Many-body effects in the absorption, gain, and luminescence spectra of semiconductor quantum-well structures

S. Schmitt-Rink
AT&T Bell Laboratories, Murray Hill, New Jersey 07974

C. Ell and H. Haug
Institut für Theoretische Physik, Universität Frankfurt, Robert-Mayer-Strasse 8, D-6000 Frankfurt am Main, Federal Republic of Germany
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Absorption, gain, and luminescence spectra of quasi-two-dimensional electron-hole plasmas in semiconductor quantum-well structures are calculated as functions of the plasma density and temperature. Self-energy corrections and the effects of multiple electron-hole scattering are evaluated for a statically screened Coulomb interaction. The Bethe-Salpeter equation for the electron-hole pair propagator is solved both numerically and analytically using a method developed by Noyes. The resulting spectra deviate considerably from the corresponding free-particle spectra, due to the strong pair fluctuations in two dimensions. Implications for the theory of quantum-well lasers are discussed.

I. INTRODUCTION

In the last few years, the optical properties of direct-gap semiconductor quantum-well structures found a growing interest due to their potential device applications in integrated optics and optical data processing.[1] In potential wells whose thickness is comparable to or smaller than the bulk exciton Bohr radius, excitonic effects are strongly enhanced due to the spatial confinement of the electrons and holes and image forces.[2] The excitons show an increased binding energy and an enhanced oscillator strength so that, e.g., they are still well resolved in room-temperature absorption spectra.[1]

Under strong laser excitation, exchange effects and the screening of the Coulomb interaction by the optically generated electron-hole pairs destabilize the exciton and an electron-hole plasma forms[3,4] (due to the short lifetime there will be *no* electron-hole liquid condensation in direct-gap materials). In this paper we investigate theoretically the optical properties of this plasma. It can be expected that even in the plasma state the effects of multiple electron-hole scattering are still very pronounced and considerably larger than in three dimensions,[5] due to the strong enhancement of (electron-hole pair) fluctuations under quantum confinement.

For clarity, we limit ourselves to an idealized model of photoexcited semiconductor quantum-well structures. Assuming the mean particle distance and the thermal wavelength to be much larger than the layer thickness, we consider only the lowest electron and hole subbands. Furthermore, we assume that these bands are parabolic. We do not address ourselves to problems which are related to the strong valence-band mixing in quantum wells. As has been discussed for modulation-doped systems,[6,7] the symmetry breaking in semiconductor quantum-well structures is strongly enhanced by the "shake-up" of the Fermi sea. Similar effects are expected for photoexcited systems.

Within our simple model, the photoexcited carriers interact via the two-dimensional (2D) Coulomb interaction $V(r)=e^2/(\epsilon_0 r)$, where ϵ_0 is the background dielectric constant of the barrier material. At low excitation intensities, the electrons e and holes h form excitons with binding energy $E_0=2me^4/\epsilon_0^2$ and Bohr radius $a_0=\epsilon_0/(2me^2)$, where m is the reduced electron-hole mass, $m^{-1}=m_e^{-1}+m_h^{-1}$. Above a critical excitation density, these excitons cease to exist and a quasimetallic electron-hole plasma forms.[4] In Sec. II we calculate the corresponding renormalization of the single-particle energies (the band-gap shrinkage) due to exchange and correlation effects. The correlation effects are treated statically, using the plasmon-pole approximation for the screened Coulomb interaction. Results for the band-gap renormalization within the latter approximation have already been reported in Refs. 4 and 8. They agree reasonably with recent experimental observations.[9,10] In Sec. III we calculate the optical spectra of highly excited quantum wells by solving the Bethe-Salpeter equation (BSE) for the e-h pair Green's function within the statically screened ladder approximation. This singular integral equation describes the multiple scattering of electrons and holes via the statically screened Coulomb interaction. Using a method developed by Noyes[11] (see also Ref. 12), the BSE is reduced to a nonsingular integral equation which is solved both analytically and numerically by matrix inversion. From the imaginary part of the e-h pair Green's function we obtain the absorption, gain, and luminescence spectra. The results of our calculation are presented in Sec. IV for various plasma densities and temperatures. As expected, the two-dimensional pair fluctuations cause strong deviations from the corresponding one-particle spectra, so that, e.g. the gain spectra bear great resemblance to bulk spectra. The relevance of our results to the theory of quantum-well lasers is discussed.

II. SCREENING AND SINGLE-PARTICLE PROPERTIES

Under high optical excitation a large concentration of electron-hole pairs is created in a semiconductor. Clearly, one has to include exchange effects as well as the screening of the long-range Coulomb forces in order to obtain a reasonable description of such a dense system of charged fermions. Various approximation schemes have been developed for that purpose. For three-dimensional (3D) systems, they have been summarized in our recent review article.[5] Some of these ideas have been extended to quasi-2D systems.[4,8] A discussion of the subband level renormalization due to exchange and correlation effects has been given and some shortcomings of standard theories (which are not so significant in three dimensions) have been pointed out.[4] In the following, we give a brief discussion of the subband level renormalization within a static approximation which is based on the random-phase approximation (RPA) and expected to yield reasonable results for the densities and temperatures of interest.

In quasi-2D e-h systems, the screened Coulomb interaction is given by

$$V_s(q,\omega)=\frac{2\pi e^2}{\epsilon(q,\omega)q}\ , \tag{1}$$

where $2\pi e^2/q$ is the bare 2D Coulomb interaction and $\epsilon(q,\omega)$ is the longitudinal dielectric function. Within the RPA $\epsilon(q,\omega)$ reads ($\hbar=k_B=1$)

$$\epsilon(q,\omega)=\epsilon_0\left[1-2V(q)\sum_{i,\mathbf{k}}\frac{f_{i0}(k)-f_{i0}(\mathbf{k}+\mathbf{q})}{\omega+i\delta+\epsilon_i(k)-\epsilon_i(\mathbf{k}+\mathbf{q})}\right]\ , \tag{2}$$

where $V(q)=2\pi e^2/(\epsilon_0 q)$ and $i=e,h$; $\epsilon_i(k)=E_g/2+k^2/(2m_i)$ are the unrenormalized single-particle energies characterized by effective electron and hole masses m_i and an energy gap E_g; $f_{i0}(k)$ are the corresponding distribution functions.

For practical purposes, the total RPA excitation spectrum can be replaced by a single effective plasmon mode, so that[4,5,8,13–17]

$$\epsilon^{-1}(q,\omega)=\epsilon_0^{-1}\left[1+\frac{\omega_p^2(q)}{(\omega+i\delta)^2-\omega^2(q)}\right]\ , \tag{3}$$

where $\omega_p(q)$ is the 2D plasma frequency, $\omega_p^2(q)=2\pi n e^2 q/(\epsilon_0 m)$, which has in contrast to the 3D plasma frequency a square-root dispersion. This result can be derived in an elementary way by combining the equation of continuity with Newton's equation.

$\omega(q)$ is the frequency of the effective plasmon mode

$$\omega^2(q)=\omega_p^2(q)\left[1+\frac{q}{\kappa}\right]+\frac{q^4}{16m^2}\ . \tag{4}$$

Here, κ is the screening wave number which is determined by the compressibility sum rule:

$$\kappa=\frac{2\pi e^2}{\epsilon_0}\sum_i\frac{\partial n}{\partial\mu_{i0}}\ , \tag{5}$$

where n is the plasma density and

$$\mu_{i0}=E_g/2+T\ln\left[\exp\left[\frac{\pi n}{m_i T}\right]-1\right] \tag{6}$$

are the unrenormalized quasichemical potentials of the electrons and holes.

The dispersion (4) reproduces the correct RPA screening behavior at long wavelengths and simulates the individual particle-hole pair excitations at short wavelengths. For $T\to 0$, the screening wave number κ reduces to a constant $\kappa_0=2e^2(m_e+m_h)/\epsilon_0$ due to the constant density of states.

Results for the exchange-correlation (xc) contributions $\mu_{i\mathrm{xc}}$ to the quasichemical potentials $\mu_i=\mu_{i0}+\mu_{i\mathrm{xc}}$ within the plasmon-pole approximation have already been reported in Refs. 4 and 8. In particular, it has been shown that for $T\to 0$ and moderate densities the corresponding exchange-correlation potential $\mu_{\mathrm{xc}}=\mu_{e\mathrm{xc}}+\mu_{h\mathrm{xc}}$ varies approximately like $\mu_{\mathrm{xc}}\sim -3.1(na_0^2)^{1/3}E_0$. This result agrees reasonably with recent experimental observations, both for photoexcited quantum wells[10] and for quantum-well lasers in which a high-density $e-h$ plasma is electrically injected.[9] The results of Refs. 4 and 8 also agree approximately with the results of more sophisticated calculations,[18–21] which include, e.g., finite thickness effects.[21] Different results have been reported in Refs. 22 and 23.

In the following, we will limit ourselves to the "static limit" of the plasmon-pole approximation, by neglecting recoil effects completely.[5,14,24] This is consistent with our static approximation for the irreducible electron-hole interaction (see Sec. III). The price we have to pay for this consistency is an *overestimate* of the correlation energy as compared to the full dynamical result.

Within this simple model, the renormalized single-particle energies are given by

$$E_i(k)=\epsilon_i(k)+\Sigma_i(k)\ , \tag{7}$$

where $\Sigma_i(k)$ are the single-particle self-energies which describe the energy the particles gain by avoiding each other. This energy can be split into a screened exchange (sx) and a Coulomb-hole (Ch) term:

$$\Sigma_i(k)=\Sigma_{i\mathrm{sx}}(k)+\Sigma_{iCh}\ , \tag{8}$$

where

$$\Sigma_{i\mathrm{sx}}(k)=-\sum_{k'}V_s(\mathbf{k}-\mathbf{k}')f_i(k') \tag{9}$$

and

$$\Sigma_{iCh}=\tfrac{1}{2}\sum_{\mathbf{k}'}[V_s(k')-V(k')]\ . \tag{10}$$

Here, $V_s(k)=V_s(k,\omega=0)$ is the statically screened Coulomb interaction as determined by Eqs. (3)–(6) and

$$f_i(k)=\frac{1}{e^{[E_i(k)-\mu_i]/T}+1} \tag{11}$$

is the quasiparticle distribution function.

Equations (8)–(10) can be derived from the dynamical self-energies given in Refs. 4 and 8 by neglecting all recoil energies with respect to the plasma frequency. Equation (10) is nothing but the (classical) self-energy of a localized charge.

Obviously, as long as exchange effects are screened efficiently, the self-energies (8) depend only weakly on momentum, so that the electron and hole mass enhancement and the decrease of the bandwidths are small. The self-energies cause mainly rigid-band shifts, the renormalized band gap being given by

$$E_g' = E_g + \sum_i \Sigma_i(0) \ . \tag{12}$$

At very high densities exchange effects become dominant, so that the effective electron and hole masses decrease and the bandwidths increase. As in three dimensions, Eq. (12) depends only weakly on the electron-hole mass ratio m_e/m_h.

In the following, we will use the $T=0$ values of the self-energies (8), which is well justified for the densities and temperatures of interest. The electron and hole self-energies are then equal (which is an artifact of the static approximation), and the exchange-correlation contributions to the quasichemical potentials are given by

$$\mu_{i\mathrm{xc}} = \Sigma_i(k_F) \ , \tag{13}$$

where $k_F = (2\pi n)^{1/2}$ is the Fermi wave number.

III. ABSORPTION, GAIN, AND LUMINESCENCE

Within our simple model, the absorption, gain, and luminescence spectra associated with optical transitions between the lowest electron and hole subbands are determined by the interband density of states

$$D(\omega) = -\mathrm{Im} \sum_{\mathbf{k},\mathbf{k}'} G(\mathbf{k},\mathbf{k}',\omega) \ , \tag{14}$$

where $G(\mathbf{k},\mathbf{k}',\omega)$ is the retarded $e-h$ pair propagator. Apart from constant factors, Eq. (14) is just the imaginary part of the retarded photon self-energy and thus determines the difference between the photon scattering rates out of and into the state ω under simultaneous creation and annihilation of $e-h$ pairs.[5] Correspondingly, the absorption, gain, and luminescence spectra are given by $D(\omega)$, $-D(\omega)$, and

$$D(\omega)/\{\exp[(\omega-\mu/T)]-1\} \ ,$$

respectively, where $\mu = \mu_e + \mu_h$ is the quasichemical potential.

Within the statically screened ladder approximation, the two-particle Green's function $G(\mathbf{k},\mathbf{k}',\omega)$ fulfills the Bethe-Salpeter equation (BSE)

$$G(\mathbf{k},\mathbf{k}',\omega) = G_0(\mathbf{k},\mathbf{k}',\omega) - \sum_{\mathbf{k}'',\mathbf{k}'''} G_0(\mathbf{k},\mathbf{k}'',\omega) V_s(\mathbf{k}''-\mathbf{k}''') \times G(\mathbf{k}''',\mathbf{k}',\omega) \ , \tag{15}$$

where

$$G_0(\mathbf{k},\mathbf{k}',\omega) = \frac{1-f_e(k)-f_h(k)}{\omega + i\delta - E_e(k) - E_h(k)} \delta_{\mathbf{k},\mathbf{k}'} \ . \tag{16}$$

This singular integral equation describes the multiple scattering of an electron-hole pair via the statically screened Coulomb interaction in the presence of the $e-h$ Fermi sea. The phase-space occupation factor $F(k) = 1 - f_e(k) - f_h(k)$ reflects the exclusion principle; the renormalized single-particle energies $E_i(k)$ are determined by (8)—(10). Clearly, self-energy and vertex corrections are treated on equal footing.

Expanding the $e-h$ pair propagator in terms of the eigenfunctions of the homogeneous part of the BSE, one obtains after some algebra[4,25]

$$G(\mathbf{k},\mathbf{k}',\omega) = \sum_n \frac{|F(k)|^{1/2} \phi_n^*(\mathbf{k}) |F(k')|^{1/2} \phi_n(\mathbf{k}')}{\omega + i\delta - E_n} \times \mathrm{sgn}(E_n - \mu) \ , \tag{17}$$

where $\phi_n(\mathbf{k})$ are the eigenfunctions of the non-Hermitean eigenvalue problem

$$\sum_{\mathbf{k}'} H_{\mathbf{k},\mathbf{k}'} \phi_n(\mathbf{k}') = E_n \phi_n(\mathbf{k}) \tag{18}$$

with

$$H_{\mathbf{k},\mathbf{k}'} = [E_e(k) + E_h(k)] \delta_{\mathbf{k},\mathbf{k}'} - \mathrm{sgn}[F(k)] \, |F(k)|^{1/2} V_s(\mathbf{k}-\mathbf{k}') \, |F(k')|^{1/2} \ . \tag{19}$$

The eigenvalues E_n of (18) are real, and the eigenfunctions $\phi_n(\mathbf{k})$ form a complete set and obey the generalized orthogonality relation

$$\sum_{\mathbf{k}} \phi_m^*(\mathbf{k}) \, \mathrm{sgn}[F(k)] \phi_n(\mathbf{k}) = \delta_{m,n} \mathrm{sgn}(E_n - \mu) \ . \tag{20}$$

In the dilute limit $\mu/T \to -\infty$, Eq. (17) reduces to

$$G(\mathbf{k},\mathbf{k}',\omega) = \sum_{\mathrm{n}} \frac{\phi_{\mathrm{n}}^*(\mathbf{k}) \phi_{\mathrm{n}}(\mathbf{k}')}{\omega + \mathrm{i}\delta - \mathrm{E_n}} \ , \tag{21}$$

where

$$E_n \phi_n(\mathbf{k}) = [\epsilon_e(k) + \epsilon_n(k)] \phi_n(\mathbf{k}) - \sum_{\mathbf{k}'} V(\mathbf{k}-\mathbf{k}') \phi_n(\mathbf{k}') \tag{22}$$

and

$$\sum_{\mathbf{k}} \phi_m^*(\mathbf{k}) \phi_n(\mathbf{k}) = \delta_{m,n} \tag{23}$$

constitute the unperturbed hydrogenic exciton problem. Equation (14) reduces then to the simple expression

$$D(\omega) = \pi \sum_n |\phi_n(r=0)|^2 \delta(\omega - E_n) \ , \tag{24}$$

which is the well-known result of Ref. 26.

In the plasma state, bound states do no longer exist so that the eigenfunctions $\phi_n(\mathbf{k})$ are scattering states $\phi_{\mathbf{p}}(\mathbf{k})$, classified according to the relative $e-h$ momentum $\mathbf{p}$. The $e-h$ pair energies are $E(p) = E_e(p) + E_h(p)$. In analogy with standard definitions, $\phi_{\mathbf{p}}(\mathbf{k})$ can be written in the form

$$\phi_{\mathbf{p}}(\mathbf{k}) = \delta_{\mathbf{k},\mathbf{p}} + \frac{\mathrm{sgn}[F(k)] \, |F(k)|^{1/2} T(\mathbf{k},\mathbf{p}) \, |F(p)|^{1/2}}{E(p) - E(k) + i\delta} \ , \tag{25}$$

where $T(\mathbf{k},\mathbf{p})$ is the on-shell T matrix defined through

$$T(\mathbf{k},\mathbf{p})\,|F(p)|^{1/2}=-\sum_{\mathbf{k}'} V_s(\mathbf{k}-\mathbf{k}')\,|F(k')|^{1/2}\phi_{\mathbf{p}}(\mathbf{k}') \;. \tag{26}$$

Equation (25) corresponds to outgoing wave boundary conditions. Substituting (26) into (25), we obtain the BSE for the on-shell T matrix

$$T(\mathbf{k},\mathbf{p})=-V_s(\mathbf{k}-\mathbf{p})-\sum_{\mathbf{k}'} V_s(\mathbf{k}-\mathbf{k}')\frac{F(k')}{E(p)-E(k')+i\delta}\times T(\mathbf{k}',\mathbf{p}) \;. \tag{27}$$

Since only s-wave scattering contributes to (14), we can replace the statically screened Coulomb interaction $V_s(\mathbf{k}-\mathbf{k}')$ in the following by its angle-averaged value

$$V_s(k,k')=\frac{1}{2\pi}\int_o^{2\pi} d\rho\, V_s(\mathbf{k}-\mathbf{k}') \;, \tag{28}$$

where ρ is the angle between $\mathbf{k}$ and $\mathbf{k}'$.

With use of a method developed by Noyes[11] (see also Ref. 12), the BSE (27) can be transformed into a new nonsingular integral equation. Multiplying (27) for $T(p,p)$ with $V_s(k,p)/V_s(p,p)$ and subtracting the result from (27), one finds

$$T(k,p)=\frac{V_s(k,p)}{V_s(p,p)}T(p,p)-\sum_{\mathbf{k}'}\left[V_s(k,k')-\frac{V_s(k,p)V_s(p,k')}{V_s(p,p)}\right]\times\frac{F(k')}{E(p)-E(k')+i\delta}T(k',p) \;. \tag{29}$$

The solution of this equation can be expressed as $T(k,p)=T(p,p)f(k,p)$, where

$$f(k,p)=\frac{V_s(k,p)}{V_s(p,p)}-\sum_{\mathbf{k}'}\left[V_s(k,k')-\frac{V_s(k,p)V_s(p,k')}{V_s(p,p)}\right]\times\frac{F(k')}{E(p)-E(k')}f(k',p) \tag{30}$$

and

$$T(p,p)=\frac{-V_s(p,p)}{1+\Lambda_1(p)} \tag{31}$$

with

$$\Lambda_1(p)=\sum_{\mathbf{k}} V_s(p,k)f(k,p)\frac{F(k)}{E(p)-E(k)+i\delta} \;. \tag{32}$$

Obviously, the kernel of (30) is finite for $k'=p$ and $f(p,p)=1$.

The oscillator strength for transitions into scattering states is determined by

$$\sum_{\mathbf{k}} |F(k)|^{1/2}\phi_p(k)=|F(p)|^{1/2}\left[1-\frac{\Lambda_2(p)}{1+\Lambda_1(p)}\right] , \tag{33}$$

where

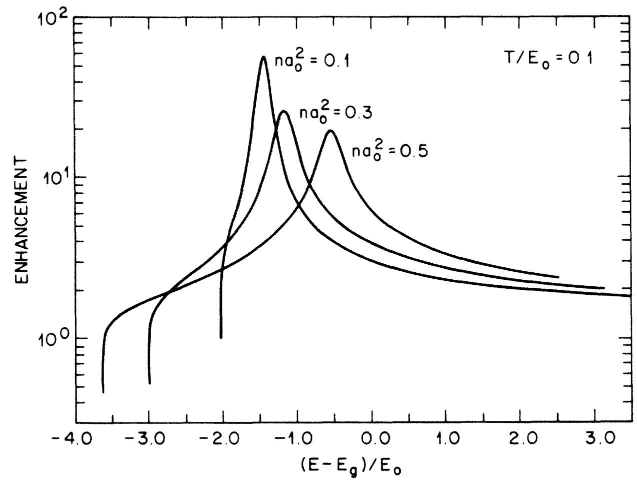

FIG. 1. Excitonic enhancement for $T/E_0=0.1$ and various plasma densities n.

$$\Lambda_2(p)=\sum_{\mathbf{k}} V_s(p,p)f(k,p)\frac{F(k)}{E(p)-E(k)+i\delta} \;. \tag{34}$$

Substituting this expression into (17) and (14), we obtain the final result for the interband density of states

$$D(\omega)=D_0(\omega)\rho(\omega) \;, \tag{35}$$

where

$$D_0(\omega)=\pi\sum_{\mathbf{p}} F(p)\delta(\omega-E(p)) \tag{36}$$

is the one-particle result (including the renormalization for the single-particle energies) and

$$\rho(\omega)=\left|1-\frac{\Lambda_2(p)}{1+\Lambda_1(p)}\right|^2_{E(p)=\omega} \tag{37}$$

is the excitonic enhancement due to the multiple $e-h$ scattering.

IV. RESULTS AND DISCUSSION

Performing the one-dimensional integration in (30) by means of a Gaussian quadrature, Eq. (30) reduces to a sys-

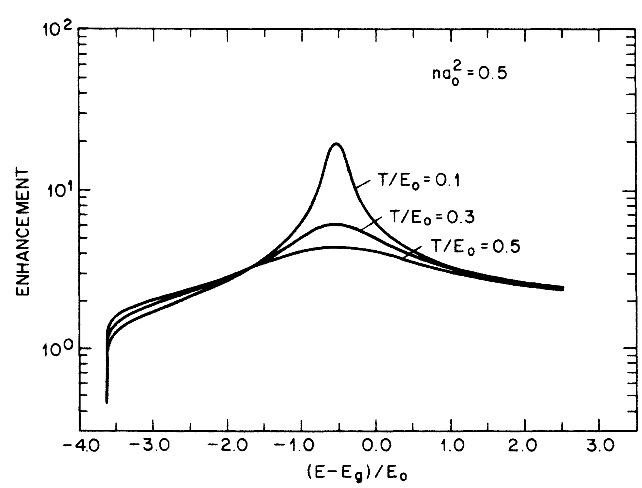

FIG. 2. Excitonic enhancement for $na_0^2=0.5$ and various plasma temperatures T.

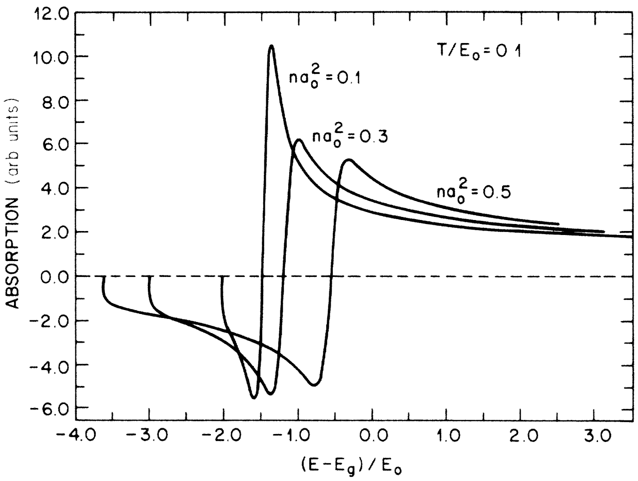

FIG. 3. Optical absorption for $T/E_0=0.1$ and various plasma densities n.

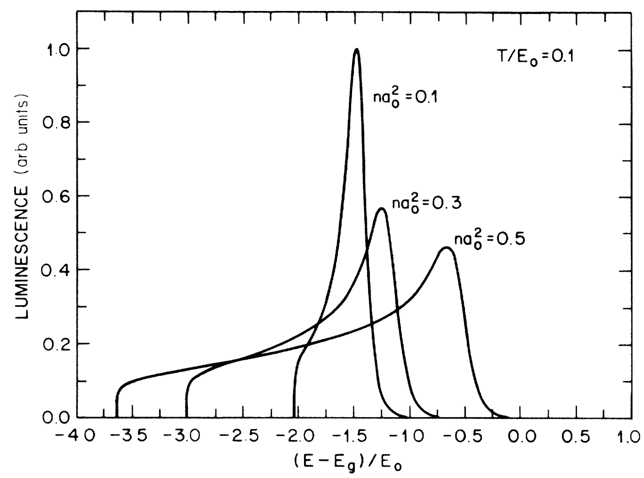

FIG. 5. Luminescence for $T/E_0=0.1$ and various plasma densities n.

tem of linear equations, which can be solved by matrix inversion. Typically, we used a mesh of 200 points in order to obtain a good overall accuracy. In the following, we will present results for $m_e/m_h=0.2$ for the set of densities $na_0^2=0.1$, 0.3, and 0.5 and temperatures $T/E_0=0.1,0.3$, and 0.5. For typical GaAs/$Al_xGa_{1-x}As$ quantum-well structures, these values correspond to densities between 10^{11} and 10^{13} cm^2 and temperatures below 150 K. In the degenerate limit, i.e., except for densities and temperatures close to the critical values for the appearance of excitons,[4] all our numerical results agree within graphical resolution with the simple analytical results obtained by keeping only the leading inhomogeneous term of (30).

The excitonic enhancement factor $\rho(\omega)$ is plotted in Figs. 1 and 2 for various plasma densities and temperatures, respectively. Obviously, the electron-hole correlation is very strong even in the plasma state. With increasing plasma densities and temperatures the maximum of $\rho(\omega)$ near the quasichemical potential μ decreases due to the decreasing e-h coupling strength and the increasing thermal broadening of the Fermi functions, respectively.

The shape of $\rho(\omega)$ reflects the tendency of the system to form e-h Cooper pairs, i.e., to undergo a transition to an excitonic insulator state. In 2D systems, there is no long-range order, as has been rigorously demonstrated in Ref. 27. However, there may be topological long-range order and a Berezinskii-Kosterlitz-Thouless transition may occur.[28,29] The excitonic enhancement reflects directly the fluctuations of the magnitude of the order parameter preceding the transition to the excitonic insulator state, similar e.g., to the enhancement of the paraconductivity in superconductors.[30] It is well known that these fluctuations are much stronger in two dimensions than in three dimensions. Hence, the quantum confinement enhances the e-h correlation both in the atomic exciton limit and in the weak-coupling plasma limit. Of course, our model breaks down for temperatures below the mean-field "critical temperature" at which $\rho(\omega)$ diverges.

Absorption and gain spectra for the same set of parameters are shown in Figs. 3 and 4. Due to the many-body effects, these spectra deviate considerably from the corresponding free-particle results. Most importantly, they are red shifted due to the renormalization of the single-

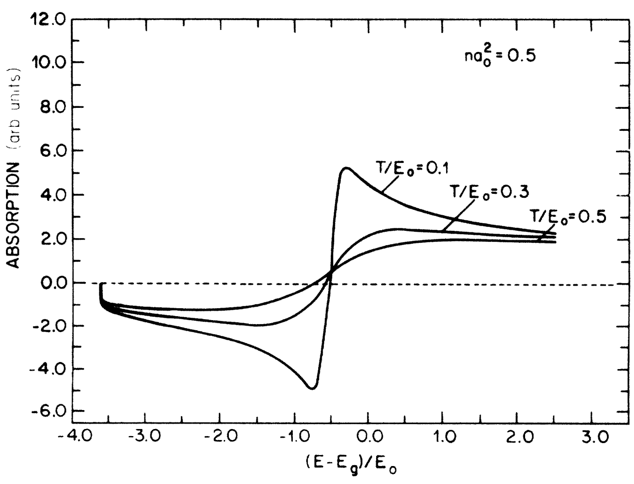

FIG. 4. Optical absorption for $na_0^2=0.5$ and various plasma temperatures T.

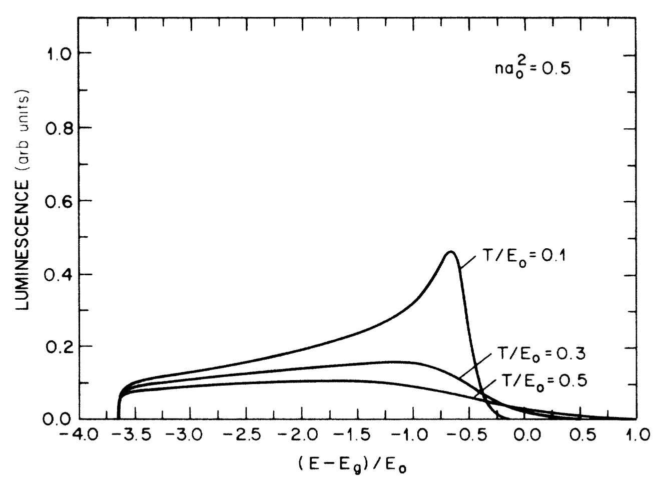

FIG. 6. Luminescence for $na_0^2=0.5$ and various plasma temperatures T.

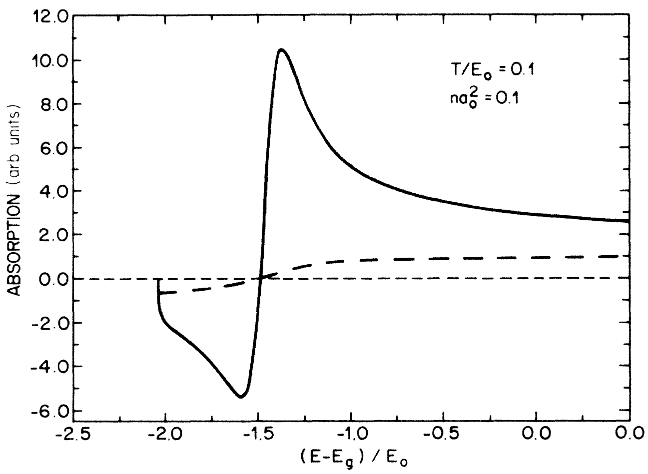

FIG. 7. Comparison of the absorption spectrum (——) with the corresponding one-particle result (— — —) for $na_0^2=0.1$ and $T/E_0=0.1$.

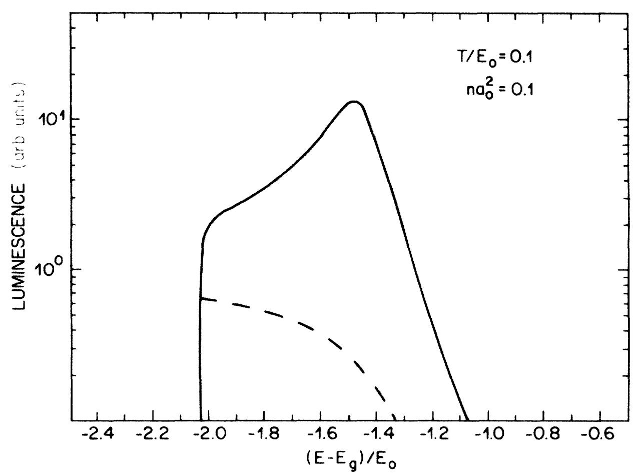

FIG. 8. Comparison of the luminescence spectrum (——) with the corresponding one-particle result (— — —) for $na_0^2=0.1$ and $T/E_0=0.1$.

particle energies and strongly enhanced near the quasichemical potential μ, i.e., near the crossover from gain to absorption, due to the *e-h* correlation. The same holds for the luminescence spectra shown in Figs. 5 and 6. With increasing temperature the effects of the excitonic enhancement decrease rapidly, whereas the renormalization of the single-particle energies is persistent (for the low temperatures studied). With increasing density the *e-h* coupling strength and thus the excitonic enhancement decrease and the renormalized band gap shifts to lower energies whereas the quasichemical potential shifts to higher energies due to the increasing band filling.

The pronounced influence of the *e-h* correlation on the optical spectra of highly excited quantum wells is demonstrated in Figs. 7 and 8, in which we compare absorption and luminescence spectra with the corresponding one-particle results (which include the renormalization of the single-particle energies). It is obvious that the *e-h* correlation can by no means be considered as a minor correction. Contrary to the one-particle spectra, the many-body spectra do *not* reflect the constant density of states of 2D electrons and holes. Instead, the spectra bear great resemblance to the corresponding bulk spectra, which have, e.g., been studied in Ref. 5.

So far, only a few measurements of the optical spectra of highly excited quantum wells have been reported[3,10,31] so that for the time being a detailed comparison of our theory and experimental results is not possible. Most likely, realistic band structures must be included in order to obtain quantitative agreement. This seems to be feasible, since in the degenerate limit (30) can be solved analytically, as has been already mentioned.

Similar conclusions as for the photoexcited systems can be drawn for quantum-well lasers, in which the electrons and holes are electrically injected.[32] Typically, the plasma densities in these systems are somewhat higher than in the photoexcited systems, so that the *e-h* correlation is not so pronounced. However, even at high densities the excitonic enhancement is still strong enough to cause considerable deviations from one-particle line shapes, as should be clear from our results.

Very recently, it has been suggested that the gain spectra of quantum-well lasers can be described by using a one-particle model without *k* selection.[33] Such a model has the advantage of being extremely simple; it is, however, not justified for direct-gap semiconductors. In order to obtain quantitative numbers for the design of laser diodes, all the effects discussed in this paper—band gap and mass renormalization *and* *e-h* correlation—have to be included.

As already mentioned, effects associated with the shake-up of the Fermi sea may also be of importance. Within our simple model, Auger-like indirect transitions will predominantly give rise to low-energy tails in the optical spectra, because at the Fermi level both electrons and holes are well-defined quasiparticles. However, as soon as the strong valence-band mixing in quantum wells is included, there may be additional effects due to the many-body coupling of hole states of different symmetry, in much the same way as in modulation-doped quantum-well structures.[6,7] In order to clarify this point, more realistic calculations are needed in the future.

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