Energy spectra and photoluminescence of charged magneto-excitons

Arladiusz Wójc*‡, John J. Quinn* and Paweł Hawrylak‡

*Department of Physics, University of Tennessee, Knoxville, Tennessee 37996, USA
‡Institute of Physics, Wroclaw University of Technology, Wroclaw 50-370, Poland
§Institute for Microstructural Sciences, National Research Council of Canada, Ottawa, Ontario K1A 0R6, Canada

Charged magneto-excitons \( X^- \) in a dilute 2D electron gas in narrow and symmetric quantum wells are studied using exact diagonalization techniques. An excited triplet \( X^- \) state with a binding energy of about 1 meV is found. This state and the singlet are the two optically active states observed in photoluminescence (PL). The interaction of \( X^- \)'s with electrons is shown to have short range, which effectively isolates bound \( X^- \) states from a dilute e–h plasma. This results in the insensitivity of PL to the filling factor \( \nu \). For the “dark” triplet \( X^- \) ground state, the oscillator strength decreases exponentially as a function of \( \nu^{-1} \) which explains why it is not seen in PL.

Keywords: Charged magneto-exciton; Low-dimensional structure; Fractional quantum Hall effect

Recent magneto-photoluminescence (PL) experiments [1] showing recombination of charged excitons \( X^- \) (two electrons bound to a valence hole) in narrow GaAs quantum wells appear to disagree completely with theoretical prediction [2]. According to theory, the singlet (spin unpolarized) state \( X^- \) is the \( X^- \) ground state (GS) at low magnetic field, while the triplet \( X^- \) is the GS at fields above 30 T. In the PL experiments, the \( X^- \) appears to be the GS for all magnetic fields, and \( X^- \) and \( X^- \) have comparable PL intensity. Here we present results of numerical diagonalization of small systems, including effects of Landau level mixing and finite well widths. We find that the energy of the lowest triplet state \( (X^-_{td}) \) behaves exactly as predicted by previous calculations, but that its PL intensity is orders of magnitude smaller than those of the \( X^- \) and an excited triplet state \( (X^-_{td}) \). We suggest that the triplet observed in PL is this bright triplet \( X^-_{bd} \) whose energy is always higher than that of \( X^- \). The dark triplet \( X^-_{bd} \) is not observed in PL, and no disagreement exists between theory and experiment.

The energy and PL spectra of the \( X^- \) are calculated by exact numerical diagonalization of the two-electron–one-hole (2e–1h) Hamiltonian [3]. In order to preserve the 2D translational symmetry of an infinite quantum well (QW) in a finite-size calculation, we use Haldane’s [4] spherical geometry. The magnetic field \( B \) perpendicular to the surface of the sphere of radius \( R \) is due to a magnetic monopole placed in the center. The monopole strength \( 2S \) is defined in the units of elementary flux \( \phi_0 = hc/e \), so that \( 4\pi R^2 B = 2S\phi_0 \), and the magnetic length is \( \lambda = R/\sqrt{S} \). The electron and hole states form degenerate angular momentum (l) shells or Landau levels (LL), and the lowest LL has \( l = S \). Our model applies to narrow and symmetric heterostructures, and the numerical results presented here are for the GaAs QW of width 11.5 nm. For such a system, only the lowest QW subband need be included, and the cyclotron motion of both electrons and holes is well described in the effective-mass approximation, with the inter-subband coupling partially taken into account through the dependence of the hole cyclotron mass on \( B \). The finite (and different) widths of electron and hole quasi-2D layers are included through effective 2D interaction potentials. The electron Zeeman energy depends on well width and \( B \). Five electron and hole LL’s are used in the calculation, and the energies obtained for different values of \( 2S \leq 20 \) are extrapolated to the limit of \( S^{-1} = (l/2R)^2 \to 0 \) (i.e. to the planar geometry), so that the finite-size and surface-curvature effects are eliminated.

The 2e–1h energy spectra (binding energy as a function of angular momentum \( L \)) calculated for \( 2S = 20 \) are shown in Fig. 1. Open and full symbols mark singlet and triplet states \( (J_e \) is the total electron spin), and each state with \( L > 0 \) represents a degenerate \( L \)-multiplet. Since the PL process (annihilation of an e–h pair and emission of a photon) occurs with conservation of angular momentum, only states from the \( L = S \) channel are radiative [5]. Recombination of other, non-radiative states requires breaking rotational symmetry (e.g., by collisions with electrons). This result is independent of the chosen spherical geometry and holds also for a planar QW, except that the definition of \( L \) is different.

Three states marked in Fig. 1 are of particular importance: \( X^-_{s} \) and \( X^-_{t} \) are the only bound radiative states, while \( X^-_{td} \) has by far the lowest energy of all non
radiative states. The radiative triplet bound state $X_{tb}^-$ is identified for the first time. The binding energies of all three $X^-$ states are extrapolated to $\lambda/R \to 0$ and plotted in Fig. 2(a) as a function of $B$. For the $X^-$, the binding energy differs from the PL energy (thin dotted line) by the Zeeman energy needed to flip one electron’s spin, and the cusp at $B \approx 42$ T is due to the change of sign of the electron $g$-factor. For the triplet states, the PL and binding energies are equal. The energies of $X^-$ and $X_{td}^-$ behave as expected: The binding of $X^-$ weakens at higher $B$ due to the “hidden symmetry” [6], which eventually leads to its unbinding in the infinite field limit [7]; the binding energy of $X_{td}^-$ changes as $e^2/\lambda \propto \sqrt{B}$; and the predicted [2] transition from the $X^-_{s}$ to the $X_{td}$ GS at $B \approx 30$ T is confirmed. The new $X_{tb}$ state remains an excited triplet state at all values of $B$, and its binding energy is smaller than that of $X^-$ by about 1.5 meV. The oscillator strengths $\tau^{-1}$ of a neutral exciton $X$ and the two radiative $X^-$ states are plotted in Fig. 2(b). In the $2e-1h$ spectrum, the well bound $X^s$ and $X_{tb}$ states share a considerable part of the total oscillator strength of one $X$, with $\tau_{tb}^{-1}$ nearly twice larger than $\tau_{s}^{-1}$.

The comparison of calculated magnitude and magnetic field dependence of the $X^-$ binding energies with the experimental PL spectra [16,17], as well as high oscillator strength of the $X_{tb}$, lead to the conclusion that the three peaks (without counting the Zeeman splittings) observed in PL experiments are due to the recombination of $X^-$, $X_{s}^-$, and $X_{tb}^-$. Due to the vanishing oscillator strength, the lowest triplet state $X_{sd}^-$ found in earlier calculations [2][1] remains undetected even at $B > 30$ T, when it is expected to be the $X^-$ GS. Only partial hole spin polarization at lower $B$ and its increase with increasing $B$ can lead to an observed [1] enhancement of the $X_{tb}$ PL intensity by up to a factor of two, while the intensity of the $X^s$ peak remains roughly unchanged.

The results in Figs. 1 and 2 are quantitatively correct for narrow and symmetrically (or remotely) doped QW’s. In strongly asymmetric QW’s or heterojunctions [1], significant difference between electron and hole QW confinements increases the $e-h$ attraction compared to the $e-e$ repulsion, and the binding energies of all three $X^-$ states contain an uncompensated $e-h$ attraction which scales with $B$ like the exciton energy ($c^2/\lambda \propto \sqrt{B}$). Nevertheless, our most important qualitative result remains valid for all structures: The triplet $X^-$ state seen in PL is the bright excited triplet state $X_{tb}^-$ and not the lowest triplet state $X_{td}^-$. To understand why the $X_{td}^-$ state remains optically inactive even in the presence of collisions, the $e-X^-$ interactions must be studied in greater detail. In Fig. 2 we plot the energy spectra of a $3e-1h$ system. As in Fig. 1, the energy is measured from the exciton energy and the open and filled circles mark multiplets with different $J_e$. In the low energy states, bound $X^-$ complexes interact with an electron through the effective pseudopotentials $V(L)$, defined as the dependence of pair interaction energy on pair angular momentum. The pair angular momentum $L$ is related to the average $e-X^-$ separation $d$, and, on a sphere) larger $L$ corresponds to smaller $d$. The allowed values of $J_e$ and $L$ can be understood by addition of spins and angular momenta of an appropriate $X^-$ and an electron. The total energy of an interacting pair is the sum of the $e-X^-$ repulsion energy $V(L)$ and the appropriate binding energy. Because of incompatible energy scales, the $e-X^-$ scattering is nearly decoupled from internal $X^-$ excitations, and $V(L)$ is similar for all pairs. The relative position of $3e-1h$ energy bands corresponding to different $X^-$ complexes depends on the involved binding energy (and hence on $B$).

In narrow ($\leq 20$ nm) QW’s, the $e-X^-$ pseudopotential retains the short-range character which results in the Laughlin correlations [13]. In a 2D electron gas, these correlations are responsible for the occurrence of the incompressible liquid states and the fractional quantum Hall effect. Similar $e-X^-$ correlations in the $e-h$ plasma limit angular momentum (and energy) of $e-X^-$ collisions and, at low density and temperature, forbid an $X^-$ from getting close to an electron, effectively isolating it from the surrounding electron gas [10,11]. This result depends critically on the short-range nature of $V(L)$, and thus on the QW thickness (in thicker QW’s, high energy collisions
occur even at low density). The Laughlin correlations at the filling factor \( \nu < m^{-1} \) (where \( \nu^{-1} \) is the number of magnetic flux quanta per electron) mean that all pair states with \( L > 2S - m \) are avoided. This relates \( \nu \) (i.e. density) to the maximum allowed \( L \) (i.e. minimum distance) for an \( e^{-X^-} \) pair.

In Fig. 4 we plot the PL oscillator strength and energy (measured from the exciton energy) calculated for some of the \( e^{-X^-} \) states marked in Fig. 3. We assume that the Zeeman energy will polarize all electron spins prior to recombination, except for those two in the \( X^+ \), and concentrate on the following three initial configurations: \( e^{-X^+} \) with \( J_e = \frac{1}{2} \), and \( e^{-X^+} \) and \( e^{-X^-} \) with \( J_e = \frac{3}{2} \). For each of the three configurations, \( \tau^{-1} \) and energy are plotted as a function of \( L \) (i.e. of \( \nu \)). The \( e^{-X^-} \) interactions have no significant effect on the PL oscillator strength and energy of an \( X^- \) at small \( L \) (i.e., at low density). This justifies a simple picture of PL in dilute \( e^{-h} \) plasmas. In this picture, recombination occurs from a single isolated bound complex and hence is virtually insensitive to \( \nu \). Quite surprisingly, the Laughlin correlations prevent increase of the \( X^- \) oscillator strength \( \tau_{td}^{-1} \) through collisions with other charges. The \( \tau_{td}^{-1} \) decreases exponentially (see insets in Fig. 4) with decreasing \( \nu \), and \( \tau_{td} \) remains ten times longer than \( \tau_s \) even at \( \nu = \frac{1}{3} \). This explains the absence of an \( X^- \) peak even in those PL spectra showing strong recombination of a higher-energy triplet state \( X^+ \).

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