SAND2000-0261J

Crossing behavior of the singlet and triplet state of the negatively charged magneto-exciton in a GaAs/AlGaAs quantum well

F. M. Munteanu[*,†], Yongmin Kim[*], C. H. Perry[*,†], D. G. Rickel[*],
J. A. Simmons[‡], and J. L. Reno[‡]
[*]*National High Magnetic Field Laboratory-Los Alamos National Laboratory Los Alamos, NM 87545*
[†]*Department of Physics, Northeastern University, Boston, MA 02115*
[‡]*Sandia National Laboratory, Albuquerque, NM 87185*

(November 22, 1999)

RECEIVED FEB 24 2000 OSTI

Polarized magneto-photoluminescence (MPL) measurements on a high mobility δ-doped GaAs/AlGaAs single quantum well from 0-60T at temperatures between 0.37-2.1K are reported. In addition to the neutral heavy hole magneto-exciton (X^0), the singlet (X_s^-) and triplet (X_t^-) states of the negatively charged magneto-exciton are observed in both polarizations. The energy dispersive and time-resolved MPL data suggest that their development is fundamentally related to the formation of the neutral magneto-exciton. At a magnetic field of 40T the singlet and the triplet states cross as a result of the role played by the higher Landau levels and higher energy subbands in their energetic evolution, confirming theoretical predictions. We also observed the formation of two higher energy peaks. One of them is completely right circularly polarized and its appearance can be considered a result of the electron-hole exchange interaction enhancement with an associated electron g-factor of 3.7 times the bulk value. The other peak completely dominates the MPL spectrum at fields around 30T. Its behavior with magnetic field and temperature indicates that it may be related to previous anomalies observed in the integer and fractional quantum Hall regimes.

INTRODUCTION

The formation and behavior with magnetic field of the negatively charged magneto-excitons (X^-) in modulation doped quantum wells (QWs) and single heterojunctions (SHJs) has been the subject of increased investigation recently[1–7]. The appearance of this particle, which consists of two electrons bound to a hole requires an excess of electrons and is facilitated by the application of a magnetic field. The two bound electrons are indistinguishable and their spin wave function can be either symmetric for the case of the triplet state (X_t^-) or antisymmetric for the case of the singlet state (X_s^-). The evolution of these states with the magnetic field is of a special interest due to the fact that at large fields, the cyclotron diameter becomes smaller than that of the X^-, which can lead to changes in their energy structure.

Experimental investigations[5,7] showed that the energy of the X_t^- and X_s^- exhibit diamagnetic shifts at low fields, in a manner similar to that displayed by the neutral heavy-hole exciton (X^0). Buhmann *et al.*[7] found that the X^0 intensity decreased as the temperature was lowered, whereas the X^- intensity was increased. The energy separation between these lines ($\cong$2meV) proved to be independent of magnetic field and temperature in the range of the magnetic field investigated (0-20T). Recently, Glasberg *et al.*[8] reported that the X_s^- binding energy exhibits a significant increase with increasing field in the range 0-7T and saturates at B=8T, while the X^+ binding energy remains nearly constant between 0-8T.

Theoretical contributions[9–12] have focused mainly on the variation of the energy of the X_t^- and X_s^- states with the magnetic field and there is uniform agreement that at zero magnetic field, only the singlet state of the two dimensional X^- will exist. However, in the high magnetic field regime the calculations become more complicated due to the formation of the Landau levels. The most convenient description under such circumstances is called the lowest-Landau-level approximation (LLL) and it essentially implies that the wave-vector basis space considered is restricted to the lowest subband (n=1) and the lowest Landau level (l=0). In this approximation, there is no bound singlet state in a magnetic field higher than zero. The only bound state will be the triplet one, a situation predicted also by Hund's rule which asserts that the triplet state will generally be at lower energy than the singlet state. The value of the binding energy of the second electron in the triplet state relative to the neutral exciton X^0 was found[12] to be $0.0545e^2/\epsilon\ell_B$ (where ℓ_B is the cyclotron radius). In contrast, experimental investigations showed evidence for the existence of the singlet state at high magnetic fields and the appearance of a triplet state that takes place after the formation of the singlet state[3–5], when the magnetic length becomes smaller than the X^0 diameter. These experimental results show that the higher energy levels play an important role in the energetic structure of the X^-. More complex calculations in which the wave-vector basis set was expanded to include both higher Landau levels and higher QW subbands were performed by Whittaker *et al.*[10]. Their results show that the triplet and the singlet state will cross at a field of about 35T in the case of a narrow QW (100Å), although no such crossing will occur, at least up to 50T, for the case of a large QW (300Å).

In the present work, we have investigated the behavior of the X_t^- and X_s^- states of the negatively charged magneto-excitons formed in a high mobility δ-doped GaAs/AlGaAs single QW with a well width of 200Å. We found that the evolution of these states with magnetic field is fundamentally related to the formation of the neutral exciton. The spectra taken in both right circular (RCP) and left circular (LCP) polarization show that the singlet and triplet states of the X^- cross at a field of about 40T, confirming the predictions of Whittaker *et al.*[10]. The value of the magnetic field where this happens does not depend on the temperature (1.5K or 370mK) or polarization. At low magnetic fields we find that the X_t^- state exhibits a more pronounced diamagnetic shift than the one displayed by the neutral exciton. For both polarizations, the X^0 line disappears at B≅30T and the X_s^- line dominates at very large magnetic fields (beyond 45T), an indication of the important role played by the localization. The difference in the binding energies between the X^0 and X_t^- states presented *maxima* at the filling factors ν=2, 1, 2/3, and 2/5 associated with the integer (IQH) and fractional (FQH) quantum Hall regimes.

Our measurements also revealed the presence of two high energy peaks, labeled H and S. While the S peak is fully RCP polarized and its appearance can be considered the result of the enhancement of the electron-hole exchange interaction (EHEXI)[13–18], the presence of the H peak in both polarizations is more difficult to understand. It appears around filling factor $\nu \approx 2$ ($B \approx 3T$) for a temperature of about 370mK and at $\nu \approx 6$ ($B \approx 1T$) for a temperature of about 1.5K. Its behavior with magnetic field and temperature has certain similarities with the some previously reported anomalies in the integer (IQHE)[19] and fractional quantum Hall regime (FQHE)[20,21].

EXPERIMENT

The remotely (700Å undoped AlGaAs buffer) δ-doped 200Å GaAs/$Al_{0.55}Ga_{0.45}$As single QW used in these studies had a dark electron density of $1.2\times10^{11}cm^{-2}$ and a mobility higher than $3\times10^{6}cm^{2}/Vs$. With constant laser illumination (700μW at 632nm) during the measurements, the 2DEG density increased to $1.58\times10^{11}cm^{-2}$. The experimental layout for the PL measurements has been described previously[22]. Using a quasi-continuous magnet, the magnetic field was varied from 0 to 60 T and then back to 0, in an interval of about 2 seconds. A high-speed CCD camera was used to acquire over 1200 spectra during the complete field swap; the field separation between two consecutive spectra was no larger than 0.165T. A single optical fiber was used to transfer the exciting laser light to the sample as well as for collection of the MPL signal. The fiber had a NA=0.16 and a diameter of 600 μm. A signal to noise ratio of 400-700 was typical for all the spectral data. Right and left circular polarizations were obtained by reversing the direction of the magnetic field. Temperature studies (between 0.37 and 1.5K) were achieved under identical conditions. For the low field studies and for the time-resolved PL measurements, a 20T superconducting magnet was used. In the latter case, the sample was excited using a femto-second Ti-sapphire laser at 810nm with a repetition rate of 76MHz, pumped with an Ar^+ laser.

RESULTS AND DISCUSSIONS

In Fig. 1 we show the unpolarized spectra taken at very low magnetic field (0-1.4T) at a temperature T=2.1K. When a magnetic field was applied, the peak located at 1.526 eV at B=0T initially moved towards *lower* energies. Bauer[23] showed that at low temperatures and zero magnetic field, the excitonic states dominate the spectrum if the carrier density is lower than $n=1.0\times10^{11}cm^{-2}$, otherwise the PL signal will be determined by the recombination of the two-dimensional electron gas (2DEG) with the photoexcited holes. In the latter case, the application of a magnetic field was found to recover the excitonic states. In our sample, the measured (from transport studies under equivalent illumination) carrier density of $n=1.58\times10^{11}cm^{-2}$ is slightly above the estimated density required for the unbinding of the exciton ($n=1.0\times10^{11}cm^{-2}$). At T=2.1K, the main peak, located at 1526 meV in the spectrum taken at B=0T, is due to the recombination of the 2DEG with the photoexcited holes. The application of a small magnetic field initially caused the peak to move towards lower energies, as a result of a gradual evolution of the band-to-band recombination into the excitonic state (shown by the X^0 symbol in Fig. 1) with a strong Coulomb energy. This type of behavior was also observed at higher temperatures (e.g. 4.2K), but was absent at temperatures of 1.5K and lower. The inset in Fig. 1 shows the energy evolution of this peak with the magnetic field. It is well known[24] that the energy of a purely 2D excitonic peak typically follows a B^2 dependence in the limit of low magnetic fields. Such behavior is not anticipated in our case at temperatures higher than 2K, owing to the presence of the 2DEG that will screen the interaction between the hole and the electron. When the magnetic field is increased, this screening becomes smaller due to the localization of the electrons, such that the excitonic binding becomes stronger.

For the spectra recorded at temperatures of 1.5K and lower, the variation of the energy with magnetic field shows the predicted diamagnetic shift. This indicates that the excitonic character is achieved as soon as the magnetic field is turned on. At these temperatures, a second peak due to the light hole excitonic (LHE) state forms at slightly higher energy, around B=0.5T. The difference in energy between LHE and heavy hole X^0 states at very low magnetic field, of about 1.9meV, is comparable with those found in other studies of QWs of about the same size[25]. These two peaks cross at a magnetic field of about 8.5T. Beyond 20T the LHE peak could no longer be resolved.

In Fig. 2 we show the results obtained at a temperature of 1.5K in both RCP (σ+) and LCP (σ-) polarization for

DISCLAIMER

two different magnetic fields (6T and 16T). The spectra at B=0T (not shown) are identical for both polarizations and have a full width at half maximum (FWHM) of about 0.66meV which reflects the high quality of the sample. For the moment we shall postpone any discussion of the higher energy peaks labeled H (in Fig. 2a, b) and S (in Fig. 2b) and focus our attention on the lower energy peaks.

At B=6T, in both polarizations, three peaks can be distinguished. The middle peak derived directly from the 2DEG recombinations at B=0T is ascribed to the neutral excitons X^0 which, due to the wide buffer (700Å), will be non-localized. The lower energy peak is generated by the triplet state of the negatively charged exciton X_t^-. In order to understand this assignment, we note that the lowest energy spin wave-functions for the X_t^- state are:

In σ- (LCP) polarization:

$$\psi^t_{LCP}(-1/2) = e\uparrow e\uparrow h\downarrow \text{ and}$$

$$\psi^t_{LCP}(-3/2) = 1/\sqrt{2}(e\uparrow e\downarrow + e\downarrow e\uparrow)h\downarrow$$

In σ+ (RCP) polarization:

$$\psi^t_{RCP}(+1/2) = e\downarrow e\downarrow h\uparrow \text{ and}$$

$$\psi^t_{RCP}(+3/2) = 1/\sqrt{2}(e\uparrow e\downarrow + e\downarrow e\uparrow)h\uparrow$$

For the σ- (LCP) polarization, the formation of the triplet state requires the presence of either two spin-up electrons or a spin-up and a spin-down electron. After recombination of one electron from the complex with the hole, the state $\psi^t_{LCP}(-1/2)$ will leave behind a spin +1/2 electron, while the state $\psi^t_{LCP}(-3/2)$ will leave behind a spin -1/2 electron. The change in the Zeeman energy for both recombinations will be the same:

$$\Delta E^t_{Z,LCP} = 1/2(|g_h| - |g_e|)\mu_B B = +1/2\Delta E_Z(X^0),$$

where $\Delta E_Z(X^0)$ is the Zeeman splitting of the neutral exciton[26]. For this reason, the peak seen in the LCP recombination will contain contributions from both states.

For the σ+ (RCP) polarization, the formation of a charged exciton requires the existence of a spin-up electron bound to a spin-down one and we can safely ignore (at least for $\nu < 2$) the $\psi^t_{RCP}(+1/2)$ state, which requires the presence of two spin-down electrons. This means that we have to expect, for filling factors $\nu < 2$, a higher population of X_t^- in the LCP polarization than in RCP polarization coming mainly from the $\psi^t_{LCP}(-1/2) = e\uparrow e\uparrow h\downarrow$ state. This happens with the simultaneous decrease in the population of the neutral exciton X^0, as can be seen in Fig. 2a. The decay of the $\psi^t_{RCP}(+3/2)$ state will leave behind a spin +1/2 electron, and the change in the Zeeman energy will be:

$$\Delta E^t_{Z,RCP} = -1/2\Delta E_Z(X^0).$$

The spectra at B=16T presented in Fig. 2b clearly show the additional presence of the singlet state X_s^- of the negatively charged exciton in both polarizations. The spin wave-functions that correspond to this state are:

In σ- (LCP) polarization:

$$\psi^s_{LCP}(-3/2) = 1/\sqrt{2}(e\uparrow e\downarrow - e\downarrow e\uparrow)h\downarrow$$

This state will recombine leaving behind a spin -1/2 electron. The change in the Zeeman energy will be:

$$\Delta E^s_{Z,LCP} = +1/2\Delta E_Z(X^0).$$

In σ+ (RCP) polarization:

$$\psi^s_{RCP}(+3/2) = 1/\sqrt{2}(e\uparrow e\downarrow - e\downarrow e\uparrow)h\uparrow.$$

This state will recombine leaving behind a spin +1/2 electron. The variation of the Zeeman energy in this case will be:

$$\Delta E^s_{Z,RCP} = -1/2\Delta E_Z(X^0).$$

The formation of the X_t^- state will be favored over the X_s^- state in the σ- (LCP) polarization, as it requires the presence of two spin-up electrons. At filling factors less than 2, the number of spin-up electrons will be larger that the number of the spin-down ones and this can explain why the X_t^- state appears at a lower field than the X_s^- state in the σ- (LCP) polarization. This argument is not valid for the σ+ (RCP) polarization, as the singlet and triplet states require the presence of the same type of particles. An alternative explanation for why the formation of the X_t^- state is favored over the formation of the X_s^- state in both polarizations could be due to a non-resonant excitation energy condition.

In Fig. 2b it can be seen that the intensities of all peaks in the RCP polarization (with the exception of the S peak) are lower than those observed in the LCP polarization, due to the depopulation of the spin-down electronic level with increasing field. The fact that we still see signal in the RCP polarization at very low temperatures can be explained by the fact that the energy of the exciting photons (632nm, 1.96eV) is much higher than the band gap plus the first Landau level, such that the photocreated electrons and holes may be excited into upper Landau levels.

The spectra taken at T=370mK show the same characteristics as those taken at T=1.5K except that the intensity of the charged excitons is higher at low temperatures, a further indication that the localization plays an important part in their formation. A similar result has been reported previously by Buhmann *et al.*[7].

Fig. 3 shows the LCP spectra at some representative magnetic fields between 20T and 60T, for T=1.5K. All the spectra have been normalized with respect the zero field data. It can be seen that with increasing magnetic field, the X_s^- state gains in intensity compared with the X_t^- state. Also, the neutral exciton loses its intensity compared with the charged excitons and beyond 30T it is almost undistinguishable from the H peak. The intensity loss could be the result of increased localization of the electrons induced by the magnetic field which causes more excitons to bind to an additional electron in order to reduce the energy of the system. The major observation is that the singlet and triplet states merge together (and possibly cross) at a magnetic field of about 40T. This behavior is shown more graphically in the Fig. 4, where it can be seen that X^0 and X_t^- show a diamagnetic shift at low magnetic fields but a roughly linear evolution at high magnetic fields. The larger diamagnetic shift experienced by X_t^- state compared with X^0 state is related to the weaker binding energy of the outer electron in the triplet state compared to the binding energy of the exciton. At very high magnetic fields (beyond 45T), the negatively charged exciton peak (X_s^-) dominates, showing the role played by the localization of the electrons in its formation. It is important to mention that the spectra taken at T=370mK show identical characteristics, with a merging of the singlet and triplet states at the same magnetic field of about 40T. This is an indication that the crossing behavior is strongly related to the Coulomb interaction[10], and is not affected by the fluctuations in the spin populations induced by thermal activation.

Fig. 5 shows the energy separation between X^0 and X_s^- peaks, $\Delta E(X^0\text{-}X_s^-)$, X^0 and X_t^-, $\Delta E(X^0\text{-}X_t^-)$, and H and X^0 peaks, $\Delta E(\text{H-}X^0)$, respectively. As we cannot follow the X^0 peak to fields higher than 32T, we have linearly extrapolated the data up to 60T. It can be seen that in the X_t^- case, $\Delta E(X^0\text{-}X_t^-)$ increases with field from 0.4 meV at 2 T to 1.2 meV at 30T; a similar behavior has been reported by others[5]. However, the energy separation between the X_s^- and X^0 states $\Delta E(X^0\text{-}X_s^-)$ decreases slowly from 2.1 meV at B=15.7T to 1.4 meV at B=58T. In the low field regime, it has been shown[5,8] that the X_s^- binding energy should increase due to the reduction of the orbits of the two electrons with magnetic field. At high magnetic fields, one might expect that the changes in the size of the orbits for both X^0 and X_s^- states to be almost the same. However, we note that the spatial wave function for the X_s^- state is symmetric, while for the X_t^- state is antisymmetric. Consequently, at high magnetic fields the orbit of the outer electron of the X_t^- state will be more affected and its binding energy will increase compared with the of X^0 and X_s^- states.

In Fig. 5 we wish to point out the oscillations in the energy separation between X^0 and X_s^- peaks that, to our knowledge, have not been reported before. $\Delta E(X^0\text{-}X_t^-)$ displays *maxima* at the filling factors ν=2, 1, 2/3 and 2/5. We associate these changes in the energy separation with the localization of the electrons in the integer and fractional quantum Hall regimes. This causes a stronger response in the X_t^- case, leading to an enhancement in its binding energy compared with the X^0 state. In contrast, the difference in the energy between H and X^0 peaks, $\Delta E(\text{H-}X^0)$, shows *minima* at the filling factors ν=1 and 2.

The experimental data that we obtained are in good agreement with the theoretical prediction[10] that a crossing between the triplet and singlet state should occur at a field around 35T in the case of a 100A QW. This behavior comes entirely from the Coulomb interaction. As calculated before, the Zeeman energy must play the same role in the measured energy of the neutral exciton and in the X_s^- and X_t^- state that is, in the LCP polarization the measured energies of the three species (X_t^-, X_s^- and X^0) will be raised by $1/2\Delta E_Z(X^0)$, while in the RCP polarization they will be lowered by the same amount due to the Zeeman interaction.

Fig. 6 shows the evolution of the intensities for the H and X^0 peaks. It can be seen that the excitonic line shows distinctive minima at ν=1, 2 and 6. A similar observation was reported by Turberfield *et al.*[20] for ν=1 as well as for some fractional filling states. They related such behavior to the expected suppression of screening at these filling factors, which will determine an increase in the excitonic binding energy and, consequently, a lower decay rate. The inset in Fig. 6 shows the intensity behavior of the X_s^- and X_t^- states. At the filling factors ν=1 and 2, the intensity of the X_t^- state shows the same minima as in the case of the neutral exciton, indicating that their formation is interrelated. The occurrence of the intensity minima at magnetic fields a little lower than the exact filling factors has been suggested[20] to be an indication of the localization effects. At B$\approx 40T$, the triplet and the singlet states can no longer be distinguished. The rapid increase in the intensity of the X_s^- state beyond this magnetic field is possibly due to an admixture of these two states. We may also expect that at these fields, the X_s^- states would be dominant as the X_t^- is an excited state.

In contrast, the H line shows the opposite behavior at these filling factors as it displays an increase in intensity at ν=1,2 and 6. In Fig. 7a we present the spectra recorded at B=30T at T=1.5K and T=370mK. In Fig. 7b we show the evolution of the ratio between the intensities of the H and X^0 peaks with magnetic fields at these two temperatures. It can be seen that the intensity of the H peak at low temperatures increases compared with the neutral exciton line, ruling out the possibility of it being generated by the changes in the population of electrons and photo-excited holes. This tendency is observed in both σ+ and σ- polarization. A comparable anomaly present in the FQHE regime has also been observed[20,21]. Recently, we reported[19] the formation of a similar H line derived from the X^0 peak in the case of the MPL spectra obtained from a $GaAs/Al_{0.3}Ga_{0.7}As$ SHJ.

Up to this point, we are unable to find a satisfactory explanation for the appearance of this peak. Its behavior does not strongly depend on polarization, as is the case for the heavy-hole exciton X^0 peak whose intensity is reduced in the RCP spectra compared to the LCP spectra. Because the sample was grown with no interruptions, we have ruled out that its formation could be related to imperfections in the size of the QW[27,28]. Deveaud *et al.*[29] showed that a single monolayer disorder can be observed in the PL spectra of GaAs/AlGaAs interfaces. They found that the smallest splitting between the peaks was 4.8meV in the case of the signal coming from adjacent regions differing in their widths by a/2 (where a is the GaAs lattice constant). This value is about one order of magnitude larger than the splitting that we observe between the H and X^0 peaks. We also have to rule out the possibility of it being created as a result of a field induced heavy hole exciton (X^0) spin-splitting, because the energy separation between the H and the X^0 peaks is almost constant over the entire range of magnetic field considered ($\Delta E \approx 0.5meV$). This energy separation, as well as the fact that the H peak appears in the integer quantum Hall regime also prevents its association with the magtneto-roton minima formation predicted by Girvin *et al.*[30]. The energy separation is predicted to scale with magnetic field as $B^{1/2}$, which would lead to an increase in the energy separation between the H and X^0 peaks with magnet field.

The results of the time-resolved PL spectroscopy measurements between 0 and 18T showed (inset of Fig. 8) a very long rise time of the PL intensity (about 1.5 ns) for X^0 and X_t^- states for a QW structure. For the neutral excitons formed in a GaAs/AlGaAs SHJ, the same phenomenon was reported by Shen *et al.*[31]. They proposed a vertical transport model to explain this result. Such a model is clearly not applicable in the case of a QW, as the size of the well limits the vertical displacement of the electrons. Damen *et al.*[32] showed that lowering the temperature will cause an increase in the rise time of the neutral exciton. This is due to the longer time required for the initial distribution of large k-vector excitons to relax towards the thermalized distribution centered on k=0. At low temperatures this will take place mainly via exciton-exciton collisions, a process that is slowed down in this regime. On the other hand, a large population of excitons with large k could play an important role in the formation of the X_t^- state over of the X_s^- state at low magnetic fields.

The life time of both the X_t^- and X^0 states (Fig. 8) increases as a function of magnetic field. Citrin[33] estimated a lifetime of about 100ps in the case of localized excitons in a GaAs/AlGaAs QW with a well-width between 50-150A. Here, the measured excitonic life-times are considerably longer (1.8ns at B=0T), supporting the hypothesis that the excitations are non-localized. Our values are comparable with those found by Ron *et al.*[34] and Feldman *et al.*[35] and they slowly increase with magnetic field. In the case of the X_t^- state, its life-time is slightly larger than that of the neutral exciton. Its behavior with magnetic field is similar to that of the X^0 state, reconfirming our supposition that its formation is indeed closely related to that of the neutral exciton.

In Fig. 2b it can be seen that a high energy peak, labeled S is present in the $\sigma+$ (RCP) polarization, but is absent in the $\sigma-$ (LCP) polarization. Its formation takes place around 9T and we could follow it up to a field of about 25T. We interpret the presence of this peak in the spectra as a result of the enhancement of the electron-hole exchange interaction. Chen et al.[13] showed that the short range of the EHEXI in a QW is considerably enhanced as a result of the confinement, compared with its value in the bulk material. For this reason, the fine structure of the HHE at B=0T consists of two transitions. The higher energy one is dipole allowed and the lower energy one (with the energy equal to the energy of the HHE in the absence of the EHEXI) is dipole forbidden. Labeling the excitonic states $|m_h, m_e>$, the dipole allowed transitions will be[13]: $|\pm 3/2, \mp 1/2>$ and $|\pm 1/2, \pm 1/2>$ for the HHE and LHE respectively. The $|+3/2, -1/2>$ and $|+1/2, +1/2>$ states should be RCP polarized, while the $|-3/2, +1/2>$ and $|-1/2, -1/2>$ are LCP polarized. The dipole forbidden transition will be given by $|\pm 3/2, \pm 1/2>$. The appearance of this peak in the PL signal[14] was related to thermalization effects. The selection rules show that for dipole forbidden states to decay, an electron- or hole-spin flip is required[16]. Bauer *et al.*[14] observed that the energy difference between the dipole allowed peak and the dipole forbidden peak (whose energy is the same of the energy of the heavy hole exciton X^0 calculated without EHEXI) is about 0.6meV, for a QW with the well width of 150Å at zero magnetic field. These results were questioned in the theoretical work presented by Andreani and Bassani[17]. The latter calculated the factor with which the short range exchange interaction is enhanced in QWs over the bulk value as a function of QW width. Their results, using a value of the bulk exchange-splitting measured earlier[36], showed that the energy splitting should be at least ten times smaller than the experimentally observed value. In the case of a 200Å QW for example, at zero magnetic field, this splitting should be about 0.04meV. Our data, when extrapolated to B=0T , showed that the energy separation between the S and the X^0 peaks is very close to zero. For this reason, we could not resolve the S line at a lower magnetic field. In Fig. 4 we show the evolution with magnetic field of all the peaks discussed before (X_s^-, X_t^-, X^0, H, LHE and S). It can be seen that the LHE and the heavy hole exciton X^0 peaks cross around 8.5T due to different in-plane masses of the light and heavy holes. Also, around B=25T, the intensity of the S peak becomes very small as a result of depopulation of the $-1/2$ electronic level.The energy separation between this line and the X^0 peak increases almost linearly with magnetic field. From these data, as well as from the fact that the initial wave function for these two peaks are $|+3/2, +1/2>$ and $|+3/2, -1/2>$, we calculate a g-factor for the electron which is 3.7 times higher than the bulk value. Such an enhancement is explained by the enhancement of the EHEXI as demonstrated by Nicholas *et al.*[37], and by Lefebvre *et al.*[38].

CONCLUSIONS

In conclusion, polarized PL measurements were performed on a high quality δ-doped GaAs/$Al_{0.55}Ga_{0.35}As$

single QW subjected to magnetic fields up to 60T. The appearance of the negatively charged magneto-exciton states (X_s^- and X_t^-) was found to be directly related to the formation of the non-localized neutral exciton state (X^0). At a magnetic field of about 40T, the X_s^- and X_t^- states crossed, confirming the important role of the higher Landau levels and higher subbands in their energy behavior. The difference in the binding energies between the X^0 and X_t^- states presented maxima at the filling factors ν=2, 1, 2/3, and 2/5 associated with the integer and fractional quantum Hall regimes. A stronger diamagnetic shift in the case of the X_t^- state than in the case of the X^0 state was considered as evidence of a different spatial extension of these states.

We also analyzed the appearance of a $\sigma+$ polarized high-energy peak (S), which we interpret as evidence for the enhancement of the enhanced EHEXI in agreement with previous theoretical work. From our data we estimate that the effective g-factor of the electron is 3.7 times the bulk value. Of equal interest is the appearance of an additional peak (H), with an energy between that of the X^0 and S peaks. Its intensity is relatively insensitive to polarization but it becomes the dominant peak beyond 30T ($\nu \approx 1/5$). It shows intensity oscillations in the IQHE and FQHE regimes that are correlated with the intensity oscillations of the X^0 peak, while the energy separation between these two lines remains almost constant. Up to now, we have no satisfactory explanation for the appearance of this peak.

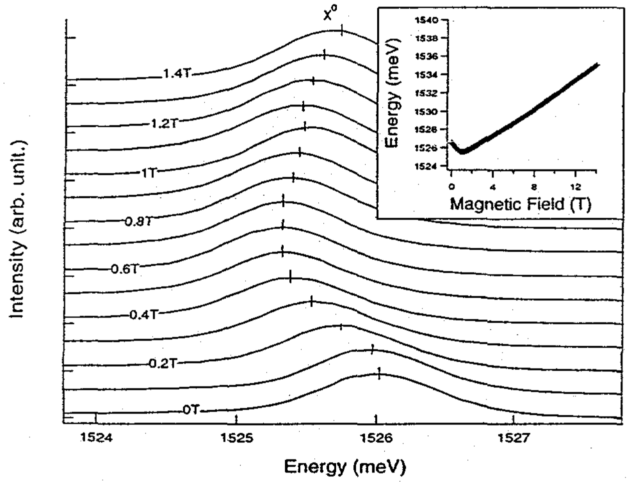
AKNOWLEDGEMENTS

The authors gratefully acknowledge the engineers and technicians at NHMFL-LANL in the operation of the 60T QC magnet. Work at NHMFL-LANL is supported by NSF Cooperative Agreement # DMR-9527035, the Department of Energy and the State of Florida. Work at Sandia National Laboratory is supported by the Department of Energy.

---

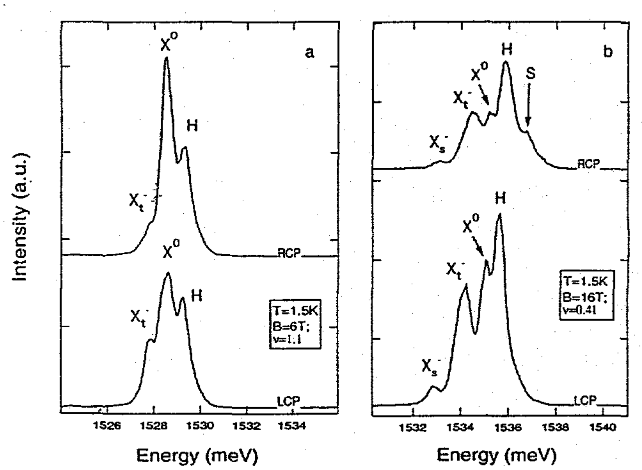
[1] K.Keng, R. T. Cox, Y. M. d'Aubigne, F. Bassani, K. Saminadayar and S. Tatarenko, Phys. Rev. Lett. **71**, 1752 (1993)
[2] G. Finkelstein, H. Shtrikman and I. Bar-Joseph, Phys. Rev. Lett.**74**, 976 (1995)
[3] G. Finkelstein, H. Shtrikman and I. Bar-Joseph, Phys. Rev. B**53**, R1709 (1996)
[4] A. J. Shields, J. L. Osborne, M. Y. Simmons, M. Pepper and D. A. Ritchie, Phys. Rev. B**52**, R5523 (1995)
[5] A. J. Shields, M. Pepper, M. Y. Simmons and D. A. Ritchie, Phys. Rev. B**52**, 7841 (1995)
[6] A. J. Shields, M. Pepper, D. A. Ritchie, M. Y. Simmons and G. A. C. Jones, Phys. Rev. B**51**, 18049 (1995)
[7] H. Buhmann, L. Mansouri, J. Wang, P. H. Benton, N. Mori, L. Eaves, M. Henini and M Potemski, Phys. Rev. B**51**, 7969(1995)
[8] S. Glasberg, G. Finkelstein, H. Shtrikman and I. Bar-Joseph, Phys. Rev. B**59**, R10425 (1999)
[9] J. R. Chapman, N. F. Johnson and V. N. Nicopoulos, Phys. Rev. B**55**, R10221 (1997)
[10] D. M. Whittaker and A. J. Shields, Phys. Rev. B**56**, 15185 (1997)
[11] J. R. Chapman , N. F. Johnson and V. N. Nicopoulos, Phys. Rev. B**57**, 1762 (1998)
[12] J. J. Palacios, D. Yoshioka and A. H. MacDonald, Phys. Rev. B**54**, R2296 (1996)
[13] Y. Chen, B. Gil, P. Lefebvre and H. Mathieu, Phys. Rev. B**37**, 6429 (1988)
[14] R. Bauer, D. Bimberg, J. Christed, D.Oertel, D.Mars, J.N. Miller, T.Fukunaga and H. Nakashima in Proceedings of the 18th International Conference on Semiconductor Physics, Stockholm, 1986, edited by O. Engstrom (World-Scientific, Singapore, 1987), p.525
[15] M. Potemski, J.C. Maan, A. Fasolino, K. Ploog and G. Weimann, Surface Sci.**229**, 151-154 (1990).
[16] H.V. van Kestern, E.C. Cosman, F.J.A.M. Greidanus, P. Dawson, K.J. Moore and C.T. Foxon, Phys. Rev. Lett.**61**, 129 (1988).
[17] L.C. Andreani and F. Bassani, Phys. Rev. B**41**, 7536 (1990)
[18] B.R. Salmassi and G.E.W. Bauer, Phys. Rev. B**39**, 1970 (1989).
[19] F. M. Munteanu, Y. Kim, C. H. Perry, D. G. Rickel, J. A. Simmons and J. L. Reno, Solid State Commun., **111**, 665 (1999).
[20] A.J. Turberfield, S. R. Haynes, P. A. Wright, R. A. Ford, R. G. Clark and J. F. Ryan, Phys. Rev. Lett.**65**, 637 (1990)
[21] D. Heiman, B. B. Goldberg, A. Pinczuk, C. W. Tu, A. C. Gossard and J. H. English, Phys. Rev. Lett.**61**, 605 (1988)
[22] C. H. Perry, Y. Kim and D. G. Rickel, Physica B**246-247**, 182 (1998)
[23] G. E. W. Bauer, Phys. Rev. B**45**, 9153 (1992)
[24] A. H. MacDonald, and D. S. Ritchie, Phys. Rev. B**33**, 8336 (1986)
[25] A.J. Shields, J.L. Osborne, D.M. Whittaker, M.Y. Simmons, M. Pepper and D.A. Ritchie Phys. Rev. B**55**, 1318, (1997); W.T. Masselink, P.J. Pearah, J. Klem, C.K. Peng and H. Morkoc, Phys. Rev. B**32**, 8027 (1985)
[26] M. J. Snelling, E. Blackwood, C. J. McDonagh, R. T. Harley and C. T. B. Foxon, Phys. Rev. B**45**, 3922 (1992)
[27] C.W. Tu, R.C. Miller, B.A. Wilson, P.M. Petroff, T.D. Harris, R.F. Kopf, S.K. Sputz and M.G. Lamont, Journal of Crystal Growth**81**, 159-163 (1987);
[28] R.C. Miller, C.W. Tu, S.K. Sputz and R.F. Kopf, Appl. Phys. Lett.**49**, 1245 (1986)
[29] B. Deveaud, J.Y. Emery, A. Chomette, B. Lambert and M. Baudet, Appl. Phys. Lett.,**45 (10)**, 1078 (1984)
[30] S.M. Girvin, A.H. MacDonald and P.M. Platzman, Phys. Rev. Lett. 54, 581 (1985) and Phys. Rev. B 33, 2481 (1986)
[31] J. X. Shen, Y. Oka, C. Y. Hu, W. Ossau, G. Landwehr, K. -J. Friedland, R. Hey, K. Ploog and G. Weimann, Phys. Rev. B**59**, 8093 (1999)
[32] T.C. Damen, J. Shah, D. Y. Oberli, D. S. Chemla, J. E. Cunningham and J. M. Kuo, Phys. Rev. B**42**, 7434 (1990)

[33] D.S. Citrin, Phys. Rev. B47, 3832 (1993) .
[34] A.Ron, H. W. Yoon, M. D. Sturge, A. Manassen, E. Cohen and L. N. Pfeiffer, Solid State Commun.97, 741 (1996)
[35] J. Feldman, G. Peter, E. O. Gobel, P. Dawson, K. Moore, C. Foxon and R. J. Elliott, Phys. Rev. Lett.59, 2337 (1987)
[36] W. Ekardt, K. Losch and D. Bimberg, Phys. Rev. B 20, 3303 (1979); R. Sooryakumar and P.E. Simmonds, Phys. Rev. B27, 4978 (1983)
[37] R.J. Nicholas, R.J. Haug, K.v. Klitzing and G. Weiman, Phys. Rev. B37, 1294 (1988)
[38] P. Lefebvre, B. Gil, J.P. Lascaray, H. Mathieu, D. Bimberg, T. Fukunaga and H. Nakashima, Phys. Rev. B37, 4171 (1988)



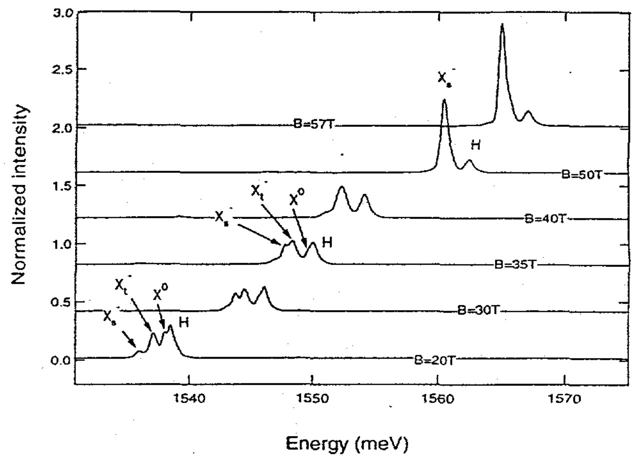
Resub PRB Munteanu et al. Fig1

FIG. 1. The unpolarized spectra taken at low magnetic fields at T=2.1K. The inset shows the variation of the energy as a function of magnetic field. The peak at 1526meV initially displays a small shift towards lower energies when the magnetic field is applied.



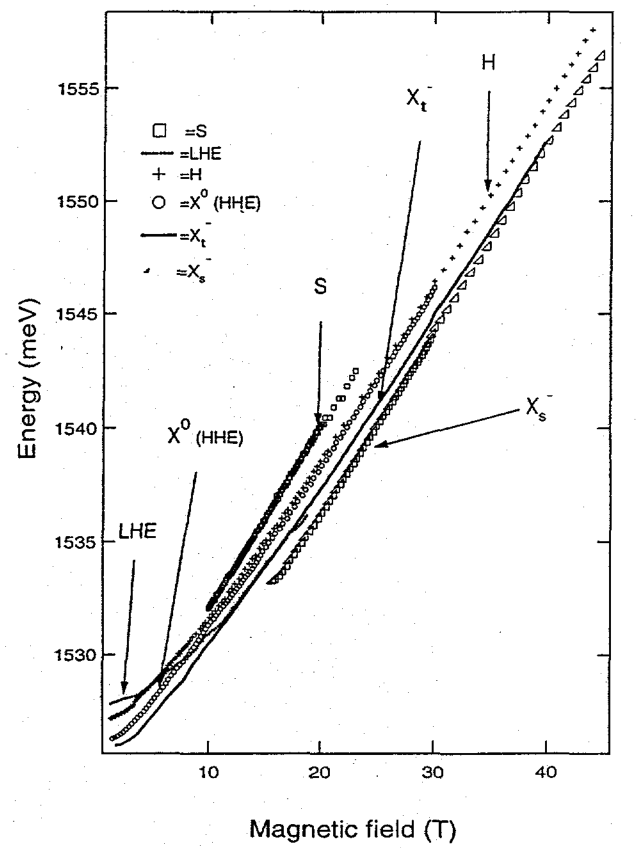
Resub PRB Munteanu et al. Fig2

FIG. 2. The σ+ (RCP) and σ- (LCP) polarized spectra at B=6T (a) and B=16T (b) taken at a temperature of 1.5K. The intensities are all normalized to the B=0T field data for the X^0 .



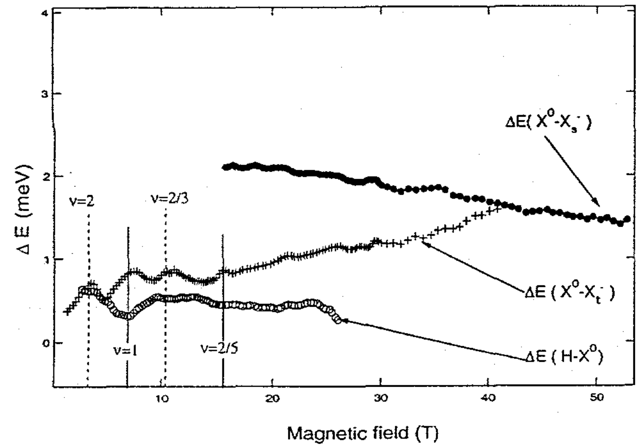
Resub PRB Munteanu et al. Fig3

FIG. 3. The LCP polarized spectra taken at a temperature of 1.5K. All the spectra were normalized with respect to the B=0T field spectrum.



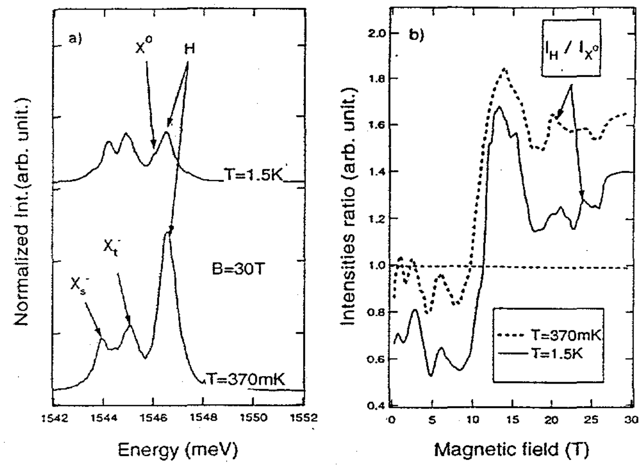
Resub PRB Munteanu et al. Fig4

FIG. 4. The energy evolution of the energy of the X_s^-, X_t^-, X^0, H, LHE and S peaks with magnetic field. For clarity, the data is shown only to 45T. Beyond this field the H and X_s^- peaks evolve linearly.



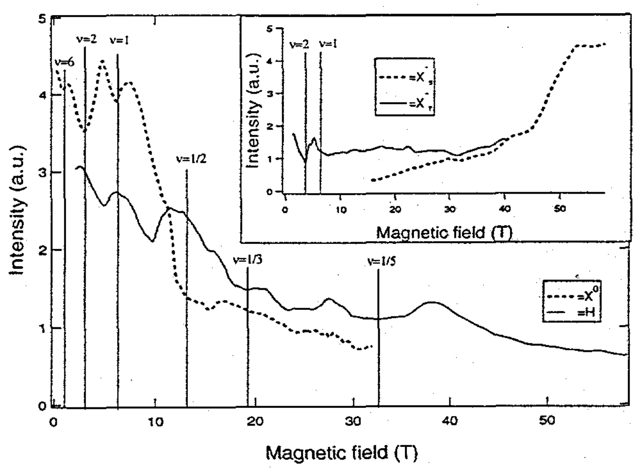
Resub PRB Munteanu et al. Fig5

FIG. 5. The energy separation between X^0 and X_s^- peaks, $\Delta E(X^0\text{-}X_s^-)$, between X^0 and X_t^-, $\Delta E(X^0\text{-}X_t^-)$ and the H and X^0 peaks, $\Delta E(H\text{-}X^0)$ are shown as a function of magnetic field. The data show oscillations at certain filling factors marked on the figure.



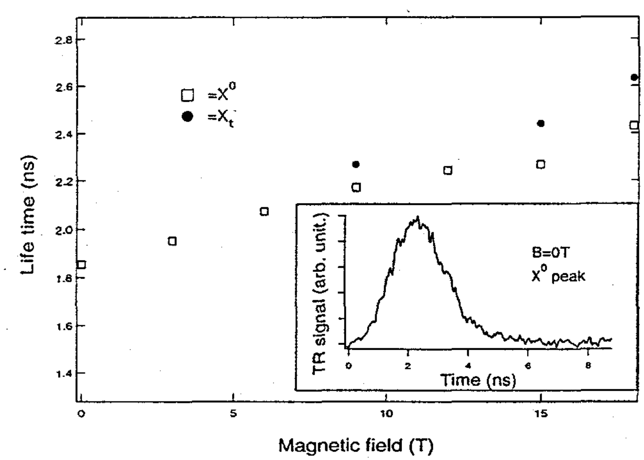
Resub PRB Munteanu et al. Fig7

FIG. 7. (a) The B=30T spectra at two temperatures 1.5K and 370mK. It should be noted that the H peak is enhanced at lower temperatures compared with the other peaks. (b) The ratio of the H peak and heavy-hole exciton X^0 intensities at two temperatures (T=1.5K and T=370mK) with magnetic field. Beyond 30T the X^0 peak could no longer be resolved.



Resub PRB Munteanu et al. Fig6

FIG. 6. The evolution of the X^0 and H peak intensities with the magnetic field. The inset shows the evolution of the X_s^- and X_t^- peaks with magnetic field. Note the interrelation of the intensities between X^0 and X_t^- and between X^0 and H as a function of filling factor.



Resub PRB Munteanu et al. Fig8

FIG. 8. The evolution of the life time of the X^0 and X_t^- peaks up to 18T. The inset shows the shape of the time-resolved (TR) signal obtained from the X^0 peak at B=0T. Of note is the long rise-time of the TR signal.