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FAR-INFRARED $s \to p^{\pm}$ INTER-EXCITON TRANSITIONS IN InGaAs/GaAs COUPLED DOUBLE QUAMTUM WELLS

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We consider $s \to p^{\pm}$ inter-exciton far-infrared (FIR) magneto-optical transitions in coupled double quantum wells (DQWs). Spatially direct (intrawell) and indirect (interwell) excitons in strained $In_xGa_{1-x}As/GaAs$ symmetric DQWs with simple valence band are considered. The evolution of transition energies and oscillator strengths with the perpendicular magnetic field B in different regimes is studied: we consider the direct regime (zero and low perpendicular electric field $\mathcal{E}$), the indirect regime (high $\mathcal{E}$), and the indirect-direct crossover (induced by increasing B at intermediate $\mathcal{E}$).

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Currently, there is considerable experimental interest in magnetoexcitons (MXs) in coupled double quantum wells (DQWs) ([1,2]). In these systems, under the application of an external perpendicular electric field $\mathcal{E}$, the ground state becomes the indirect (interwell) exciton which has large radiative lifetimes. This opens a possibility to observe collective effects in neutral e–h systems in high B at low excitonic temperatures. Spectroscopic identification of collective effects requires a detailed knowledge of magneto-optics of excitons. Theory of interband magneto-optics of MXs in DQWs developed by us [3] is in good agreement with experiment [1]. Intraband FIR magneto-spectroscopy 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 theoretically study FIR magneto-optical transitions of such excitons in $B > 1\,\mathrm{T}$.

We consider symmetric strained $In_xGa_{1-x}As/GaAs$ DQWs. Due to the strain, the light-hole band is split off in such structures. As a result, the peculiarities in the excitonic spectra in a DQW are entirely due to the coupling through the barrier and 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, at zero and low B, the intermediate barrier regime is realized [1 3]. 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, $p^{\pm}, \ldots$ 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 state splits to the symmetric (S) and antisymmetric (A) under the simultaneous inversion $\{z_e \to -z_e, z_h \to -z_h\}$ states [4,3]. In the wide barrier regime, the excitonic symmetric-antisymmetric splittings $\Delta_X \simeq \Delta_e\Delta_h/\delta E_{DI}$

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are determined by the two-particle e-h co-tunneling through the barrier [3]. Thus, Δ_X are suppressed by the excitonic effects ($\sim \delta E_{DI}^{-1}$), in particular, decrease with increasing B.

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

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

Here "+" is for the left ($\Delta l_z = 1$) and "−" is for the right ($\Delta l_z = -1$) circularly polarized light; $\pi_j^{\pm} = \pi_{jx} \pm i\pi_{jy}$ ($\pi_j = -i\hbar\nabla_{\rho_j} - \frac{e_j}{c}\mathbf{A}_j$, $j = e, h$) 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 magnetic [6] 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 eigenstates, we use the expansion [3] of the wave functions of MXs in a DQW with angular momentum projection l_z (e.g., ≥ 0)

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

Here $l_B = (\hbar c/eB)^{1/2}$, $\zeta_i(z_e)$ are the electron and $\xi_j(z_h)$, hole wave functions describing the free perpendicular motion in a DQW in $\mathcal{E}$. $\phi_{nm}(\rho)$ are factored wave functions in B; $\rho = \rho_e - \rho_h$ is the in-plane e–h separation. For the MX, the quantum numbers n and m correspond to the e and h Landau levels, respectively; 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 more details see Ref.3. In high magnetic fields, $1s$ excitonic states are predominantly formed by the orbital ϕ_{00} corresponding to zero 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, $2p^+$ ($2p^-$) excitonic states are formed predominantly by the orbital ϕ_{10} (ϕ_{01}) with small admixture of $\phi_{n+1\,n}$ ($\phi_{n\,n+1}$) 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), in 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

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

(Note that excitons with $\mathbf{K} \neq 0$ cannot be characterized in B by the projection of angular momentum of the relative e–h motion l_z. As a result, the selection rules are rather relaxed. The effects associated at finite temperatures with FIR absorption by $2D$ MXs with $\mathbf{K} \neq 0$ will be discussed elsewhere [7]. In high B, the inter-exciton transition $1s \to 2p^+$ ($1s \to 2p^-$) can be thought of as the electron (hole) cyclotron resonance (CR) $\phi_{00} \to \phi_{10}$ ($\phi_{00} \to \phi_{01}$), 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 inter-exciton transition energies are *larger* than the bare CR energies $\hbar\omega_{ce(h)} = eB/m_{e(h)}$. In the strictly-2$D$ case and high magnetic field limit, the binding energies of the $1s$ and $2p^{\pm}$ MXs are [8] $E_{00} = E_0$ 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 $1s \to 2p^{\pm}$ in this limit are

$$E_{1s\to 2p^+} = \hbar\omega_{ce} + \frac{1}{2}E_0\,,\ E_{1s\to 2p^-} = \hbar\omega_{ch} + \frac{1}{2}E_0\,. \tag{4}$$

Excitonic FIR transitions are affected by the Coulomb interparticle e–h interaction. Kohn's theorem [9] is inapplicable because the charge-to-mass ratios are different for the e and h. It can be shown however that, at any B (for simple parabolic bands and excitons with $\mathbf{K} = 0$), the *difference* $E_{s\to p^+} - E_{s\to p^-} = \hbar\omega_{ce} - \hbar\omega_{ch}$ does not depend on the e–h interaction (cf. with the theorem [10] for one-component electron systems in B).

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

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

of the transitions from the ground symmetric $1s$ state $D_{00}S$ to different $p^{\pm}$ states of quasi-2D excitons in $In_{0.2}Ga_{0.8}As/GaAs$ DQWs with the well and barrier widths $L_1 = L_2 = L_b = 60$ Å is shown in Fig.1. Let us discuss the $1s \to p^+$ transitions. As expected, the most strong $D_{00}S \to D_{10}S$, $D_{00}S \to I_{10}S$ transitions to the first electron Landau level lie above the electron CR (the long-dashed line). Only for the $D_{00}S \to D_{10}S$ transition, $|f|^2$ increases with B. This is because in high B the initial state: (i) becomes essentially direct; thus the transition $D_{00}S \to I_{10}S$ does not gain the intensity, and (ii) it belongs predominantly to zero Landau levels; thus transitions to $3p^+, \ldots$ states are weak $\sim (l_B/a_B)^2$. At $B < 4$ 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 visible redistribution of the oscillator strengths (cf. with the interband transitions. [1,3]) The behavior of the $1s \to p^-$ transitions [Fig.1*b*] are qualitatively similar, except that heavy-hole nonparabolicity results in a strong sublinear behavior of 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 energies and dipole matrix elements of the $1s \to p^+$ transitions with increasing B is shown at electric field $\mathcal{E}$=7 kV/cm in Fig.2*b*, and at $\mathcal{E}$=17.2 kV/cm in Fig.2*b*. Note that the perpendicular $\mathcal{E}$ breaks inversion symmetry thus making all $s \to p^{\pm}$ transitions in a symmetric DQW allowed. At higher $\mathcal{E}$ [Fig.2*b*], the initial $1s$ state is the indirect exciton I_{00}^-. No indirect-direct crossover is induced in the considered range of magnetic fields $B < 16$ T. Thus, only the transition $I_{00}^- \to I_{10}^-$ has a large dipole transition matrix element which rapidly increases (approximately linear) with B. Transitions to all other higher-lying states are strongly suppressed.

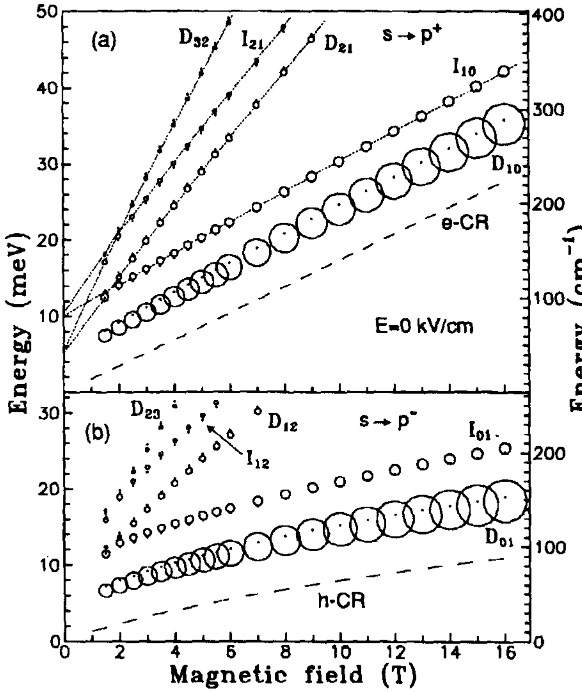

Fig.1.The evolution of the transitions from the ground symmetric $1s$ state $D_{00}S$ to excited p^+ states (a), and to p^- states (b) with B at $\mathcal{E}=0$. The area of open circles is proportional to the dipole transition matrix elements $|f|^2$, Eq. (5). The lines with short dashes are guides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

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 $I_{00}^- \to D_{10}^+$ transition has a noticeable dipole matrix element even at intermediate B. At $B > 10$ T this transition rapidly starts to gain the oscillator strength. It 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. Considering conventional probability distributions

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

it can be shown that at $\mathcal{E}=7$ kV/cm and $B=10$ T the ground state in the zero Landau levels (i.e., the initial state in the transitions considered), is predominantly spatially direct, 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 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 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 $D_{00}^- \to I_{10}^-$ 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,

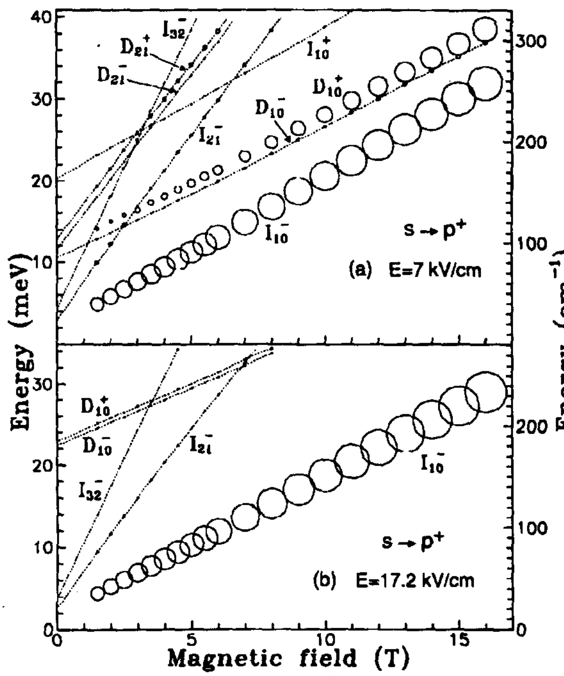

Fig.2. The evolution of the $1s \to p^+$ transition energies and dipole matrix elements with the magnetic field B at two different electric fields: $\mathcal{E} = 7$ kV/cm (a), and $\mathcal{E} = 17.2$ kV/cm (b). Labeling shows the predominant character of the final states

respectively. Thus only the $D_{00}^- \to D_{10}^+$ transition is strong, while the $D_{00}^- \to D_{10}^-$ transition is rather weak (small spatial overlap).

It should be noted that the total FIR absorption reflects not only 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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