Influence of external electric and magnetic fields on the excitonic absorption spectra of quantum-well wires

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The quantum-confined Stark effect and the influence of magnetic confinement on the optical absorption spectra for a highly excited quantum-well wire are analyzed. The shifts of the exciton resonance and the changes of its oscillator strength are evaluated. The calculated spectra of magnetoabsorption and of magnetoluminescence for a modulation-doped quantum wire both agree qualitatively with corresponding recent measurements.

I. INTRODUCTION

Semiconductor quantum-well wires are a subject of considerable current interest. Since the early work with strain-induced[1,2] or with intermixed wires,[3] more promising wire structures have been obtained, e.g., with V grooving,[4] etching, and overgrowing.[5] Typical one-dimensional (1D) effects such as blueshifts of the single-particle energies,[6] center-of-mass quantization of excitons,[7] and enlarged exciton binding energies[4] have already been reported. But in most experimental situations, particularly under high excitation,[5] several subbands (often not even resolved) are occupied. By using two-photon absorption techniques,[8] $2p$ excitons have been created which feel a stronger confinement due to their larger radius.

Naturally, the most striking one-dimensional effects are expected in the quasi-1D limit where only the lowest subband is involved. Theoretical studies of the linear optical spectra of quantum-well wires showed[9] that excitonic effects completely regularize the square-root singularity of the one-dimensional density of states at the band gap. For the multisubband case, it has been shown that band-mixing effects with[10] and without[11] a magnetic field lead to nonparabolic dispersions and to an optical anisotropic polarization dependence.[12,13]

The single quantum-well wire considered here may be built from a quasi-two-dimensional (2D) quantum well (grown in the z direction) by introducing an additional lateral confinement. We assume that the effective-mass approximation remains valid and do not consider valence-band mixing explicitly. For sufficiently thin quantum wells (which are, however, still wide enough to keep the carriers inside), the originally degenerate heavy- and light-hole energies at $k=0$ split due to their different confinement energies. Therefore, the upper 2D valence band has mainly heavy-hole character. By introducing a further lateral confinement, the 2D heavy-hole band splits further in one-dimensional subbands. Although any one-dimensional valence subband has contributions from heavy- and light-hole bulk bands, it can display a dominant heavy- or light-hole character.[10] Since we only want to take into account the lowest one-dimensional subband, we neglect band-mixing effects and treat the valence band as a pure heavy-hole band. The shape of the lateral confinement potential in the x direction is assumed to be parabolic. Particularly for intermixed wires, laser-induced diffused wires,[14] nonplanar-grown V-grooved wires, or modulation-doped wires,[15] the assumption of a smooth lateral confinement potential is more realistic than a square-well potential. Without external fields, an electron-hole symmetry is assumed, leading to a local charge neutrality. This symmetry would be exact for a confinement potential with infinite walls. The one-dimensional Coulomb interaction is described by averaging the bare three-dimensional (3D) Coulomb interaction (at $z=0$) with the lateral envelope wave function.[9]

II. QUANTUM-CONFINED STARK EFFECT

In this section we investigate the influence of a static, homogeneous electric field E perpendicular to the wire axis. As long as the lateral confinement is strong enough to keep the particles in the wire, a quantum-confined Stark effect results which is well known from quantum wells.[16] Within the parabolic confinement approximation, the electric field causes a spatial separation of the electron (hole) wave functions and a lowering of the single-particle energy.[17] Without many-body corrections, the single-particle energies are given by (omitting the band gap E_g^{2D} of the quantum well and putting $\hbar=1$)

$$\epsilon_k^{j_0}=\frac{k^2}{2m_j}-\frac{e^2E^2}{2m_j\Omega_j^2}+\frac{\Omega_j}{2},\quad j=e,h\ , \tag{1}$$

where k is the wave number along the wire axis (y axis), m_j is the electron (hole) mass, and Ω_j is the intersubband spacing in the parabolic confinement potential. The first term is the usual kinetic energy in the wire direction, the second term describes the quadratic Stark effect, and the last term is the lateral confinement energy. The corresponding single-particle wave functions

$$\psi_k^j(x,y)=\frac{1}{\sqrt{L_y}}e^{iky}\left[\frac{m_j\Omega_j}{\pi}\right]^{1/4}e^{-(m_j\Omega_j/2)(x+e_jE/m_j\Omega_j^2)^2} \tag{2}$$

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are plane waves in the wire direction (L_y is the wire length) and shifted Gaussian functions in the x direction. In the presence of an (e.g., optically generated) electron-hole plasma, the single-particle energies are shifted by exchange-correlation effects and Hartree corrections.[16] The Hartree potential ϕ results from the spatial separation of the electrons and holes. The new wave functions and particle energies have to be determined by solving Poisson's and Schrödinger's equations self-consistently. As an approximation, we use an iterative procedure. The charge density $\rho(\mathbf{r})$ reads

$$\rho(\mathbf{r})=\frac{Ne}{L_y}\delta(z)[|\psi^e(x)|^2-|\psi^h(x)|^2] . \tag{3}$$

Here, the charge density is assumed to be uniform along the wire axis (N is the number of electrons) and to be located in the quantum well at $z=0$. From the (formal) solution of Poisson's equation, one gets for the Hartree potential ϕ,

$$\phi(x,y,z=0)=\frac{Ne}{\epsilon_0 L_y}\int_{-\infty}^{\infty}\frac{dq}{|q|}\int_{-\infty}^{\infty}dx' e^{iq(x-x')} \times[|\psi^e(x')|^2 -|\psi^h(x')|^2] . \tag{4}$$

Here, ϵ_0 represents the static dielectric constant of the well material. Treating the Hartree potential as a perturbation with respect to the shifted lateral confinement potential, one finds for the renormalization of the single-particle energies,

$$\delta\epsilon_H^j=\frac{2Ne^2}{\epsilon_0 L_y}\int_0^{\infty}dq\, e^{-q^2/4m\Omega}\frac{1-\cos(q\Delta)}{q}, \quad \Delta=\frac{m_e+m_h}{(m\Omega)^2}eE . \tag{5}$$

In the low-field limit one gets

$$\delta\epsilon_H^j=\frac{2(m_e+m_h)^2e^4}{\epsilon_0 m^3\Omega^3}\frac{N}{L_y}E^2 . \tag{6}$$

Here, $\Omega=\Omega_e+\Omega_h$ is the total subband spacing, m is the reduced electron-hole mass, and Δ is a measure of the field-induced electron-hole separation. The Hartree contribution is in lowest order (like the usual Stark shift) quadratic in the electric field but with opposite sign. Naturally, the Hartree contribution vanishes for zero plasma density. As mentioned above, the single-particle energies are also renormalized by exchange and correlation effects. In Ref. 18 it is shown that the screening in one-dimensional systems is of minor importance compared to phase-space filling effects. The exchange contribution to the single-particle energies is given by

$$\delta\epsilon_{\mathrm{xc}}^j=-\sum_q V^{jj}(q)f_q^j , \tag{7}$$

where f_q^j is the Fermi distribution function and $V^{jj}(q)$ an effective one-dimensional interaction:

$$V^{jj}(q)=\frac{e^2}{\epsilon_0 L_y}e^{q^2/4m_j\Omega_j}K_0\left[\frac{q^2}{4m_j\Omega_j}\right] . \tag{8}$$

K_0 is a modified Bessel function. The electron-hole interaction in the presence of a lateral external field is described by

$$V^{\mathrm{eh}}(q)=-\frac{2e^2}{\epsilon_0 L_y}\int_0^{\infty}dk\, e^{-k^2/2m\Omega}\frac{\cos(k\Delta)}{(k^2+q^2)^{1/2}} . \tag{9}$$

The reduced overlap between the lateral electron and hole wave functions causes a reduction of the electron-hole interaction ($\Delta=0$ if no field is present) and a decrease of the optical dipole transition element d_{eh} (assumed to be a constant in k space):

$$d_{\mathrm{eh}}(E)=d_{\mathrm{eh}}(E=0)e^{-m\Delta^2\Omega/2} . \tag{10}$$

For weak lateral confinement or in type-II quantum wires,[19] additional contributions from optical interband transitions of subbands with different quantum numbers have to be taken into account. Since we consider a simple two-band model, we neglect this effect. The optical spectra of a quantum-well wire with a quasiequilibrium electron-hole plasma are obtained from the solutions of the equation for the interband polarization P_k:[20]

$$i\frac{d}{dt}P_k=(\epsilon_k^e+\epsilon_{-k}^h-i\gamma)P_k -(1-f_k^e-f_k^h) \times\left[d_{\mathrm{eh}}(E)E_t e^{-i\omega t}-\sum_{k'}V^{\mathrm{eh}}(k-k')P_{k'}\right] . \tag{11}$$

The first term on the right-hand side (rhs) of the polarization equation is the unperturbed time development (including the renormalized single-particle energies $\epsilon_k^j=\epsilon_k^{j0}+\delta\epsilon_H^j+\delta\epsilon_{\mathrm{xc}}^j$), where a phenomenological lifetime $1/\gamma$ has been included. The inhomogeneous second term shows how the polarization is driven by the laser test field E_t. The last term gives rise to excitonic effects. In explicit calculations we use the material parameters of GaAs and measure all energies and lengths in units of the three-dimensional exciton Rydberg E_0 and exciton Bohr radius a_0, respectively.

Figure 1 shows the influence of an external field on the linear (i.e., without a plasma) optical absorption

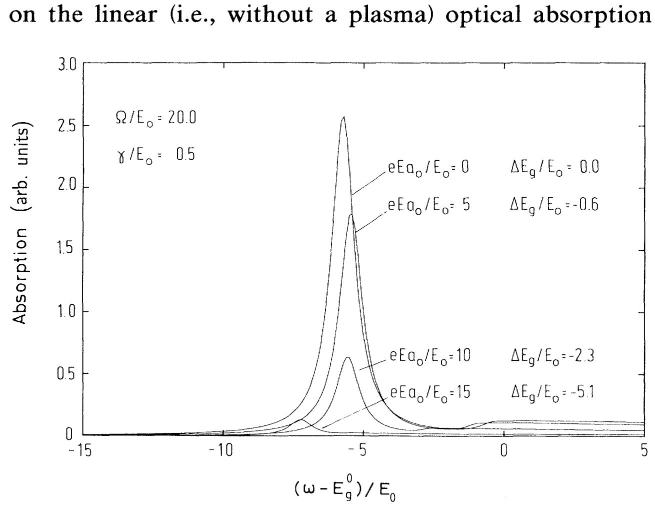

FIG. 1. Linear excitonic absorption spectra for various lateral electric fields.

spectra. The absorption coefficient is proportional to the imaginary part of the optical susceptibility $\chi = d/V\sum_k P_k/E_t$ (V is the wire volume). In order to be consistent with the assumed one-subband limit, we have chosen a rather large intersubband spacing $\Omega \simeq 84$ meV which has been realized recently.[21] Without field, one gets a strong excitonic absorption energetically well below the one-dimensional band gap E_g^0. By switching on the electric field, the excitonic absorption bleaches due to the reduced electron-hole overlap. For low fields we observe a small blueshift of the excitonic resonance. This indicates that in quantum wires the reduction of the exciton binding energy due to the reduced overlap can be larger than the redshift of the band gap (denoted by ΔE_g). Due to a compensation of the Stark shift and the reduction of the binding energy, the excitonic resonance shifts much less than the band gap. Physically, this effect represents the charge neutrality of an exciton. For large electric fields the Stark effect dominates so that the absorption peak shifts red. In corresponding experimental spectra on quantum wells[16] and variational calculations of the exciton position,[22] a monotonous redshift of the exciton absorption has been found. This redshift is accompanied by a rapid broadening due to inhomogeneous well-width fluctuations. Due to the two-dimensional confinement in quantum wires, this broadening effect is expected to be further enhanced in real quantum-well wires.

Figure 2 shows the influence of a thermal electron-hole plasma on the excitonic absorption spectra. The solid lines are calculated without field, the dashed lines with an electric field. In both cases the presence of the plasma causes a bleaching of the excitonic resonance. In the absence of a field, one gets a small redshift of the resonance due to the incomplete compensation of exchange renormalization and reduction of the binding energy by Coulomb blocking.[18] The situation changes drastically in the presence of an external field. Due to the violation of charge neutrality, the Hartree term exists and shifts the excitonic absorption peak to higher energies for increasing plasma densities. This feature may also influence the measurements of gap renormalization in realistic quantum-well wires, because imperfections and impurities in the lateral barrier region may result in internal electric fields.

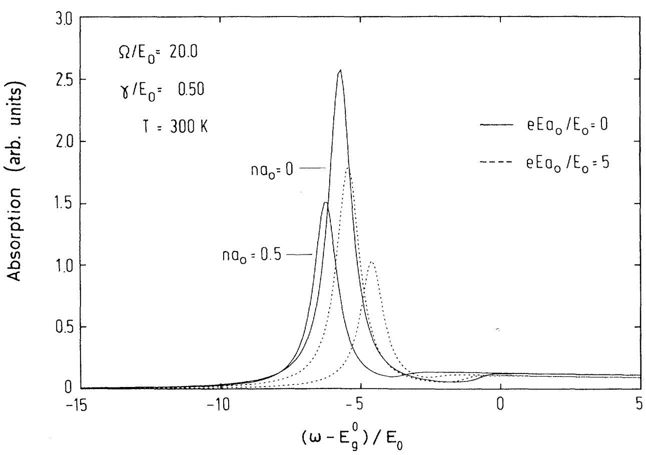

FIG. 2. Bleaching of the excitonic resonance due to an electron-hole plasma without (solid lines) and with (dashed lines) a lateral electric field.

III. MAGNETO-OPTICAL SPECTRA

Next we discuss the influence of an external magnetic field **B** on the optical absorption spectra. The magnetic field is oriented perpendicular to the quantum-well layer (i.e., **B** parallel to the growth axis z of the quantum well). Using again the parabolic lateral confinement model, one obtains for the single-particle energies:[21]

$$\epsilon_{kn}^j = \frac{\Omega_j^2}{\Omega_j^2+\omega_c^{j\,2}} \frac{k^2}{2m_j} + (\Omega_j^2+\omega_c^{j\,2})^{1/2}(n+\tfrac{1}{2}) \pm g_j \mu_B B \;, \qquad n=0,1,\ldots \;, \tag{12}$$

where $\omega_c^j = e_j B/m_j c$ are the individual ($j=e,h$) cyclotron energies, $\mu_B = e/2m_0 c$ is the Bohr magneton, and g_j are the effective Landé factors.[23] The magnetic field enhances the subband spacing,[24] increases the effective mass in the wire direction, and splits the otherwise spin-degenerated states. Again, only the ground state ($n=0$) is taken into account. The corresponding single-particle wave functions are now given by

$$\psi_k^j(x,y) = \frac{1}{\sqrt{L_y}} e^{iky} \left[\frac{m_j \Omega_j^{\rm eff}}{\pi} \right]^{1/4} \times \exp\left[-m_j \Omega_j^{\rm eff}/2 \left[x - \frac{k}{m_j} \frac{\omega_j^c}{\Omega_j^2+\omega_c^{j\,2}} \right]^2 \right] . \tag{13}$$

They consist of a plane wave in the wire direction and k ("velocity") dependent shifted ground-state oscillator functions. $\Omega_j^{\rm eff} = (\Omega_j^2+\omega_c^{j\,2})^{1/2}$ is an effective intersubband spacing according to Eq. (12). The one-dimensional Fourier-transformed Coulomb interaction $V_{kk'}^{jj'}(q)$ between two particles with wave numbers k and k' becomes explicitly velocity dependent:

$$V_{kk'}^{jj'}(q) = \int \frac{dq'}{2\pi} dx\, dx' e^{i[q'(x-x')]} \Phi_{k'+q}^{*j'}(x') \times \Phi_{k'}^{j'}(x') \frac{2\pi e_j e_{j'}}{\epsilon_0 L_y (q^2+q'^2)^{1/2}} \Phi_{k-q}^{*j}(x) \Phi_k^j(x) \;, \tag{14}$$

where Φ is the lateral part of the wave function Ψ. Without fields, we assume again the electron-hole symmetry ($m_e\Omega_e = m_h\Omega_h$). Therefore, one obtains from Eq. (14),

$$V_{kk'}^{jj'}(q) = \frac{e_j e_{j'}}{\epsilon_0 L_y} e^{-(\omega_c/\Omega)^2(q^2/2m\Omega)} \times \int_{-\infty}^{\infty} \left[\frac{dx\, e^{-x^2}}{\left[x^2 + \frac{q^2}{2m\Omega} \right]^{1/2}} \right] \times \cos\left[2\frac{\omega_c}{\Omega} \left[\frac{k'-k+q}{\sqrt{2m\Omega}} \right] x \right] . \tag{15}$$

Here,

$$\Omega=(\Omega_e^{\mathrm{eff}^2}+\Omega_h^{\mathrm{eff}^2})^{1/2}$$

denotes the total electron-hole subband spacing (including the magnetic field) and ω_c is the cyclotron energy calculated with the reduced electron-hole mass. Within the Hartree-Fock approximation for the interband polarization, only special wave-number combinations [e.g., $V_{k,k-q}(q)$ and $V_{0,-q}(q)$] are needed. Due to symmetry relations, these interactions depend only on q. Therefore, one obtains for the effective one-dimensional Coulomb interaction in the presence of a magnetic field the analytical expression

$$V^{jj'}(q)=\frac{e_j e_{j'}}{\epsilon_0 L_y} e^{-(\omega_c/\Omega)^2(q^2/2m\Omega)+q^2/4m\Omega} K_0\left[\frac{q^2}{4m\Omega}\right] . \tag{16}$$

For zero magnetic field, one recovers Eq. (8). For large magnetic fields (i.e., $\Omega\rightarrow\infty$) the Coulomb interaction shows a logarithmic divergence due to the magnetic confinement. Equations (16) and (12) show that changes in the optical absorption spectra can only be expected if the cyclotron energy is of the same order of magnitude as the lateral subband spacing. Because we treat the strong confinement limit, changes in the optical spectra will be seen only for large magnetic fields where the magnetic confinement becomes comparable with the wire confinement. In order to see magnetic effects, we have chosen a relatively large magnetic field of $B=20$ T. In Ref. 21 quantum wires with an estimated subband spacing of 80 meV have been investigated in magnetic fields up to 40 T. Figure 3 shows the bleaching of the excitonic absorption due to an electron-hole plasma with (solid lines) and without (dashed lines) a magnetic field. The magnetic confinement acts in a twofold manner: First, it increases the single-particle energies. Second, the magnetic confinement enhances the effective Coulomb interaction. In all spectra the excitonic absorption peak shifts to higher energies in increasing magnetic fields.

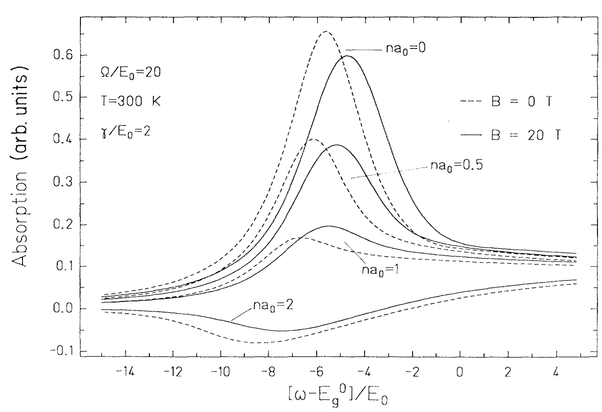

FIG. 3. Magnetoabsorption spectra with (solid lines) and without (dashed lines) a magnetic field for various plasma densities.

This indicates that the single-particle confinement is larger than the energy gain by the enhanced Coulomb interaction and shows that the continuum states (above the band gap) are more influenced by the magnetic field than the bound state. Total changes in the oscillator strength of the excitonic absorption due to a magnetic field are caused by three partly compensating effects (a)–(c): (a) The Zeeman term produces a small subband splitting (although not visible) and therefore an accompanying hardly visible broadening and bleaching. The enlarged effective Coulomb interaction (b) and the enhanced single-particle density of states (c) lead to an increasing oscillator strength. In the linear and low plasma-density spectra, one observes a small decrease of the oscillator strength due to the magnetic field. For higher plasma densities the exciton resonance is enhanced by the magnetic field because of the (now dominating) enlarged density of states. This effect also causes a higher density of transparency, so that the optical gain in the high excitation spectrum is reduced by the field. Because of the chosen large broadening γ, the Zeeman splitting cannot be seen in the spectra in agreement with corresponding experimental observations.[21]

Up to now, we have discussed a high-field effect because of the assumed large subband spacing. Next, we apply our model to a doped quantum wire. Plaut *et al.*[15] measured magnetoluminescence spectra of a modulation-doped quantum wire. The observed luminescence process is due to a recombination of free electrons with photoexcited holes bound to acceptors in a separated Be δ layer. Therefore, we change our model in the following way: The Be impurity level is modeled by an infinity hole mass. The reduction of the Coulomb interaction between the electrons and the photoexcited holes is taken into account by reducing their interaction by a factor $\exp(-|qd|)$, where d is an average electron-hole distance. The Zeeman splitting is neglected because it is a fine-structure effect and the dominant contribution stems from the hole splitting.[11] Naturally, luminescence can only occur in the presence of holes. Therefore, we calculate the luminescence rate per hole and take the limit of vanishing hole density. For the calculation of luminescence spectra, it is useful to write Eq. (11) in terms of a susceptibility

$$\chi_k(\omega)=\chi_k^0(\omega)-\chi_k^0(\omega)\sum_{k'} V^{\mathrm{eh}}(k-k')\chi_{k'}(\omega) . \tag{17}$$

In order to extract the Bose factor between emission and absorption, we use the free susceptibility χ^0 in the spectral representation:[18]

$$\chi_k^0(\omega)=\int_{-\infty}^{\infty}\frac{d\omega'}{\pi}\frac{\gamma}{(\omega'-\epsilon_k)^2+\gamma^2}\frac{1-f^e(\omega')}{\omega-\omega'+i\delta} \tag{18}$$

$$\simeq \frac{\omega-\epsilon_k}{(\omega-\epsilon_k)^2+\gamma^2}[1-f^e(\epsilon_k)] - i\frac{\gamma}{(\omega-\epsilon_k)^2+\gamma^2}[1-f^e(\omega)] . \tag{19}$$

Finally, the spontaneous emission rate $R(\omega)$ is expressed by

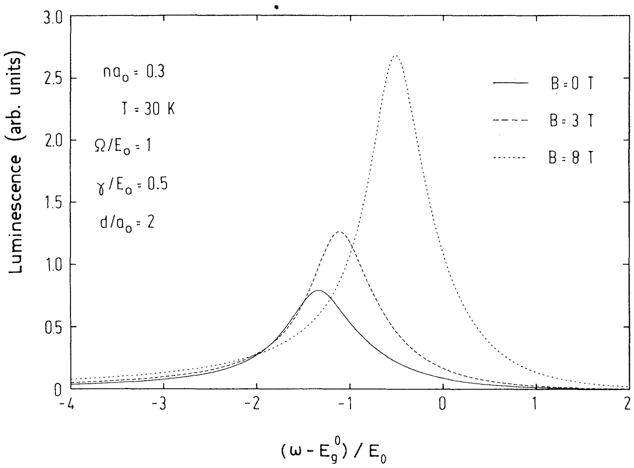

FIG. 4. Magnetoluminescence spectra of an n-doped quantum wire for various magnetic fields.

$$R(\omega) = \text{const} \times e^{-\beta(\omega-\mu)} \text{Im} \sum_k \chi_k(\omega) \ , \tag{20}$$

where μ is the chemical potential and β the inverse thermal energy of the electrons. Figure 4 shows the luminescence spectra for a fixed electron density and temperature and various magnetic fields. The intersubband spacing in this modulation-doped wire is relatively small ($\simeq 4.2$ meV). Therefore, changes in the optical spectra are already obtained for low magnetic fields. Increasing the magnetic field up to 8 T, a blueshift due to the Lorentz term and an increasing oscillator strength of the optical transition are obtained. The increasing oscillator strength results from the enlarged one-particle density of states and the enhancement of the Coulomb interaction. The quantum wire results differ strongly from those obtained for quantum wells both experimentally[15] and theoretically.[25] Here, a splitting of the Landau levels has been observed. In order to reproduce the experimental results, we have taken a relatively high electron temperature of 30 K (the lattice temperature has been only 1.5 K). For low electron temperatures we would obtain a pronounced Coulomb enhancement at the Fermi edge. Such a Fermi-edge singularity has been observed in similar structures[26] recently.

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