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Effect of a magnetic field on the excitonic luminescence line shape in a quantum well

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The effect of magnetic field on the excitonic photoluminescence line shape has been studied in a high-quality single GaAs-$Al_xGa_{1-x}As$ quantum well grown by the molecular-beam-epitaxy technique. An increase of magnetic field from 0 to 6 T has been found to result in (1) a decrease in the Lorentzian contribution Γ_0 to the line shape from $\Gamma_0(0\ \text{T})=0.504\pm0.01$ meV to $\Gamma_0(6\ \text{T})=0.336\pm0.01$ meV due to the formation of a quasi-zero-dimensional density of states. This leads, in turn, to an increase in the exciton dephasing time due to the inhibition of the carrier relaxation, and (2) an increase in the Gaussian contribution from $\sigma(0\ \text{T})=0.24$ meV to $\sigma(6\ \text{T})=0.39$ meV, attributed to the shrinking of the exciton wave function in real space; the last effect causing the exciton to become more responsive to the statistical potential fluctuations at the quantum-well interfaces.

I. INTRODUCTION

Excitonic photoluminescence (PL) spectra result from the formation, under optical excitation, of free electrons and holes which subsequently form excitons which, in their turn, emit a well-defined radiation energy upon their collapse. In a homogeneous medium at absolute zero temperature, the distribution function of electrons and holes is reestablished by a rearrangement of carriers within a certain nonzero scattering time τ (final-state interaction),[1] resulting in a Lorentzian broadening of the emission line with a broadening parameter $\Gamma_0=\hbar/\tau$. The resulting homogeneous linewidth Γ_0 is essentially the rate at which a newly created state "forgets" the phase with which it was created.

If excitons exist in a disordered medium, the energy radiated by each exciton will be influenced by the spatial position of the exciton. As an example of a disordered medium, we can consider the interface of a quantum-well (QW) structure, which has been found to consist of randomly distributed islands extending over a one- or two-monolayer (ML) distance from the ideal atomically abrupt interface.[2-4] The resulting distribution of excitonic energies is Gaussian and, therefore, the interface roughness will result in an inhomogeneous Gaussian broadening of the excitonic emission line shape.

The application of an external magnetic field to the excitonic system is expected to influence both (1) the Lorentzian contribution to the emission line shape due to the magnetic-field-induced drastic changes in the densities of states of the carriers which form the exciton, and (2) the Gaussian contribution owing to the decrease of the effective volume of the exciton (an information volume of the optical probe) relative to the lateral dimensions of the potential fluctuations.

In this work we demonstrate that the application of the magnetic field to the system of excitons confined two dimensionally in a plane of the QW results in a narrowing of the Lorentzian contribution and a simultaneous broadening of the Gaussian contribution to the excitonic emission line shape. The observed narrowing of the Lorentzian contribution is explained by an increase of the carrier relaxation time, which is caused by the formation of orthogonal quantum states [zero-dimensional (0D) system] in the magnetic field, whereas the broadening of the Gaussian contribution is attributed to an increased sensitivity of the exciton wave function to the potential fluctuations at the QW interfaces, caused by a decrease of the information volume of the exciton.

II. EXPERIMENT

We have studied a high-quality undoped GaAs QW sample with well width $L_z=170$ Å, which was grown on a CrO-doped [001]-oriented GaAs substrate by molecular-beam epitaxy (MBE) technique without growth interruption. The GaAs well layer was sandwiched between barrier layers of composition $Al_{0.23}Ga_{0.77}As$ with a thickness of 500 Å.

Photoluminescence (PL) was excited by the 514.5-nm line of a cw Ar-ion laser, with the sample being located in a superconducting magnetic. A magnetic field of magnitude up to 6 T was applied perpendicular to the plane of the QW, and the exciting and emitted light direction was parallel to the direction of the magnetic field (Faraday configuration). The temperature of the sample was about 6 K, and the density of photoexcited carriers per unit plane N_S, estimated from the value of the excitation power 160 mW/cm^2 and the excitonic radiative lifetime for our QW in the range 200–300 ps,[5] was of the order of $N_S\sim10^8$ cm^{-2}. The spectrometer resolution was about 0.7 Å.

III. RESULTS AND DISCUSSION

Typical PL spectra of the QW sample exhibit emissions resulting from the recombination of electrons (el) with both heavy (hh) and light (lh) holes. The emissions show

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a monotonous increase in intensity with increasing magnetic field (a 2.2-fold increase with magnetic field changing from 0 to 6 T) due to the buildup of the excitonic oscillator strength, caused by the increased probability of finding an electron and a hole at the same point in two-dimensional (2D) space due to the shrinking of the exciton wave function in nonzero magnetic field.[6] The full width at half maximum (FWHM) of the emissions is seen to be as small as ~ 0.5 meV, which speaks of the high optical quality of the samples under investigation.

A. Diamagnetic shift

The diamagnetic (Langevin) shift of the emissions, reflecting a change in the energy of the ground states of the hh and lh excitons under the influence of the magnetic field, is plotted in Fig. 1, the lh-exciton emission also showing a nonlinear Zeeman spin splitting in agreement with the results of Ossau *et al.*[7] To analyze the results for the diamagnetic shift, it should be noted that the relative motion of an optically excited electron-hole (*e-h*) pair in the magnetic field is determined by an effective potential consisting of two terms, the first representing the electrostatic Coulomb *e-h* interaction, and the second being imposed by the magnetic field. In the absence of the Coulomb interaction the *e-h* eigenstates would be energetically discrete Landau levels. The electrostatic interaction correlates these Landau-level states into magnetoexcitons, the inter-Landau-level *e-h* Coulomb interaction having been shown to persist up to very high magnetic fields $\lambda=8$.[8] The dimensionless parameter λ is a measure of the effect of the magnetic field on the exciton

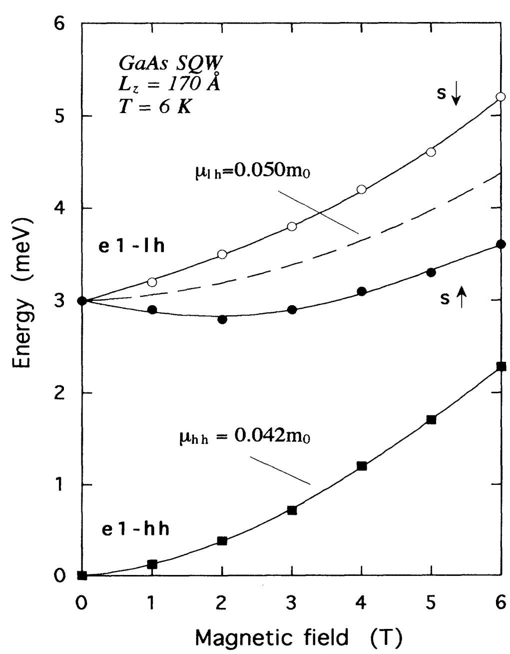

FIG. 1. Relative diamagnetic shift of the emission energy for the hh and lh excitons as a function of the magnetic field. The dashed line shows the theoretical result for zero spin splitting of the lh exciton.

wave function determined by the ratio of the cyclotron energy of electrons and holes $\hbar\omega c_{h,l}=e\hbar H/c\mu_{h,l}$, and the exciton binding energy $R^*_{h,l}=e^4\mu_{y,1}/2e\hbar^2$, i.e.,

$$\lambda_{h,l}=\hbar\omega_{c_{h,l}}/2R^*_{h,l} \ , \qquad (1)$$

where $\mu_{h,l}$ stands for the reduced mass of the exciton corresponding to heavy- (μ_h) or light- (μ_l) hole bands, in the plane of the quantum well. The values of $\mu_{h,l}$ can be expressed in terms of the well-known Kohn-Luttinger parameters γ_1 and γ_2 as[9]

$$\mu_{h,l^{-1}}=m_e^{-1}+(\gamma_1\pm\gamma_2)/m_0 \ , \qquad (2)$$

where m_0 is the free-electron mass. Taking the values of the Kohn-Luttinger parameters $\gamma_1=6.98$ and $\gamma_2=2.1$,[10] we obtain the in-plane exciton masses $\mu_h=0.042m_0$ and $\mu_l=0.050m_0$. Using Eq. (1), and taking the values of the exciton binding energies, calculated for the QW with a barrier of finite height with composition $x=0.23$, as $R_h^*=6.9$ meV and $R_l^*=7.4$ meV,[11,12] we then obtain for the QW under investigation that the magnetic energy is equal to the Coulomb energy ($\lambda=1$) at magnetic fields of 5 and 6.4 T for the hh- and lh-excitons, respectively. Therefore, for the magnetic fields used in this study ($\leq 6T$), we are in the regime of the intermediate magnetic field, and we must treat both the Coulomb and magnetic-field energy on an equal basis. For this reason we are unable to fit our experimental results on the diamagnetic shift to the expression[6]

$$\Delta E_{h,l}(1s)=3e^2\rho^2H^2/16\mu_{h,l} \ , \qquad (3)$$

valid only in the low-magnetic-field limit, for any choice of the in-plane exciton mass $\mu_{h,l}$. The parameter ρ in Eq. (3) is a variational parameter of the hydrogenic *s*-state wave function $\phi_{ex}(\mathbf{r})=(2/\pi)^{1/2}\rho^{-1}\exp(-r/\rho)$ used in calculations,[6] this parameter being a measure of the lateral spread of the exciton wave function.

The observed diamagnetic shifts at magnetic fields higher than 2 T are found to be no longer proportional to the square of the magnetic field, as predicted by Eq. (3), but gradually bend down approaching a linear dependence. The reason for such behavior is that the exciton wave function in nonzero magnetic field shrinks in real space, i.e., the parameter ρ decreases. Hence it is necessary to use the general expression for the change in energy of the exciton ground state[7]

$$\Delta E_{h,l}(1s)=(e^2H^2/8\mu_{h,l})\langle\Psi|x^2+y^2|\Psi\rangle_{h,l} \ , \qquad (4)$$

where the term in brackets is the expectation value of the squared lateral dimensions of the hh and lh excitons, respectively. In three dimensions the expectation value is $2a_B^2$, where a_B is the exciton Bohr radius $a_B=e\hbar^2/\mu_{h,l}e^2$, whereas for a perfectly 2D exciton the expectation value becomes $(\frac{3}{8})a_B^2$.[13] Since the expectation value of the lateral dimensions of the exciton is determined solely by its in-plane mass $\mu_{h,l}$, one can fit the experimental results for the diamagnetic shift to Eq. (4) in order to deduce the values of $\mu_{h,l}$, while retaining the 3D value of a_B.[13] On the other hand, one can use the experimental results and Eq. (4) to determine the magnetic-field dependence of the lateral exciton dimensions while keeping the zero-field value of $\mu_{h,l}$. Using our results for the diamagnetic shifts and the values of $\mu_h = 0.042m_0$ and $\mu_l = 0.050m_0$ we, therefore, calculated the values of $\langle \Psi | x^2 + y^2 | \Psi \rangle$ for both hh and lh excitons as a function of the magnetic field, the results for the hh exciton being used in the following emission line-shape analysis (Sec. III C).

B. Excitonic emission line shape: Lorentz-Gaussian profile

The line-shape function $I(E)$ of the excitonic emission from the QW as a function of photon energy E is given by a convolution integral as[3]

$$I(E) \sim f(E) \int_0^\infty D_{\rm ex}(E,E') B(E', E_{\rm ex}) dE' , \qquad (5)$$

where $D(E)$ is the density of states (DOS), $B(E)$ is a Lorentzian broadening function due to the final state interaction, and $f(E)$ is the thermal distribution function given, at low excitation, by a Maxwell-Boltzmann distribution. The el-hh exciton recombination occurs at the energy

$$E_{\rm ex} = E_g + E_{e1}(L_z) + E_h(L_z) - R_h^*(L_z) , \qquad (6)$$

where E_g is the 3D band gap of the QW material, E_{e1} and E_h stand for L_z-dependent energies of the lowest 2D electron and heavy-hole subbands, respectively, and R_h^* is the hh exciton binding energy. In the MBE-grown QW samples the well width L_z has been found to be statistically distributed over a certain "information volume" of exciton (which is equal to the exciton volume in case of the excitonic recombination), the excitons being localized in randomly distributed atomically smooth interface islands with the lateral dimensions $\delta_2 \geq 2a_B$, extending 1- or 2-ML distance δ_1 from the ideal atomically abrupt interface ("atomically smooth island" picture),[14] or being scattered by interface roughness on an atomic scale ("microroughness" picture).[15] It has also been shown recently that both the above pictures maybe inadequate, the real QW interface consisting of ~60-nm to ~1-μm size islands which themselves have substantial microroughness.[16] Since L_z is statistically distributed, the energies E_{e1}, E_h, and $E_{\rm ex}$ will also be distributed statistically, and $E_{\rm ex}$ in Eqs. (5) and (6) has to be replaced by its probability distribution $P(E_{\rm ex}, \langle E_{\rm ex} \rangle)$ and its expectation value $\langle E_{\rm ex} \rangle$, respectively, with a Gaussian being a good approximation for this probability distribution.[17,18] It should be noted here that the Gaussian contribution to the excitonic emission line shape, caused by an alloy disorder in the QW barrier layers,[5] can be neglected in our case due to the negligible amplitude of the hh-exciton wave function in the barriers for the QW layer with $L_z = 170$ Å.[11]

The experimentally measured line shape of the excitonic emission is, therefore, a convolution of two distribution functions; (1) the Lorentzian contribution with the full width at half maximum (FWHM) Γ_0, arising from the final-state interaction, and, at nonzero temperature, the exciton-phonon interaction; and (2) the Gaussian contribution with the FWHM σ, due to the potential fluctuations at the QW interface. This Lorentz-Gaussian function, as it follows from Eq. (5), needs to be further convoluted with the quasi-2D DOS, which changes from a step function to a δ function by application of the magnetic field, and then multiplied by a Maxwell-Boltzmann factor, which is characterized by an exciton temperature.

The detailed PL spectra of the $e1$-hh exciton emission for three magnitudes of the external magnetic field are shown in Fig. 2, together with the Lorentzian fits of the emission line shapes. It can be seen that at zero magnetic field the emission line shape is almost "pure" Lorentzian,

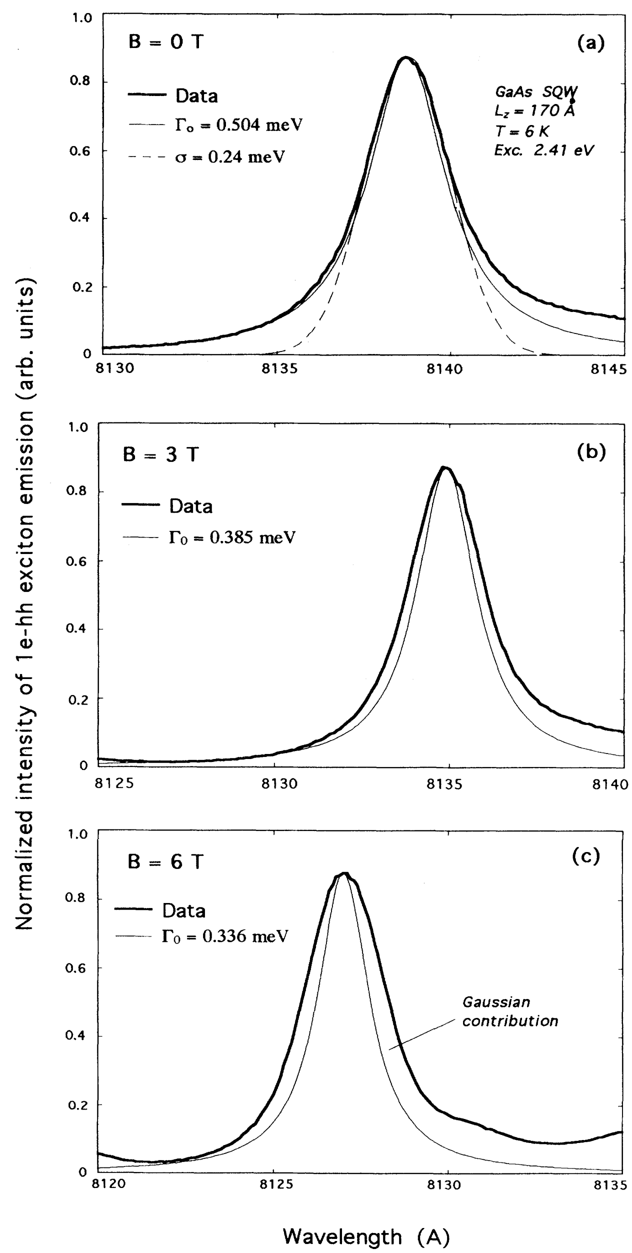

FIG. 2. Detailed PL spectra of the $e1$-hh exciton emissions from the SQW sample together with Lorentzian fits of the wings for magnitudes of the magnetic field 0 T (a), 3 T (b), and 6 T (c). For spectrum (a), an attempted, but obviously unsuccessful, Gaussian fit is also shown.

with a FWHM $\Gamma_0=0.504\pm0.01$ meV. Therefore, it is concluded that the emission line is predominantly lifetime broadened. The exciton temperature, estimated from the expression $\Gamma_0=kT_{ex}\ln 2$,[19] was found to be $T_{ex}=8.5$ K, which is just a few degrees higher than the temperature of the crystal lattice $T_c\sim 6$ K. Our experiment, therefore, can be considered as having been carried out in the low excitation regime.

It has been found from resonant Raleigh scattering and hole-burning experiments,[20,21] as well as from theoretical calculations,[22] that within an inhomogeneous excitonic emission line shape there are two types of excitons; i.e., the low-energy excitons, which are localized in the potential fluctuations at the interface and have a very narrow homogeneous linewidth $\Gamma_0\sim 0.04$ meV; and delocalized ones with $\Gamma_0\sim 2$ meV (for $L_z=170$ Å). The great difference in the values of Γ_0 for localized and delocalized excitons is caused by fewer and weaker dephasing mechanisms for localized excitons (only phonon scattering or tunneling) in comparison with those for the delocalized ones, which in addition are scattered elastically by the potential fluctuations as they move through the QW.

Considering the values of $\Gamma_0=0.504$ meV, $T_{ex}=8.5$ K, and $L_z=170$ Å, it follows from the results of Takagahara[22] that in our case the lifetime broadening arises from both the localized and delocalized excitons having an averaged dephasing time $\tau=\hbar/\Gamma_0\sim 1.3$ ps, the dominant mechanism of dephasing relaxation for these excitons being the scattering by acoustic phonons.

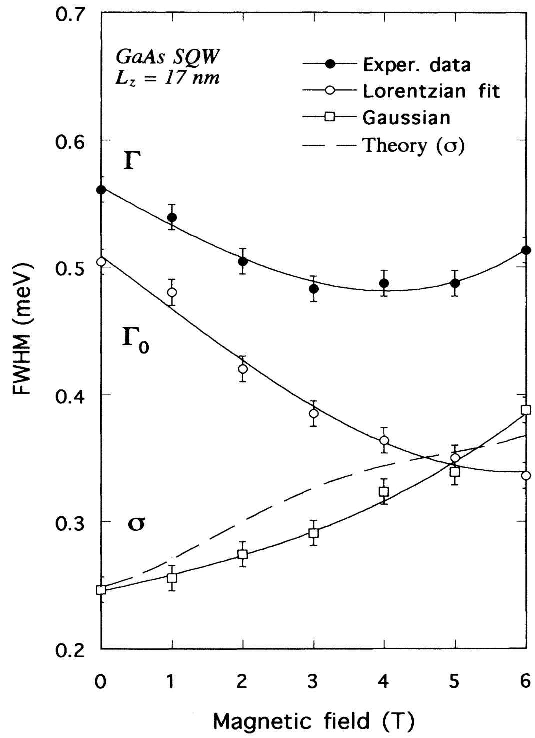

FIG. 3. The magnetic-field dependence of the experimentally measured excitonic emission linewidth Γ together with that of the Lorentzian (Γ_0) and Gaussian [$\sigma=(\Gamma^2-\Gamma_0^2)^{1/2}$] contributions. The theoretical curve for $\sigma(H)$, calculated on the basis of the interface model, developed by Singh and co-workers (Refs. 2 and 4) is also shown.

C. Emission line shape under magnetic field

1. Lorentzian contribution

The application of an external magnetic field results in a considerable narrowing of the Lorentzian contribution to the emission line shape (Fig. 2), the broadening parameter Γ_0 of the Lorentzian function having been determined by fitting our experimental data on the wings of the line shapes, since the wings of Lorentz-Gaussian profiles can be described by a pure Lorentzian function.[23] The resulting dependence of Γ_0 on the magnetic field is shown in Fig. 3, the value of Γ_0 decreasing from $\Gamma_0(0\ \mathrm{T})=0.504\pm0.01$ meV down to $\Gamma_0(6\ \mathrm{T})=0.336\pm0.01$ meV.

There are several mechanisms which, in general, may be responsible for the change in the linewidth of the Lorentzian contribution to the excitonic emission line shape, caused by the application of the magnetic field: (1) shrinking of the exciton wave function due to the increasing overlap of the *e-h* pair, which forms the exciton;[6,17] (2) modulation of the number of optically created carriers, caused by the formation of Landau levels from the background steplike 2D DOS, the number of carriers being dependent on the energy position of the photons of the excitation light in relation to that of Landau levels;[24,25] (3) filling-factor-dependent modulation of Landau-level linewidth;[26,27] and (4) formation of a zero-dimensional DOS for carriers which form magnetoexcitons,[28] this formation *in itself* being the origin of the change in Γ_0.

The first three mechanisms can, however, be neglected due to the following.

(1) The increased *e-h* overlap has been found to cause a reduction of the exciton radiative lifetime, with the recombination probability P_{rad} varying as $\langle\Psi|x^2+y^2|\Psi\rangle^{-1}$ [the expectation value of the exciton squared lateral dimensions determined in Eq. (4)].[29,30] The increase in P_{rad} will cause an increase in the number of excitons, which recombine before they cool down to the temperature of the lattice. Under increasing magnetic field we therefore can expect the broadening of the Lorentzian contribution due to the buildup of T_{ex},[31] which is contrary to our experimental results.

It should also be noted that the effect an interaction of the exciton with a radiation field, resulting in a rapid radiative decay of the exciton in low-dimensional systems (about 3 ps for a GaAs QW),[22] has been neglected in our study since this effect can be considered only as a small (short-lived) contribution to the integrated PL intensity in our experiment carried out under cw excitation.

(2) Since the photon energy of the excitation light $h\nu=2.41$ eV used in our study is higher than the band gap of the barrier material $E_{Al_xGa_{1-x}As}$ $(x=0.23)=1.8$ eV, the carriers are generated, under optical excitation, in the barrier material, not in the QW itself. They subsequently diffuse into the QW layer, where they thermalize

and form excitons which collapse within an averaged radiative lifetime. If the linewidth of the exciton emission had depended on the concentration of carriers generated in the barrier layers (this concentration being modulated by the magnetic field), then we would have observed the change in Γ_0 by changing the excitation light intensity. However, we have not detected any discernible change in the linewidth of emission by changing the excitation intensity (maintaining the low excitation regime) in the range 10–160 mW/cm^2, which rules out mechanism (2).

(3) The filling-factor-dependent modulation of the Landau-level linewidth can be neglected in our case due to the very low concentration of carriers ($N_S \sim 10^8$ cm^{-2}), giving a Landau-level filling factor $\nu = N_S h/eH$ of about 0.5, even at magnetic field h as small as 0.01 T. All optically created carriers are, therefore, at the bottom of the first Landau level at any magnetic fields used in this study.

As a result of the above considerations, we attribute the observed decrease in the Lorentzian linewidth Γ_0 to the magnetic-field-induced drastic changes in the density of states [$D(E)$ in Eq. (5)] for carriers which form magnetoexcitons, i.e., to the transformation of the steplike 2D DOS given (per unit area) by[33,34]

$$D_i^{2D}(E) = \frac{m_i}{\pi\hbar^2} \sum_{n=1}^{\infty} H[E - E_{n_i}] , \tag{7}$$

where m_i ($i = e$, l, and h) stands for the effective mass of electron, light hole, or heavy hole in the direction perpendicular to the QW layer, respectively [is not to be confused with the in-plane exciton mass $\mu_{h,l}$ given by Eq. (2)], $H[x]$ is the Heaviside function, and E_{n_i} denotes the quantized energy levels; into a discrete 0D DOS given by[28]

$$D_i^{0D}(E) = (\hbar\omega_c) \left[\frac{2m_i}{\hbar^2} \right] \times \sum_{k=0}^{\infty} \sum_{n=1}^{\infty} \delta[E - E_{n_i} - (k + 1/2)\hbar\omega_c] , \tag{8}$$

where the indexes n and k denote the nth 2D subband due to the electrical confinement and kth Landau level, respectively.

For the e1-hh exciton, formed from the carriers which occupy the lowest electrical ($n = 1$) and magnetic ($k = 0$) subbands, Eq. (8) can be rewritten in a simple form

$$D_{ex}(E) = \frac{2eH}{c\hbar} \left[\frac{M_h}{\mu_h} \right] \delta \left[E - E_{1e} - E_{hh} - \frac{e\hbar H}{2c\mu_h} \right] , \tag{9}$$

where $M_h = m_e + m_{hh}$, and $\delta[x]$ is the delta function. Equation (9) is valid to the extent that the broadening of the Landau levels can be neglected. To take into account the lifetime and thermal broadening of the Landau levels, we should replace the δ function in Eq. (9) by a Lorentzian. Then, denoting $E' = e\hbar H/2c\mu_h$ and $E_0 = E - E_{1e} - E_{hh}$ (i.e., by taking the energy of the exciton at $H = 0$ as zero energy), we obtain

$$D_{ex}(E) = \frac{2eH}{c\hbar} \left[\frac{M_h}{\mu_h} \right] \frac{\Gamma_0}{(E_0 - E')^2 + \Gamma_0^2} , \tag{10}$$

where $\Gamma_0 = \hbar/\tau$ is the FWHM of the Lorentzian contribution, and τ is the scattering time.

The experimentally observed decrease in the value of Γ_0 in a magnetic field can, therefore, be attributed to the increase in the exciton dephasing time τ. The application of the magnetic field, therefore, causes not only a decrease in the exciton dimensions, but also a simultaneous increase of the coherence volume of exciton.[31] The increase in the dephasing time τ with increasing magnetic field is due to the suppression of the carrier thermalization, which is caused by the formation of a discrete 0D (atomiclike) DOS, the last conclusion being in agreement with the results of recent studies of dynamics of the magnetoexcitons in the QW,[35,36] as well as with theoretical calculations.[37,38]

2. Gaussian contribution

Returning back to Fig. 2, we can see that the increase of the magnetic field results not only in the above-discussed narrowing of the Lorentzian contribution to the emission line shape, but also in a growing misfit between the observed line shape and the Lorentzian contribution, deduced from the fitting at the wings of the spectra. This misfit is caused by an increase of the Gaussian contribution to the emission line shape. The magnetic-field dependence of the Gaussian contribution to the emission linewidth σ, shown in Fig. 3, has been deconvoluted from the values of the emission FWHM Γ and the Lorentzian FWHM Γ_0 as $\sigma = (\Gamma^2 - \Gamma_0^2)^{1/2}$, i.e., by assuming, in the first approximation, an independence of the lifetime and statistical broadenings. It can be seen that the value of σ exhibits a 60% increase from $\sigma(0\ \text{T}) = 0.24$ meV to $\sigma(6\ \text{T}) = 0.39$ meV.

Due to the very small amplitude of the hh-exciton wave function in the barrier layers for the QW with $L_z = 170$ Å,[11] we can neglect any Gaussian contribution to the emission line shape arising from an alloy disorder in the barrier layers.[39–42] We therefore attributed the Gaussian contribution to the emission line shape to the statistical potential fluctuations at the QW interfaces, a simplified theory of the effect of these fluctuations on the emission linewidth having been established by Singh and Bajaj.[2,4]

Their theory is based on the experimentally verified model of the nonideal QW interface as consisting of randomly distributing islands with lateral dimensions δ_2, extending a distance $z = \delta_1$ from the ideal (average) interface at $z = 0$ ($L_z = L_z^0$).[15] The theory then correlates the standard deviation σ_0 of the Gaussian distribution function, which describes the line shape of the exciton emission due to statistical fluctuations in L_z, with the atomic scale parameters of the interface as

$$\sigma_0 = \left[\sqrt{p(1-p)} \frac{\delta_1 \delta_2}{2R_{ex}} \right] \frac{\delta E_{ex}}{\delta L_z} \bigg|_{L_z = L_z^0} , \tag{11}$$

where p is the averaged coverage of the interface, and $R_{ex} = \langle \Psi | x^2 + y^2 | \Psi \rangle^{1/2}$ is the lateral extension of the exciton wave function [see Eq. (4)]. Equation (11) was obtained assuming that the quality of the QW is determined primarily by the quality of one interface, due to the poorer quality of the inverted interface, formed by the growth of GaAs on $Al_xGa_{1-x}As$, in comparison with that of the normal interface, formed by the growth of $Al_xGa_{1-x}As$ on top of GaAs.[43,44]

Since the binding energy of the exciton is not expected to change much with small fluctuations of L_z values, from Eq. (6) one obtains that for the e1-hh exciton $(\delta E_{ex}/\delta L_z) = (\delta E_{eh}/\delta L_z)$, where $E_{eh} = E_{1e} + E_h$ is the L_z-dependent sum of the local energies of an electron and a hole, forming excitons in their respective 2D subbands. In the infinitely high barrier approximation, we obtain for the 1e-hh exciton that $E_{eh} = \pi^2 \hbar^2 / 2\mu' L_z^2$, where μ' is the reduced exciton mass in the direction perpendicular to the QW layer [which is not to be confused with the in-plane exciton mass $\mu_{h,l}$ given in Eq. (2)]. Then, taking into account that the FWHM σ of the Gaussian distribution is 2.35 times its standard deviation σ_0,[39] we obtain the FWHM as

$$\sigma = 1.18 \frac{\pi^2 h^2}{\mu' L_z^3} \frac{\delta_1 \delta_2}{R_{ex}} \sqrt{p(1-p)} \,. \tag{12}$$

We note that the width of the Gaussian line shape is determined, for given QW geometrical parameters δ_1, L_z, and p, by the lateral spread of the exciton wave function R_{ex} relative to that of the interface islands δ_2. Since the application of magnetic field is expected to cause the shrinking of the exciton wave function, i.e., the decrease in the value of the exciton lateral dimensions $R_{ex} = \langle \Psi | x^2 + y^2 | \Psi \rangle^{1/2}$, one can anticipate an increase in σ due to the magnetic field, in perfect agreement with the experimental results of Figs. 2 and 3.

We have calculated the perpendicular hh-exciton mass from the Kohn-Luttinger parameters $\gamma_1 = 6.98$ and $\gamma_1 = 2.1$,[10] using the expression $\mu'^{-1} = m_e^{-1} + m_0^{-1}(\gamma_1 - 2\gamma_2)$,[9] as being equal to $\mu' = 0.057 m_0$, as well as the lateral exciton dimensions at zero magnetic field $R_{ex} = \varepsilon \hbar^2 / \mu_h e^2$, taking the hh-exciton in-plane mass $\mu_h = 0.042 m_0$ (see Sec. III A), as $R_{ex} = 160$ Å, the value of R_{ex} being in agreement with the theoretical results of Campi.[45] Then, taking the value of δ_1 as being equal to 1 ML ($\delta_1 = 2.83$ Å), and assuming the most probable value of $p = 0.5$ from Eq. (12) and the deconvoluted value of $\sigma(0\ \mathrm{T}) = 0.24$ meV, we calculated the lateral dimensions of the interface islands for our QW to be $\delta_2 = 100$ Å.

Since in Eq. (12) only one parameter, namely R_{ex}, depends on the magnetic field, we can use the values of the parameters μ' and δ_2, determined for zero magnetic field, as well as the values of R_{ex} for various magnitudes of the magnetic field, deduced from the experimental data on the diamagnetic shift in Sec. III A, to calculate the magnetic-field dependence of the linewidth of the Gaussian contribution to the exciton emission line shape. The results of such calculations are shown in Fig. 3. It can be seen that, although the range of change in the values of σ, expected from the theoretical considerations with the magnetic field changing from 0 to 6 T, agrees well with the values of σ, deconvoluted from the experimental results, the forms of the theoretical and experimental curves are different. This discrepancy has been attributed to the simplifying assumptions made in the course of theoretical treatment. For one, the substantial microroughness on the atomic scale inside the interface islands[16] may be taken as an explanation for the higher deconvoluted values of σ in comparison with those expected from the interface islands model, at a magnetic field higher than 5 T. Therefore, for a consistent description of the Gaussian contribution to the excitonic luminescence line shape, one needs to go beyond the simple model of Singh and Bajaj[2,4] and take the Fourier transform of the 2D interface roughness.[16,46]

IV. CONCLUSION

In conclusion, we have studied the effect of magnetic field on the e1-hh excitonic photoluminescence line shape, being a convolution of the Lorentzian and Gaussian contributions, in a single GaAs-$Al_xGa_{x-1}As$ quantum well grown by MBE. An increase of magnetic field from 0 to 6 T has been found to result in (1) a decrease in the Lorentzian contribution Γ_0 to the line shape from $\Gamma_0(0\ \mathrm{T}) = 0.504 \pm 0.01$ meV to $\Gamma_0(6\ \mathrm{T}) = 0.336 \pm 0.01$ meV due to the formation of a quasi-zero-dimensional density of states with a following increase in the exciton dephasing time due to the inhibition of the carrier relaxation in the 0D case; and (2) an increase in the Gaussian contribution from $\sigma(0\ \mathrm{T}) = 0.24$ meV to $\sigma(6\ \mathrm{T}) = 0.39$ meV, attributed to the shrinking of the exciton lateral dimensions $\langle \Psi | x^2 + y^2 | \Psi \rangle$, the last effect causing the exciton to become more responsive to the statistical potential fluctuations at the quantum-well interfaces.

[1] P. T. Landsberg, Phys. Status Solidi **15**, 623 (1966).

[2] J. Singh, K. K. Bajaj, and S. Chaudhuri, Appl. Phys. Lett. **44**, 805 (1984).

[3] D. Bimberg, J. Christen, T. Fukunaga, H. Nakashima, D. E. Mars, and J. N. Miller, J. Vac. Sci. Technol. B **5**, 1191 (1987).

[4] J. Singh and K. K. Bajaj, J. Appl. Phys. **57**, 5433 (1985).

[5] J. Kusano, Y. Segawa, Y. Aoyagi, and S. Namba, Solid State Commun. **65**, 925 (1988).

[6] M. Sugawara, N. Okazaki, T. Fujii, and S. Yamazaki, Phys. Rev. B **48**, 8848 (1993).

[7] W. Ossau, B. Jakel, E. Bangert, and G. Landwehr, Surf. Sci. **174**, 188 (1986).

[8] C. Stafford, S. Schmitt-Rink, and W. Schaefer, Phys. Rev. B **41**, 10 000 (1990).

[9] M. Altarelli and N. O. Lipari, Phys. Rev. B **9**, 1733 (1974).

[10] K. Hess, D. Bimberg, N. O. Lipari, J. U. Fishbach, and M. Altarelli, in *Proceedings of the 13th International Conference on Physics of Semiconductors, Rome, 1976*, edited by F. G.

Fumi (North-Holland, Amsterdam, 1977), p. 142.
[11]R. L. Greene, K. K. Bajaj, and D. E. Phelps, Phys. Rev. B **29**, 1807 (1984).
[12]R. L. Greene, K. K. Bajaj, and D. E. Phelps, Solid State Commun. **45**, 831 (1983).
[13]G. Bastard, E. E. Mendez, L. L. Chang, and L. Esaki, Phys. Rev. B **26**, 1974 (1982).
[14]C. Wiesbuch, R. C. Miller, R. Dingle, A. C. Gossard, and W. Wiegmann, Solid State Commun. **37**, 219 (1981).
[15]D. Bimberg, D. Mars, J. N. Miller, R. Bauer, and D. Oertel, J. Vac. Sci. Technol. B **4**, 1014 (1986).
[16]C. A. Warnick and R. F. Kopf, Appl. Phys. Lett. **60**, 386 (1992).
[17]C. Colvard, D. Bimberg, K. Alavi, C. Maier, and N. Nouri, Phys. Rev. B **39**, 3419 (1989).
[18]M. Sugawara, T. Fujii, S. Yamazaki, and K. Nakajima, Phys. Rev. B **42**, 9587 (1990).
[19]J. Christen and D. Bimberg, Phys. Rev. B **42**, 7213 (1990).
[20]J. Hegarty, M. D. Sturge, C. Wiesbuch, A. C. Gossard, and W. Wiegmann, Phys. Rev. Lett. **49**, 930 (1982).
[21]J. Hegarty and M. D. Sturge, J. Opt. Soc. Am. B **2**, 1143 (1985).
[22]T. Takagahara, Phys. Rev. B **32**, 7013 (1985).
[23]J. Humlicek, E. Schmidt, L. Bosanek, R. Svehla, and K. Ploog, Phys. Rev. B **48**, 5241 (1993).
[24]Th. Englert, J. C. Maan, Ch. Uihlein, D. C. Tsui, and A. C. Gossard, Solid State Commun. **46**, 545 (1983).
[25]D. Heitmann, M. Ziesmann, and L. L. Chang, Phys. Rev. B **34**, 7463 (1986).
[26]M. Potemski, Y. H. Zhang, M. Hausenblas, G. Martinez, and K. Ploog, Surf. Sci. **263**, 591 (1992).
[27]M. S. Scolnick, J. M. Rorison, K. J. Nash, D. J. Mowbray, P. T. Tapster, S. J. Bass, and A. D. Pitt, Phys. Rev. Lett. **58**, 2130 (1987).
[28]Y. Arakawa and A. Yariv, IEEE J. Quantum Electron. **22**, 1887 (1986).
[29]J. Christen, D. Bimberg, A. Steckenborn, and G. Weimann, Appl. Phys. Lett. **44**, 84 (1984).
[30]E. O. Gobel, H. Jung, J. Kuhl, and K. Ploog, Phys. Rev. Lett. **51**, 1588 (1983).
[31]J. Feldmann, G. Peter, E. O. Gobel, P. Dawson, K. Moore, C. Foxon, and R. J. Elliott, Phys. Rev. Lett. **59**, 2337 (1987).
[32]E. Hanamura, Phys. Rev. B **38**, 1228 (1988).
[33]C. W. J. Beenakker and H. van Houven, in *Solid State Physics,* edited by H. Ehrenreich and D. Turnbull (Academic, New York, 1991), Vol. 44, p. 1.
[34]A. Ishihara, in *Solid State Physics,* edited by H. Ehrenreich and D. Turnbull (Academic, New York, 1989), Vol. 42, p. 271.
[35]J. B. Stark, W. H. Knox, and D. S. Chemla, Phys. Rev. Lett. **68**, 3080 (1992).
[36]J. B. Stark, W. H. Knox, D. S. Chemla, W. Schafer, S. Schmitt-Rink, and C. Stafford, Phys. Rev. Lett. **65**, 3033 (1990).
[37]U. Bockelmann and G. Bastard, Phys. Rev. B **42**, 8947 (1990).
[38]T. Inoshita and H. Sakaki, Phys. Rev. B **47**, 7260 (1992).
[39]E. F. Schubert, E. O. Gobel, Y. Horikoshi, K. Ploog, and H. J. Quesser, Phys. Rev. B **30**, 813 (1984).
[40]J. Singh and K. K. Bajaj, Appl. Phys. Lett. **44**, 1075 (1984).
[41]S. M. Lee and K. K. Bajaj, Appl. Phys. Lett. **60**, 853 (1992).
[42]J. Singh and K. K. Bajaj, Appl. Phys. Lett. **48**, 1077 (1986).
[43]R. F. Kopf, E. F. Schubert, T. D. Harris, and R. S. Becker, Appl. Phys. Lett. **58**, 631 (1991).
[44]H. Markos, T. J. Drummond, and R. Fisher, J. Appl. Phys. **53**, 1032 (1982).
[45]D. Campi and C. Villavecchia, IEEE J. Quantum Electron. **28**, 1765 (1992).
[46]C. A. Warnick, W. Y. Jan, A. Ourmard, and T. D. Harris, Appl. Phys. Lett. **56**, 2666 (1990).