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Energy spectra and photoluminescence of charged magneto-excitons

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Abstract

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 ν. For the "dark" triplet X^- ground state, the oscillator strength decreases exponentially as a function of ν^{-1} which explains why it is not seen in PL.

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Recent magneto-photoluminescence (PL) experiments [1] showing recombination of charged excitons X^- (two electrons bound to a valence hole) in narrow GaAs quantum wells (QW) appear to disagree completely with theoretical prediction [2]. According to theory, the singlet (spin unpolarized) state X_s^- is the X^- ground state (GS) at low magnetic field, while the triplet X_t^- is the GS at fields above 30 T. In the PL experiments, the X_s^- appears to be the GS for all magnetic fields, and X_s^- and X_t^- 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_s^- and an excited triplet state (X_{tb}^-). We suggest that the triplet observed in PL is this bright triplet X_{tb}^- whose energy is always higher than that of X_s^-. The dark triplet X_{td}^- 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

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two-electron–one-hole (2e–1h) Hamiltonian [3–5]. In order to preserve the 2D translational symmetry of an infinite QW in a finite-size calculation, we use Haldane's [6,7] 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 $l = S$.

We use a model in which the QW is symmetric and relatively narrow. The numerical results presented here are for a GaAs QW of width 11.5 nm. For this 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. Inter-subband coupling is partially taken into account through the dependence of the hole cyclotron mass on B. The finite (and different) widths of electron and hole envelope functions are included through effective 2D interaction potentials [8]. The Zeeman energy depends on well width and B. Five electron and hole LLs are used in the calculation, and the energies obtained for different values of $2S$ are extrapolated to the limit of $S^{-1} = (\lambda/R)^2 \to 0$ (i.e. to the planar geometry), so that the finite-size and surface-curvature effects are eliminated.

The 2e–1h energy spectra (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 [3,4]. 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 [7].

The occurrence of a strict PL selection rule at finite B may seem surprising, since the "hidden symmetry" [9,10] that forbids the X_{td}^- recombination in the lowest LL does not hold when the mixing with higher LLs is included. However, it is both the hidden symmetry and the above-mentioned angular momentum conservation that independently forbid the X_{td}^- recombination, and the latter remains valid at finite B. The hidden symmetry follows from equal magnitude of e–e, e–h, and h–h interactions in the lowest LL and leads to decoupling of optically active ($L = 0$) excitons from the remaining e–h system. The number N_X of such decoupled excitons is conserved, and only the so-called "multiplicative" e–h states which have $N_X > 0$ are radiative (the recombination occurs from a decoupled X at the bare exciton energy). The "non-multiplicative" states have $N_X = 0$ and are not radiative. They include the only bound X^- state of an 2e–1h system in the lowest LL. At a finite B, this state is the X_{td}^- in Fig. 1. Although the hidden symmetry and resulting N_X conservation law no longer hold at finite B, the X_{td}^- recombination remains strictly forbidden because of the independently conserved L.

We expect breaking of both symmetries for real experimental situations. Different electron and hole effective masses and finite well widths cause e–e, e–h, and h–h interactions to differ. Furthermore, e–X_{td}^- scattering during the recombination in the presence of excess electrons can relax the strict conservation of the X^- angular momentum in the radiative decay. However, for narrow and symmetric QWs containing a relatively small number of excess electrons, the symmetries may be only weakly broken and some remnant of the strict conservation laws may survive.

Three states marked in Fig. 1 are of particular importance: X_s^- and X_{tb}^- are the only strongly 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. From the mapping [7] between the 2D algebra of the total angular momentum on a sphere (L and L_z) and the pair of good quantum numbers on a plane (total, center-of-mass, and relative angular momenta: M, M_{CM}, and $M_{REL} = M - M_{CM}$), we find that on a plane the states X_s^- and X_{tb}^- have $M_{REL} = 0$, while X_{td}^- has $M_{REL} = 1$. 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_s^-, 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_s^- and X_{td}^- behave

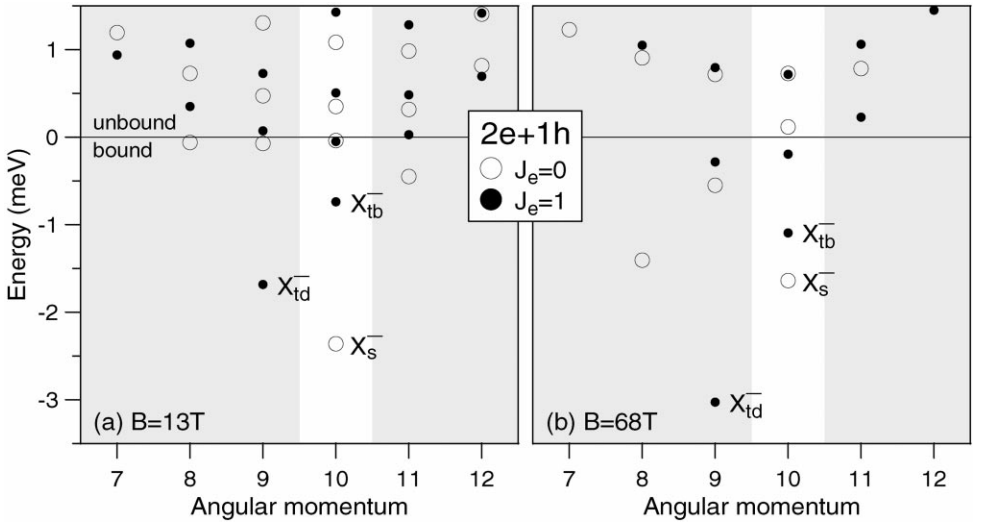

Fig. 1. The energy spectra (binding energy versus angular momentum) of the 2e–1h system on a Haldane sphere with the Landau level degeneracy of $2S+1=21$. The parameters are appropriate for the 11.5 nm GaAs quantum well.

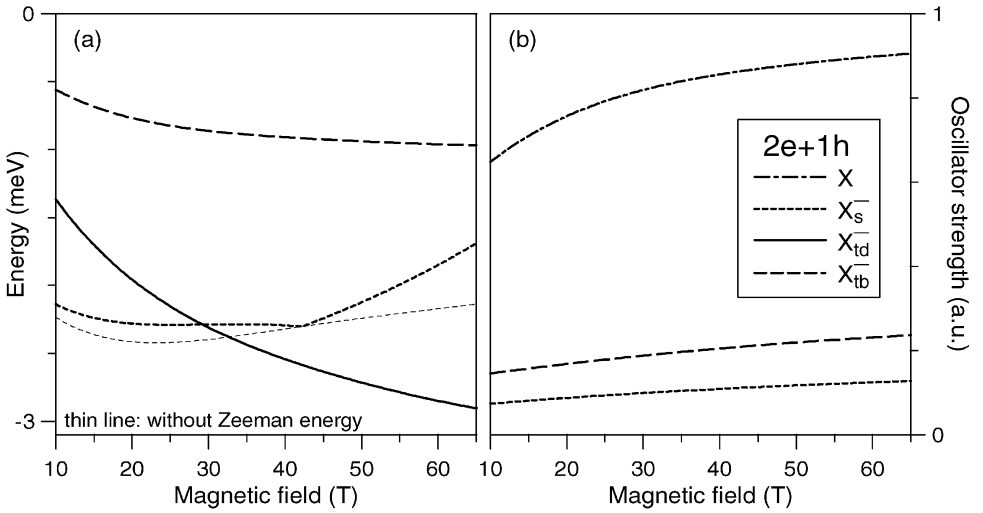

Fig. 2. The X^- binding energies (a) and oscillator strengths (b) in the 11.5 nm GaAs quantum well plotted as a function of the magnetic field.

as expected: The binding of X_s^- weakens at higher B due to the "hidden symmetry" [9,10], which eventually leads to its unbinding in the infinite field limit [11]; 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_s^- by about 1.5 meV. The oscillator strengths τ^{-1} of a neutral exciton X and the two radiative X^- states are plotted in Fig. 2(b). In the 2e–1h spectrum, the strongly bound X_s^- and X_{tb}^- states share a considerable part of the total oscillator strength of one X, with τ_{tb}^{-1} nearly twice larger than τ_s^{-1}.

The comparison of calculated magnitude and magnetic field dependence of the X^- binding energies with the experimental PL spectra [1,12–14], as well as high oscillator strength of the X_{tb}^-, lead to the

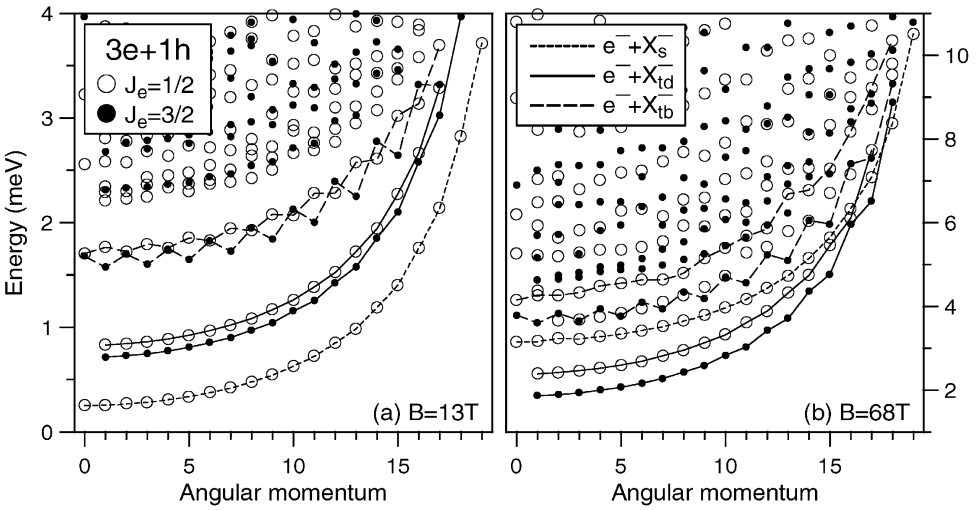

Fig. 3. The energy spectra (energy versus angular momentum) of the 3e–1h system on a Haldane sphere with the Landau level degeneracy of $2S+1=21$. The parameters are appropriate for the 11.5 nm GaAs quantum well.

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_{td}^- found in earlier calculations [2,15,16] 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 doped QWs. In strongly asymmetric QWs or heterojunctions [17], significant difference between electron and hole QW confinements alters the relative strengths of e–h attraction and e–e repulsion, and the binding energies of all three X^- states contain additional uncompensated e–e or e–h interaction. 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. 3 we plot the energy spectra of an 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 ($\leqslant 20$ nm) QWs, the e–X^- pseudopotential retains the short-range character which results in the Laughlin correlations [18]. 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,

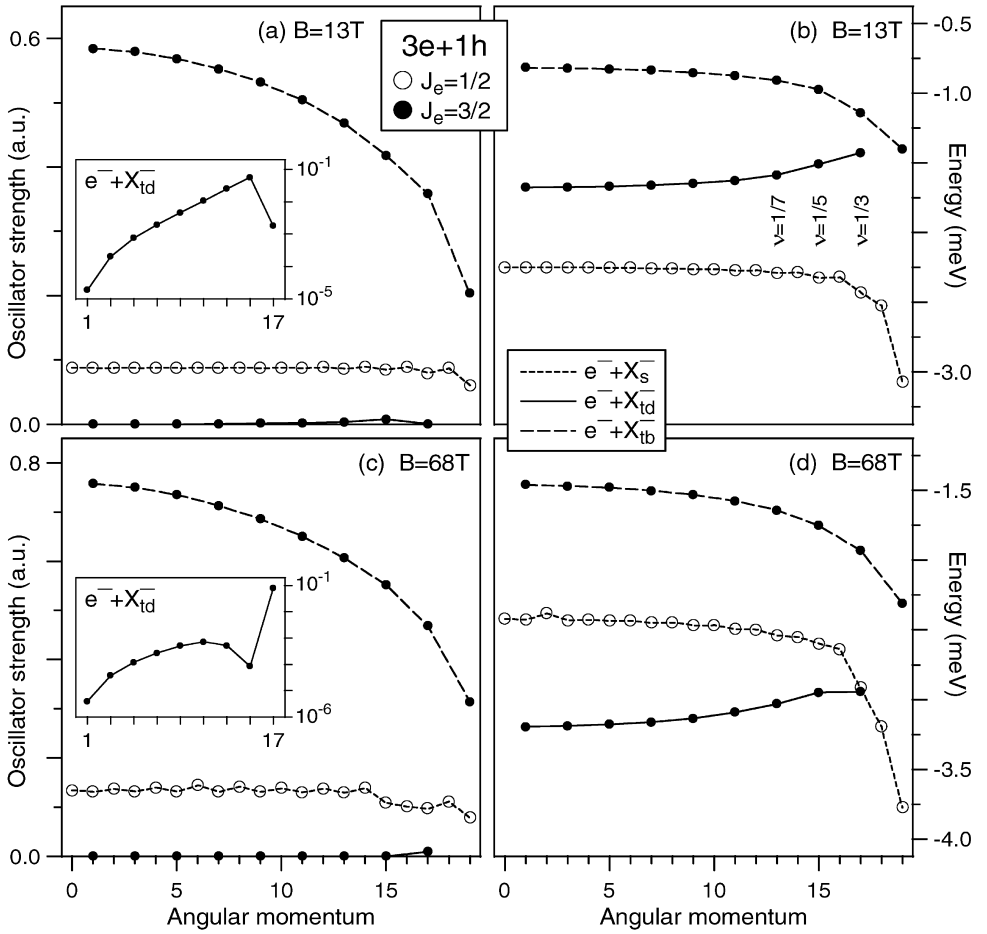

Fig. 4. The oscillator strengths (left) and recombination energies (right) of an X^- interacting with an electron on a Haldane sphere with the Landau level degeneracy of $2S+1=21$, plotted as a function of the e–X^- pair angular momentum. The parameters are appropriate for the 11.5 nm GaAs quantum well.

effectively isolating it from the surrounding electron gas [3–5]. This result depends critically on the short-range nature of $V(L)$, and thus on the well width (in wider wells, high-energy collisions occur even at low density). The Laughlin correlations at the filling factor $\nu \leqslant m^{-1}$ (where ν^{-1} is the number of magnetic flux quanta per electron) mean that all pair states with $L > 2S - m$ are avoided [19–22]. This relates ν (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_s^-, and concentrate on the following three initial configurations: e–X_s^- with $J_e = \frac{1}{2}$, and e–X_{tb}^- and e–X_{td}^- with $J_e = \frac{3}{2}$. For each of the three configurations, τ^{-1} and energy are plotted as a function of L (i.e. of ν). 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 ν. Somewhat surprisingly,

the Laughlin correlations prevent substantial increase of the X_{td}^- oscillator strength τ_{td}^{-1} through collisions with other charges. The τ_{td}^{-1} decreases exponentially (see insets in Fig. 4) with decreasing ν, because the $n = 1, 2, \ldots$ largest pseudopotential coefficients are avoided when $\nu \leqslant (2n+1)^{-1}$. As a result, τ_{td} remains ten times longer than τ_s even at $\nu = \frac{1}{3}$. This explains the absence of an X_{td}^- peak even in those PL spectra [1,12–14,17] showing strong recombination of a higher-energy triplet state X_{tb}^-.

Acknowledgements

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