Negatively charged magnetoexcitons in quantum dots

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We study the interaction of an electron with an exciton in a disk-shaped quantum dot subjected to magnetic and electric fields perpendicular to the plane of the dot. The absorption spectrum of magnetoexcitons and negatively charged magnetoexcitons is calculated. The differences between neutral and negatively charged exciton spectra, as a function of the polarization of light and of the applied electric field, allow for the detection of an electron bound to an exciton in a dot.

I. INTRODUCTION

The rapid progress in semiconductor technology in recent years has allowed for the fabrication of low-dimensional electronic nanostructures (quantum dots, wires, etc.). In these nanostructures, the lateral confinement of a quasi-two-dimensional (2D) electron gas leads to discrete energy spectrum observable in low-temperature experiments.[1] While a considerable amount of work has been done both on magnetoexcitons (X) in nanostructures[2–8] and on extrinsic recombination in quantum dots,[9] the interaction of mobile excitons with free carriers[5] is not yet well understood. A number of recent experiments investigated the effect of free carriers on magnetoexcitons in quantum dots. The anomalies in magnetoabsorption spectra observed recently in disordered quantum wells, by Kheng *et al.*,[2] have been attributed to negatively charged excitons. Zrenner *et al.*[3] studied magnetoexcitons in vertically tunable quantum dots containing free carriers. The applied electric field allowed control over the spatial separation between electrons and holes, tuning the ratio $\alpha = V_{e\text{-}h}/V_{e\text{-}e}$ of the electron-hole ($V_{e\text{-}h}$) to electron-electron ($V_{e\text{-}e}$) interaction. A control over the spin of the excited electron is provided by different polarizations of light.[2] Theoretically, the effect of a low free carrier concentration on excitons in bulk semiconductors has been studied by Lampert,[10] who suggested the binding of the second electron by the exciton and the formation of a negatively charged exciton (X^-). The binding of the second electron depended strongly on the hole mass.[11] For an infinite hole mass, the problem reduces to the familiar problem of D^-.[12] The major difference between the D^- and the X^- is the behavior in a magnetic field. A strong magnetic field stabilizes the binding of the second electron in the D^- complex. However, for the mobile X^- in a very strong magnetic field, hidden symmetries[13] prevent the coupling of an exciton to free carriers. The hidden symmetry implies essentially that, in the lowest Landau level, there is an exact decoupling of excitons from the excess of free electrons. To break the hidden symmetry and make excitons a useful probe of free carriers, one must either account for Landau level mixing, break the balance between the electron-electron and electron-hole interactions, or laterally confine carriers.

In this work, the coupling of a mobile magnetoexciton and a free carrier, both confined in the disk-shaped quantum dot, and subjected to magnetic and electric fields perpendicular to the plane of the dot, is examined in detail. We show that the presence of a free carrier in the dot significantly changes the low-energy absorption spectrum of an exciton, allowing for the detection of free carriers.

II. THE MODEL

We consider a system of N (N=1,2) quasi-two-dimensional electrons and a single hole, in the presence of the perpendicular magnetic field, confined in a quantum dot. The hole may move on a plane separated from the plane of electrons by a distance d. The effective separation d may be tuned by an applied electric field. The Hamiltonian describing a magnetoexciton and free carriers can be written as

$$H = \frac{1}{2m_h}\left(\mathbf{p}^h - \frac{e}{c}\mathbf{A}(\mathbf{r}^h)\right)^2 + V_c^h(\mathbf{r}^h) + g^h\mu S_z^h + \sum_{i=1}^{N}\left[\frac{1}{2m_e}\left(\mathbf{p}_i^e + \frac{e}{c}\mathbf{A}(\mathbf{r}_i^e)\right)^2 + V_c^e(\mathbf{r}_i^e) + g\mu S_z^i + V_{e\text{-}h}(|\mathbf{r}_i^e - \mathbf{r}^h|)\right] + \sum_{i<j;i,j=1}^{N} V_{e\text{-}e}(|\mathbf{r}_i^e - \mathbf{r}_j^e|), \quad (1)$$

where m_e, m_h are the effective masses, $\mathbf{r}^{e,h}$ are positions, $p^{e,h}$ are momenta, $\mathbf{A} = \mathbf{B}\times\mathbf{r}/2$ is a vector potential in symmetrical gauge, V_c are the lateral confining potentials, $V_{e\text{-}h}$ is the electron-hole attraction, $V_{e\text{-}e}$ is the electron-electron repulsion, g is the effective g factor, μ is the Bohr magneton, S_z^i is the z component of the ith particle spin, and we take $\hbar = 1$.

III. MAGNETOEXCITONS

Let us begin with a single neutral exciton: X ($N=1$). There are two possible ways of diagonalizing the exciton Hamiltonian. For weak confinement, when the continuous part of the energy spectrum corresponding to scattered electron-hole states plays an important role, the Hamiltonian of the free magnetoexciton should be diagonalized first,[7] and then matrix elements of the confining potentials calculated in the exciton center-of-mass and relative basis will be diagonalized. For strong confinement, one first calculates the states of the confined particles and, next, treats the electron-hole interaction as a perturbation.[8]

In the case of strong parabolic confinement, $V_{e,h} = 1/2m_{e(h)}\omega^2_{e(h)}(r^{e(h)})^2$, the eigenstates $|m,n\rangle$ and eigenenergies $E_{m,n}$ of a noninteracting electrons (holes) are those of a pair of two decoupled harmonic oscillators,[9] $E^{e,h}_{mn} = \Omega_{e(h),+}(n+\frac{1}{2})+\Omega_{e(h),-}(m+\frac{1}{2})$. These have characteristic frequencies $\Omega_{e(h),\pm} = [\sqrt{\omega^2_{e(h),c}+4\omega^2_{e(h),0}} \pm \omega_{e(h),c}]/2$, where $\omega_{e(h),c} = eB/m_{e,h}c$ is the cyclotron frequency (we set $\hbar = 1$). For zero magnetic field $\Omega_{e(h),+} = \Omega_{e(h),-} = \omega_{e(h),0}$, while for high magnetic fields $\Omega_{e(h),+} \simeq \omega_{e(h),c}$ and $\Omega_{e(h),-} \ll \omega_{e(h),c}$ are the inter- and intra-Landau-level frequencies, respectively. The two-particle states are the products of electron and hole single-particle states, $|n_e,m_e,n_h,m_h\rangle = |n_e,m_e\rangle|n_h,m_h\rangle$, with angular momenta $R_e = m_e - n_e$ and $R_h = n_h - m_h$ for an electron and a hole, respectively. The harmonic oscillator basis allows us to conveniently calculate the electron-hole Coulomb interaction matrix elements:

$$\langle n'_e m'_e, n'_h m'_h | V_{e\text{-}h} | n_h m_h, n_e m_e \rangle$$

$$= \frac{-E_0\sqrt{\Omega_{e\text{-}h}}\delta_{R,R'}(-1)^{s_e+s'_e}\Omega_{e\text{-}h}^{(s_e+s'_e+s_h+s'_h)/2}}{\sqrt{n_e!n'_e!n_h!n'_h!m_e!m'_e!m_h!m'_h!\Omega_e^{(s_e+s'_e)}\Omega_h^{(s_h+s'_h)}}}$$
$$\times \sum_{p_e=0}^{min(n_e,n'_e)} \sum_{p_h=0}^{min(n_h,n'_h)} \sum_{l_e=0}^{min(m_e,m'_e)} \sum_{l_h=0}^{min(m_h,m'_h)} \binom{n_e}{p_e}\binom{n'_e}{p_e}\binom{n_h}{p_h}\binom{n'_h}{p_h}\binom{m_e}{l_e}\binom{m'_e}{l_e}\binom{m_h}{l_h}\binom{m'_h}{l_h}$$
$$\times p_e!p_h!l_e!l_h!\left(\sqrt{\frac{\Omega_{e\text{-}h}}{\Omega_h}}\right)^{(s_h+s'_h-2p_h-2l_h)}\left(\sqrt{\frac{\Omega_{e\text{-}h}}{\Omega_e}}\right)^{(s_e+s'_e-2p_e-2l_e)}\left(\frac{-1}{2}\right)^p\frac{\Gamma(p+1/2)}{\sqrt{(2\pi)}}, \quad (2)$$

where $R = R_e + R_h$ is the total angular momentum of the pair, $E_0 = \mathcal{R}\sqrt{2\pi}a_0/l_0$ ($\mathcal{R} = e^2/2\epsilon a_0$ is the effective Rydberg, $a_0 = \hbar^2/me^2$ is the Bohr radius, and $l_0 = \sqrt{\hbar c/eB}$ is the magnetic length), $\Omega_{e(h)} = \sqrt{1+4(\omega_{e(h),0}/\omega_{e(h),c})^2}$, $2\Omega^{-1}_{e\text{-}h} = \Omega_e^{-1}+\Omega_h^{-1}$, $p = n'_e + m_e + n_h + m'_h - p_e - l_e - p_h - l_h$, and $s = m+n$.

The absorption $A(\omega)$ from the initial state $|i\rangle$, as a function of the photon frequency ω, is given by Fermi's golden rule: $A(\omega) = \sum_f |\langle f|\mathbf{P}|i\rangle|^2\delta(E_f - E_i - \omega)$. Denoting the projections of the final state $|f\rangle$ on the basis vectors as $C^f_{n_e m_e n_h m_h} = \langle f|n_e m_e n_h m_h\rangle$ and taking the transition matrix element proportional to the probability amplitude of finding an electron and a hole at the same position and with opposite spins $\langle f|\mathbf{P}|i\rangle \sim \int\int d\mathbf{r}_e d\mathbf{r}_h \phi^f(\mathbf{r}_e,\mathbf{r}_h)\delta(\mathbf{r}_e - \mathbf{r}_h)$, the absorption $A(\omega)$ for zero-angular-momentum magnetoexcitons with polarizations $(\sigma_\pm, L)$ is given by $A(\omega) = \sum_f |\sum_{n,m} C^f_{nm,nm}|^2\delta(E_f - E_i - \omega)$. The coefficients $C^f_{n_e m_e n_h m_h}$ are obtained from a direct diagonalization of the electron-hole Hamiltonian for different polarizations of light, which control the spin configuration of an electron and a hole. Circularly polarized light excites the electron-hole pairs with well defined spin projections $|\uparrow_e\downarrow_h\rangle$, $|\downarrow_e\uparrow_h\rangle$, while linearly polarized light (L) leads to a creation of states with symmetrical spin functions $(|\uparrow_e\downarrow_h\rangle + |\downarrow_e\uparrow_h\rangle)/\sqrt{2}$ (antisymmetrical spin states are not optically active).

While the magnetoexciton spectrum has been calculated by, e.g., Halonen *et al.*[7] and by Dzyubenko and Sivachenko,[8] it is interesting to ask to what extent the exciton absorption spectrum $A(\omega)$ reflects the electron/hole energy spectrum. The most interesting feature of the electron (hole) energy spectrum $E^{e,h}_{mn} = \Omega_{e(h),+}(n+\frac{1}{2}) + \Omega_{e(h),-}(m+\frac{1}{2})$ is the partial restoration of symmetries by the magnetic field. This restoration of symmetries leads to gaps in the density of states for the special values of the magnetic field at which $r = \Omega_+/\Omega_-$ is an integer. If the confining length $l_{e,h} = \sqrt{1/(2m_{e,h}\omega_{e,h})}$ is the same for electrons and holes, then $r_e = r_h$ and degenerate peaks separated by gaps will appear in the electron-hole spectrum. When $V_{e\text{-}h}$ is strong, the energy quantization of the system is determined by the interaction, and no special behavior at any value of magnetic field can be expected. This is illustrated in Fig. 1 where we show the absorption spectra $A(\omega)$ for magnetic fields corresponding to $r = 4.7$ and $r = 5$, as a function of the ratio $\alpha = V_{e\text{-}h}/V_{e\text{-}e}$. For better illustration Gaussian broadening of peaks has been used. The effective masses ($m_e = 0.068, m_h = 0.340$) and dielectric constant ($\epsilon = 12.5$) are for GaAs, and confining lengths $l_e = l_h = 150$ Å yield confining energies: $\omega_e = 2.49$ meV, $\omega_h = 0.498$ meV, for electrons and for holes. The absorption spectrum for $\alpha = 0$, $r = 5$ in Fig. 1 indeed consists of a series of regularly spaced peaks, while the density of states for $\alpha = 0$, $r = 4.7$ has irregularly spaced peaks leading to a dense irregular spectrum at higher en-

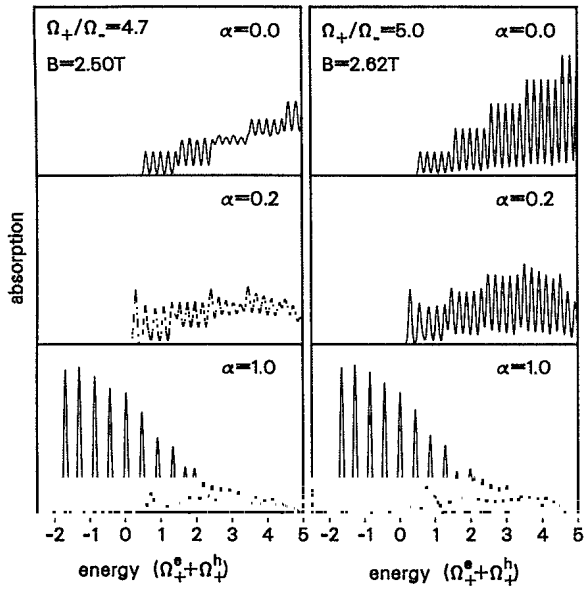

FIG. 1. The magnetoexciton absorption spectra $A(\omega)$ for magnetic fields corresponding to $r = 4.7$ and $r = 5$ as a function of the ratio $\alpha = V_{e\text{-}h}/V_{e\text{-}e}$. The restoration of symmetries by a magnetic field at $r = \Omega_+/\Omega_- = 5$ is visible in the absorption spectrum of X when $\alpha = 0$, and vanishes for $\alpha = 1$.

ergies. The regular structure of the absorption spectrum for $\alpha = 0$, $r = 5$ survives for small electron-electron–hole interaction strength α. For $\alpha = 1$ the quantization is determined by the electron-hole attraction. The spectra for $r = 4.7$ and $r = 5$ are hardly distinguishable. Hence for the weak electron-hole interaction (in comparison with single-particle quantization) one could observe the magnetic field induced level crossings and symmetries of electronic states in quantum dots.

IV. NEGATIVELY CHARGED MAGNETOEXCITONS

Let us now discuss the negatively charged exciton (X^-). We will discuss two cases: (a) excitation of the electron-hole pair in a quantum dot containing an electron and (b) charging of a quantum dot containing an exciton with an additional electron. Let us start with the first case. We assume that prior to the creation of an exciton the initial electron is in the lowest energy state with spin down. After injection of an exciton there are three possible final-state spin configurations, corresponding to linear (L) and two circular ($\sigma_\pm$) light polarizations. In our notation, σ_- generates a pure spin-triplet electronic configuration $|\downarrow_e\downarrow_e\uparrow_h\rangle$; while σ_+ , by injecting an electron with spin up, creates a configuration $|\downarrow_e\uparrow_e\downarrow_h\rangle$, which is a linear combination of spin-singlet and spin-triplet states. The final-state basis for each polarization can be expanded in terms of the following two-electron–one-hole orbitals and spin functions:

$$
\begin{aligned}
|e1,e2,h\rangle_\pm &= [|n_{e1}m_{e1}n_{e2}m_{e2}n_hm_h\rangle \\
&\quad \pm |n_{e2}m_{e2}n_{e1}m_{e1}n_hm_h\rangle]/\sqrt{2}, \\
|e1,e2,h\rangle_+ &= |n_{e1}m_{e1}n_{e1}m_{e1}n_hm_h\rangle \quad (e1 = e2), \\
T_0 &= [|\downarrow_{e1}\uparrow_{e2}\downarrow_h\rangle + |\downarrow_{e2}\uparrow_{e1}\downarrow_h\rangle]/\sqrt{2}, \\
S_0 &= [|\downarrow_{e1}\uparrow_{e2}\downarrow_h\rangle - |\downarrow_{e2}\uparrow_{e1}\downarrow_h\rangle]/\sqrt{2}, \\
T_1 &= [|\downarrow_{e1}\downarrow_{e2}\uparrow_h\rangle]. \qquad (3)
\end{aligned}
$$

Let us denote the expansion coefficients of exact states in the above basis by $C^{f,\sigma_\pm,L}_{n_{e1}m_{e1}n_{e1}m_{e1}n_hm_h}$. The coefficients C are obtained by diagonalizing the three-body Hamiltonian, Eq. (1), in the basis of noninteracting, properly antisymmetrized three-particle states, Eq. (3), for each spin configuration. The two-body electron-electron Coulomb matrix elements are calculated in a similar way to the electron-hole matrix elements, given in Eq. (2). The transition matrix element for the creation of X^- is proportional to the probability amplitude of finding an electron and a hole in the same place with opposite spins, while the second electron has the same spin and position as the initial-state electron. Hence, the transition probability can be written as $\langle f|\mathbf{P}|i\rangle \sim \int\int\int d\mathbf{r}_{e1}d\mathbf{r}_{e2}d\mathbf{r}_h\phi^f(\mathbf{r}_{e1},\mathbf{r}_{e2},\mathbf{r}_h)\delta(\mathbf{r}_{e2}-\mathbf{r}_h)\phi^i(\mathbf{r}_{e1})$; where the initial-state electron occupies the state $|m_0,n_0,\downarrow\rangle$, and spin restrictions related to light polarizations are implicit. Since Coulomb interactions do not mix states with different spin configurations, each final state has purely triplet or singlet configuration. We can identify these spin components (S,T) in the absorption spectrum for each different polarization:

$$
\begin{aligned}
A^{\sigma-}_T(\omega) &= \sum_f \left| \sum_{nm\neq n_0m_0} \eta C^{f,T}_{n_0m_0,nm,nm} \right|^2 \delta(E_f - E_i - \omega), \\
A^{\sigma+}_T(\omega) &= \frac{1}{2}\sum_f \left| \sum_{nm\neq n_0m_0} \eta C^{f,T}_{n_0m_0,nm,nm} \right|^2 \delta(E_f - E_i - \omega), \\
A^{\sigma+}_S(\omega) &= \sum_f \left| C^{f,S}_{n_0m_0,n_0m_0,n_0m_0} + \frac{1}{\sqrt{2}} \sum_{nm\neq n_0m_0} \eta C^{f,S}_{n_0m_0,nm,nm} \right|^2 \delta(E_f - E_i - \omega), \\
A^{L}_T(\omega) &= \frac{3}{4}\sum_f \left| \sum_{nm\neq n_0m_0} \eta C^{f,S}_{n_0m_0,nm,nm} \right|^2 \delta(E_f - E_i - \omega), \\
A^{L}_S(\omega) &= \sum_f \left| C^{f,S}_{n_0m_0,n_0m_0,n_0m_0} + \frac{1}{2} \sum_{nm\neq n_0m_0} \eta C^{f,S}_{n_0m_0,nm,nm} \right|^2 \delta(E_f - E_i - \omega), \qquad (4)
\end{aligned}
$$

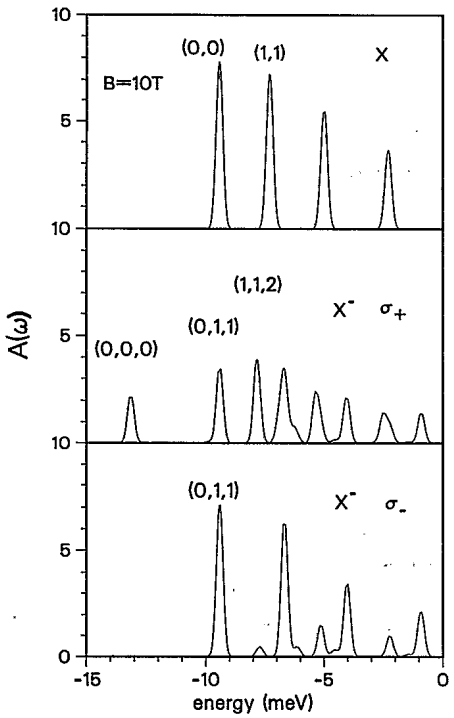

FIG. 2. The low-energy absorption spectra for X and X^- for two different light polarizations $\sigma_\pm$. Peaks have been broadened with Gaussians of width 0.15 meV and labeled according to the lowest Landau level orbitals of noninteracting particles.

where $\eta = -1$ if $n, m < n_0, m_0$, $\eta = 1$ otherwise. The comparison of the absorption spectra of the exciton X and of the negatively charged exciton X^-, for the magnetic field $B = 10$ T, is shown in Fig. 2. The initial-state electron occupies the ground state $|0, 0, \downarrow\rangle$. The two lowest Landau levels (LL) for electrons and 6 LL for holes have been included with 20 states per Landau level, giving a basis of ≈ 4000 states. The energy scale includes the zero-motion energy of excited particles and the energy gap between the valence and conduction bands has been set to zero. Figure 2 represents the low-energy portion of the absorption spectra for $\sigma_\pm$ polarizations, the group of peaks corresponding to excitations of the electron to the first LL. For clarity, we have labeled some of the peaks by single-particle states which significantly contribute to the actual eigenstates of X and X^-, e.g., the $|00, 01, 01\rangle$ state as $(0, 1, 1)$. The $(0, 1, 1)$ denotes an electron in the orbital $m = 0$, an electron in the orbital $m = 1$, and a hole in the orbital $m = 1$, with the net angular momentum $R = 0 + 1 - 1 = 0$. As can be seen from Fig. 2 the structure of peaks is richer for X^- than for X. This is a consequence of a bigger number of final states. In the basis, there are not only states corresponding to one electron in the initial state and the electron-hole pairs excited to the orbitals with opposite angular momentum $|00, 0m, 0m\rangle$, but there are the scattered states: $|0m', 0m'', nm\rangle$ as well. However, as the latter are orthogonal to the initial one-electron state, they give only a small contribution to the absorption spectrum through mixing with allowed states via Coulomb interactions. As a result, we obtain additional small peaks in the higher energy part of the spectrum, e.g., as $(0, 1, 1)$ and $(1, 1, 2)$ doublets for the σ_+ polarization. The difference between the energies of the X^- states for σ_+ and σ_- polarizations is due to the Pauli exclusion principle, which forbids photoexcitation of a second electron into an already occupied state. As a result, the lowest X^- peak, in the σ_- polarization, $(0, 1, 1)$ corresponds to the second $(1, 1)$ peak in the X spectrum. The energy of the X^- is, however, lowered by the exchange energy of two spin-triplet electrons, and is almost identical to the exciton energy (a manifestation of hidden symmetry). The lowest energy peak $(0, 0, 0)$ in the σ_+ polarization corresponds to a singlet electronic state bound to a hole, or an exciton and a captured electron. The difference in the X^- and X energies gives the binding energy of an extra electron $E_b \approx 3.7$ meV. The small, but nevertheless sizable, binding energy is related to broken hidden symmetries as a result of spin, confinement, finite hole mass, and Landau level mixing. Let us now concentrate on the competition of the electron-electron and electron-hole interactions in the negatively charged exciton. This competition is controlled by the ratio α of the electron-hole to the electron-electron interaction. The ratio α can be tuned by spatially separating electrons and holes in the applied perpendicular electric field. In the absence of the electron-hole interactions ($\alpha = 0$), the two-electron motion decouples into the center of mass (c.m.) and relative (rel) motion[9] and the eigenstates of the two-electron Hamiltonian can be defined by c.m. and rel angular momenta: $R_{\text{c.m.}}$ and R_{rel}. $R_{\text{c.m.}}$ takes on any integer values while R_{rel} is even (odd) for the spin-singlet (triplet) configuration. The ground state of the two-electron system is driven by the increasing magnetic field through a series of increasing angular-momentum states: $R_{\text{rel}} = 0, 2, 4, \ldots$. Increasing angular momentum simply means that electrons move away from the center of the dot in order to minimize their mutual repulsion. A similar behavior is observed for a larger number of electrons. The electron-hole interaction mixes these different electronic states and may lead to a completely different ground state (GS). It turns out that for strong magnetic fields (or when Coulomb attraction and repulsion between corresponding orbitals are equal) the GS for zero total angular momentum is built mostly of the state $|00, 00, 00\rangle$ for the electronic spin-singlet (S) configuration and of the state $|00, 01, 01\rangle$ for the electronic spin triplet (T) configuration. The attracting potential of the hole cancels the electron-electron repulsion and all particles tend to sit close to each other in the center of the dot. In Fig. 3 we show the evolution of the X^- energy spectra from that dominated by the electron-electron interactions ($\alpha = 0$) to that with balanced electron-electron and electron-hole interactions ($\alpha = 1$). Figure 3(a) shows the low-angular-momentum energy spectra for two interacting electrons (E_e) in singlet spin configuration. It also shows a free hole (E_h) ($\alpha = 0$) in a magnetic field $B = 10$ T. The electronic GS has the angular momentum $R_e = 4$. The inset on the right hand side shows the three particle joint density of states of the negatively charged exciton: $E_{X^-} = E_e + E_h$. Figure 3(b) represents the formation of

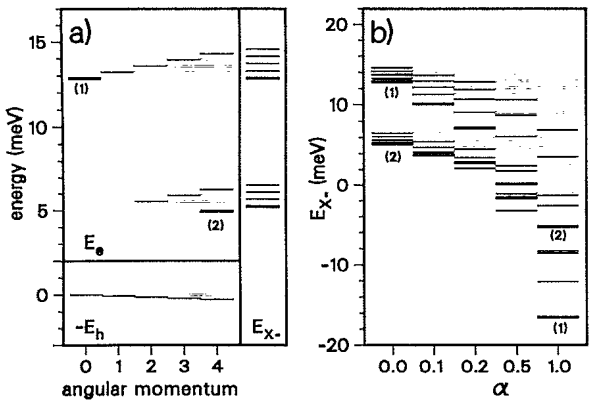

FIG. 3. The low-angular-momentum energy spectra in the lowest Landau level: (a) for two electrons (E_e) in spin singlet configuration, single hole (E_h), and $E_{X^-} = E_e + E_h$, i.e., no electron-hole interaction, zero total angular momentum, electrons are in spin singlet; (b) for X^- of zero total angular momentum, as a function of $\alpha = V_{e\text{-}h}/V_{e\text{-}e}$. Labels (1) and (2) denote levels originating from corresponding noninteracting levels with energies in frame (a).

the GS as the electron-hole Coulomb potential is being switched on. A thick line marks states Φ originating from the ground states for $\alpha = 0$ (Φ_0^α) and for $\alpha = 1$ (Φ_1^α). It can be seen that the energies of these states interchange at approximately $\alpha = 0.5$. For $\alpha = 1$ their overlap with states $|0,4,4\rangle_{\text{c.m.-rel}}$ and $|0,0,0\rangle_{\text{c.m.-rel}}$ is approximately $|\langle\Phi_0^1|0,4,4\rangle_{\text{c.m.-rel}}\rangle|^2 = 0.59$ and $|\langle\Phi_1^1|0,0,0\rangle_{\text{c.m.-rel}}\rangle|^2 = 0.61$. The effect of tuning the electron-hole interaction on the low-energy part of the X^- absorption spectrum is shown in Fig. 4. The large difference between X and X^- GS energies appears, as now the electron-hole attraction can no longer compensate the electron-electron repulsion in X^-. Since the eigenstates of X^- (especially the GS) depend strongly on $V_{e\text{-}h}/V_{e\text{-}e}$ (see Fig. 3), the structure of the peaks for X^- is also different from that given in Fig. 3. In fact, the lowest energy state corresponds to an electronic state with an angular momentum $R_e = 2$, while the second peak corresponds to the two-electron ground state with $R_e = 4$. The two states (1) and (2) marked in Fig. 3 are also indicated in the spectrum. Hence, using the electric field one could map out the two-electron energy spectrum responsible for the energy structure shown in Fig. 3(a). Initially, we examined the situation where an exciton is created in the presence of an electron occupying the zero-angular-momentum state. Since an exciton carries no angular momentum, only zero total angular-momentum states could be probed this way. One can, of course, first inject an exciton and next inject free carriers with finite angular momentum. It is therefore interesting to discuss the interaction between an exciton and a carrier with finite angular momentum R. This procedure allows us to examine in detail the role of hidden symmetries. The hidden symmetry implies the existence in the lowest Landau level of an exact eigenstate of an interacting spin-polarized 2D electron-hole (e-h) fluid. This eigenstate consists of a Bose condensed gas of noninteracting excitons, decoupled from a fluid of excess electrons. If excitons are created in the electron (hole) fluid occupying initially its ground state, the Bose condensed state turns out to be the GS of the entire system. However, if electrons are injected into the exciton gas with a finite angular momentum R, the Bose condensed state may be an excited state. This is illustrated in Fig. 5, which shows the GS energies of X^- vs total angular momentum R. The zero motion energy has not been included. The energy is measured in $\frac{E_0}{\sqrt{2}} = \langle 00,00|V_{e\text{-}e}|00,00\rangle$. Data for noninteracting (solid line) and interacting exciton-electron systems (circles for electronic spin singlet S, and squares for electronic spin-

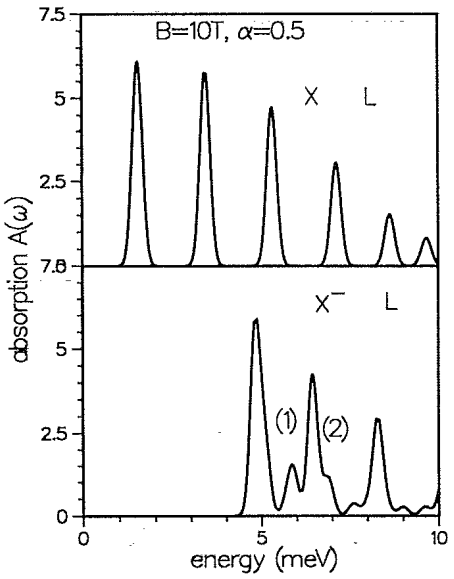

FIG. 4. The low-energy absorption spectra for X and X^- corresponding to energies from 3(b) for $\alpha = 0.5$ (linearly polarized light). Labels (1) and (2) denote levels from 3(b).

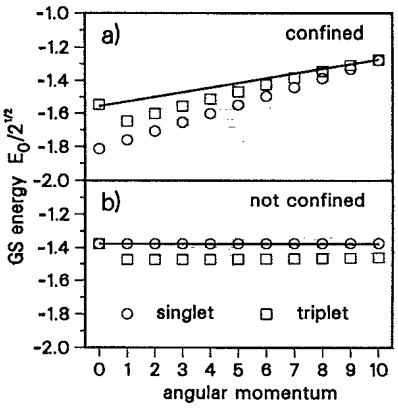

FIG. 5. The ground-state energies of noninteracting (solid line) and interacting (circles for electrons in spin-singlet configuration, squares for spin triplet) exciton-electron system vs total angular momentum: (a) confined in a quantum dot, $B = 10$ T and (b) confined to a plane, lowest LL.

triplet T configuration) have been plotted separately for both laterally confined and free 2D systems. In the latter case, only the lowest LL has been included, which corresponds to the infinite magnetic field limit. As mentioned before, for $R = 0$, the second electron can be bound in X^- in the S configuration only in the presence of confinement. The formation of T-like (spin-polarized electrons) X^- is, however, possible for $R > 0$ (independent of confinement), due to attractive exchange interaction. The Bose condensed state is then not a GS but the first excited state. On the other hand, in the absence of confinement, for S ($R > 0$) and T ($R = 0$) configurations, the exciton does not interact with a carrier. The discrepancy between the $R = 0$ and $R > 0$ cases for T clearly follows from the fact that in the former case the hole is not allowed to occupy the $|00\rangle$ orbital.

V. CONCLUSIONS

We studied the interaction of an electron with an exciton in a quantum dot subjected to magnetic and electric fields perpendicular to the plane of the dot. For balanced electron-hole and electron-electron interactions, the ground state of a negatively charged exciton corresponds to all particles residing in the center of the dot. As a result of the smearing of the charge of the hole, an interaction of the second electron with the charge neutral exciton is fairly weak; the binding energy of the electron by the exciton is small. However, the absorption spectrum for the negatively charged exciton is highly sensitive to the polarization of light, controlling the spins of excited particles. The difference in energies between the X and X^- states can be significantly enlarged by applying an electric field perpendicular to the dot. This leads to spatial separation of electrons and holes and a weakening of the electron-hole attraction, which is neccessary to optically probe different electronic configurations in the three-body complex.

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