K} = 0$ excitons are optically active. Here, assuming low temperatures, we consider FIR-active transitions between excitons with $\mathbf{K} = 0$.

To calculate the eigenstates, we use the expansion³ of the wave functions of MXs with angular momentum projection l_z (e.g., ≥ 0) in a DQW:

$$\Psi_{\mathbf{K}=0, l_z}(\mathbf{r}_e, \mathbf{r}_h) = \exp\left(\frac{i[\boldsymbol{\rho} \times \mathbf{R}]_z}{2l_B^2}\right) \sum_{i,j=1,2} \sum_n A_{ijn}(l_z) \zeta_i(z_e) \xi_j(z_h) \phi_{n+l_z n}(\boldsymbol{\rho}). \quad (2)$$

Here $l_B = (\hbar c / eB)^{1/2}$, $\zeta_i(z_e)$ are the electron and $\xi_j(z_h)$ the hole wave functions describing the free transverse motion in a DQW in field \mathcal{E} . The $\phi_{nm}(\boldsymbol{\rho})$ are factored wave functions in B ; $\boldsymbol{\rho} = \boldsymbol{\rho}_e - \boldsymbol{\rho}_h$ is the in-plane $e-h$ separation. For MXs the quantum numbers n and m correspond to the e and h Landau levels, respectively, and the angular momentum projection $l_z = n - m$. We include, depending on B , from ten to 36 Landau level orbitals ϕ_{nm} ; for the values of the effective masses used and other details see Ref. 3. In high magnetic fields, the $1s$ excitonic states are predominantly formed by the orbital ϕ_{00} corresponding to the zeroth e and h Landau levels. Due to the Coulomb $e-h$ interaction, there is also a small $\sim l_B / a_{Be(h)} \ll 1$ [$a_{Be(h)} = \epsilon \hbar^2 / m_{e(h)} e^2$] admixture of higher Landau levels ϕ_{nn} . Similarly, the $2p^+$ ($2p^-$) excitonic states are formed predominantly by the orbital ϕ_{10} (ϕ_{01}) with small admixtures of ϕ_{n+1n} (ϕ_{nn+1}) states. We label excitonic states by the high-field quantum numbers (D_{nm}, I_{nm}) ; at $\mathcal{E} = 0$ we also indicate the symmetry under inversion (S or A), and at high \mathcal{E} , the lower (D_{nm}^-, I_{nm}^-) or the upper (D_{nm}^+, I_{nm}^+) branches in the spectra.^{1,3}

For $\mathbf{K} = 0$ excitons in B the selection rules are

$$\langle \Psi'_{\mathbf{K}=0, l'_z} | \delta \hat{V}^\pm | \Psi_{\mathbf{K}=0, l_z} \rangle \sim \delta_{l'_z, l_z \pm 1}. \quad (3)$$

(Note that excitons with $\mathbf{K} \neq 0$ cannot be characterized in B by the projection l_z of the angular momentum of the relative $e-h$ motion. As a result, the selection rules are considerably relaxed. The effects associated at finite temperatures with FIR absorption by 2D MXs with $\mathbf{K} \neq 0$ will be discussed elsewhere.⁷ At high B the interexciton transition $1s \rightarrow 2p^+$ ($1s \rightarrow 2p^-$) can be thought of as the electron (hole) cyclotron resonance (CR) $\phi_{00} \rightarrow \phi_{10}$ ($\phi_{00} \rightarrow \phi_{01}$), modified by excitonic effects. The main excitonic effect is the renormalization of the transition energies. Indeed, the initial $1s$ state is more strongly bound than the final $2p^\pm$ states. Hence, the interexciton transition energies are *larger* than the bare CR energies $\hbar\omega_{ce(h)} = eB/m_{e(h)}$. In the strictly-2D case and the high-magnetic-field limit, the binding energies of the $1s$ and $2p^\pm$ MXs are⁸ $E_{00} = E_0$ and $E_{10} = E_{01} = \frac{1}{2}E_0$, respectively; here $E_0 = \sqrt{\pi/2}e^2/\epsilon l_B \sim \sqrt{B}$. Hence, the transition energies $1s \rightarrow 2p^\pm$ in this limit are

$$E_{1s \rightarrow 2p^+} = \hbar\omega_{ce} + \frac{1}{2}E_0, \quad E_{1s \rightarrow 2p^-} = \hbar\omega_{ch} + \frac{1}{2}E_0. \quad (4)$$

Excitonic FIR transitions are affected by the Coulomb interparticle $e-h$ interaction. Kohn's theorem⁹ is applicable because the charge-to-mass ratios are different for the electron and hole. It can be shown however that, at any B (for simple parabolic bands and excitons with $\mathbf{K} = 0$), the *difference* $E_{s \rightarrow p^+} - E_{s \rightarrow p^-} = \hbar\omega_{ce} - \hbar\omega_{ch}$ does not depend on the $e-h$ interaction (cf. with the theorem¹⁰ for one-compound electron systems in a field B).

The evolution with B (at $\mathcal{E} = 0$) of the energies and dipole matrix elements

$$|f|^2 \sim \left| \left\langle \Psi_{\mathbf{K}=0, p^\pm} \left| \left(\frac{\pi_e^\pm}{m_e} - \frac{\pi_h^\pm}{m_h} \right) \right| \Psi_{\mathbf{K}=0, s} \right\rangle \right|^2 \quad (5)$$

of the transitions from the ground symmetric $1s$ state $D_{00}S$ to different p^\pm states of quasi-2D excitons in $\text{In}_{0.2}\text{Ga}_{0.8}\text{As}/\text{GaAs}$ DQWs with the well and barrier widths $L_1 = L_2 = L_b = 60 \text{ \AA}$ is shown in Fig. 1. Let us discuss the $1s \rightarrow p^+$ transitions. As expected, the strongest $D_{00}S \rightarrow D_{10}S$, $D_{00}S \rightarrow I_{10}S$ transitions to the first electron Landau level lie above the electron CR (the long-dashed line). Only for the $D_{00}S \rightarrow D_{10}S$ transition does $|f|^2$ increase with B . This is because in high fields B the initial state: (i) becomes essentially direct; thus the transition $D_{00}S \rightarrow I_{10}S$ does not gain in intensity, and (ii) it belongs predominantly to the zeroth Landau levels; thus transitions to $3p^+, \dots$ states are weak $\sim (l_B/a_B)^2$. At $B < 4 \text{ T}$ the behavior is complicated by a number of anticrossings between the direct np^+ and indirect $n'p^+$ excitons with $n' > n$. This leads to small splittings between lines and to a visible redistribution of the oscillator strengths (cf. with the interband transitions.^{1,3}). The behavior of the $1s \rightarrow p^-$ transitions (Fig. 1b) are qualitatively similar, except that heavy-hole nonparabolicity results in a strongly sublinear behavior of the transition energies versus B .

The binding energies of the indirect excitons do not grow so fast with B as for the direct excitons. As a result, at a fixed high \mathcal{E} , increasing B induces reentrant Landau-level dependent indirect-direct crossovers^{1,3} (see also Ref. 5). These reentrant transitions show up in the excitonic FIR spectra. The evolution of the energies and dipole matrix elements of the $1s \rightarrow p^+$ transitions with increasing B at electric $\mathcal{E} = 7 \text{ kV/cm}$ is shown in Fig. 2a and that for $\mathcal{E} = 17.2 \text{ kV/cm}$ is shown in Fig. 2b. Note that the transverse field \mathcal{E} breaks

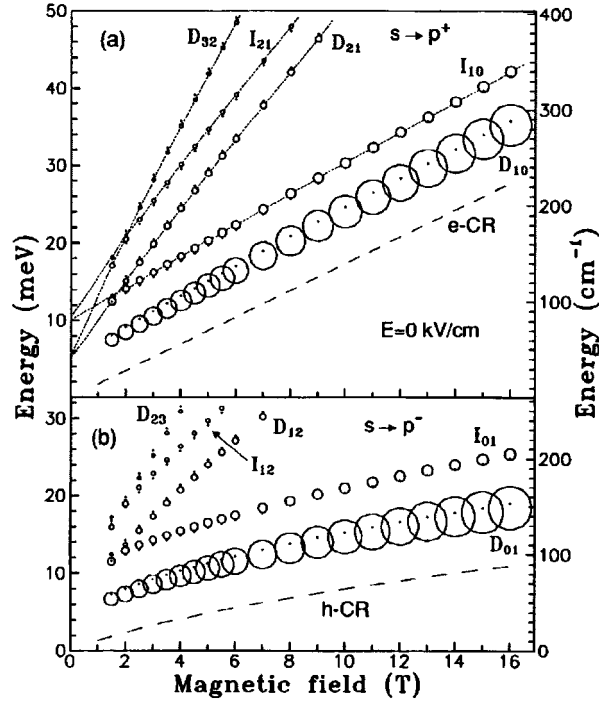


FIG. 1. The evolution of the transitions from the ground symmetric $1s$ state $D_{00}S$ to excited p^+ states (a) and to p^- states (b) with the magnetic field B at $\mathcal{E}=0$. The area of open circles is proportional to the dipole transition matrix elements $|f|^2$, Eq. (5). The lines with short dashes are guides to the eye showing positions of several weak transitions. The plots are labeled with the predominant character of the final states. Small dots correspond to forbidden ($S \rightarrow A$) transitions to antisymmetric states.

the inversion symmetry, thus making all $s \rightarrow p^\pm$ transitions in a symmetric DQW allowed. At higher \mathcal{E} (Fig. 2b) the initial $1s$ state is the indirect exciton I_{00}^- . No indirect-direct crossover is induced in the range of magnetic fields considered, $B < 16$ T. Thus only the transition $I_{00}^- \rightarrow I_{10}^-$ has a large dipole transition matrix element, which increases rapidly (approximately linearly) with B . Transitions to all other higher-lying states are strongly suppressed.

At lower electric fields (Fig. 2a), the energy separations between the I_{10}^- and D_{10}^\pm states are smaller. Also, the direct and indirect states are more mixed.³ Thus the $I_{00}^- \rightarrow D_{10}^+$ transition has a noticeable dipole matrix element even at intermediate B . At $B > 10$ T this transition rapidly starts to gain in oscillator strength. This is due to the indirect-direct crossover for the initial state (the ground state gradually changes^{1,3} from I_{00}^- to D_{00}^-). For the $2p^\pm$ states, since the excitonic effects are diminished, such a crossover occurs at much higher B . Note that the transition to the D_{10}^- final state remains rather weak. This is due to rather different spatial characteristics of the (almost degenerate) D_{10}^- and D_{10}^+ excitons. By considering the conventional probability distributions

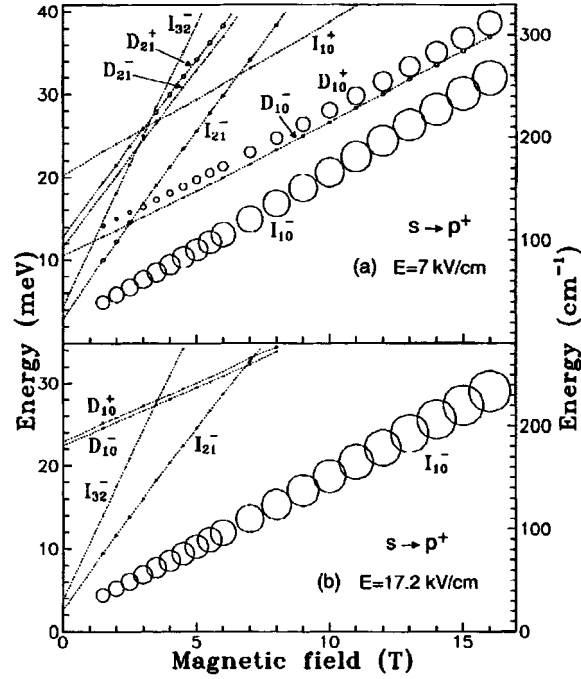


FIG. 2. The evolution of the $1s \rightarrow p^+$ transition energies and dipole matrix elements with the magnetic field B at two different electric fields: $\mathcal{E} = 7$ kV/cm (a), and $\mathcal{E} = 17.2$ kV/cm (b). The plots are labeled with the predominant character of the final states.

$$P_{\mathbf{K}=0, l_z}(z_e, z_h) = \int d^2\rho |\Psi_{\mathbf{K}=0, l_z}(\mathbf{r}_e, \mathbf{r}_h)|^2, \quad (6)$$

it can be shown that at $\mathcal{E} = 7$ kV/cm and $B = 10$ T the ground state in the zeroth Landau levels (i.e., the initial state in the transitions considered) is predominantly spatially direct, and we label it D_{00}^- . It is the intrawell exciton that resides in the left QW ($z_e, z_h < 0$), with a large admixture of the indirect component ($z_e > 0, z_h < 0$) and an extremely small admixture of the direct component in the right well ($z_e, z_h > 0$). At the same time, the ground $2p^+$ state, due to diminished excitonic effects, is polarized in \mathcal{E} : it is predominantly the spatially indirect exciton I_{10}^- ($z_e > 0, z_h < 0$), with a small admixture of the direct ($z_e, z_h < 0$) component. The dipole matrix element of the $D_{00}^- \rightarrow I_{10}^-$ transition is large due to a large spatial overlap between the two states. The next two excited $2p^+$ states (D_{10}^- and D_{10}^+), are predominantly direct excitons that reside in the right and in the left QWs, respectively. Thus only the $D_{00}^- \rightarrow D_{10}^+$ transition is strong, while the $D_{00}^- \rightarrow D_{10}^-$ transition is rather weak (small spatial overlap).

It should be noted that the total FIR absorption not only reflects the dipole transition matrix elements $|f|^2$ but also depends on the population of the excitonic state. The latter is determined by the excitonic radiative lifetimes, which should be strongly influenced by

the direct–indirect crossovers. Our theoretical predictions can be helpful, in particular, for separating these two effects.

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