Influence of interparticle interactions on the effective masses of carriers in a quasi-two-dimensional electron–hole plasma in quantum wells in $In_{0.53}Ga_{0.47}As/InP$

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The photoemission from a homogeneous quasi-two-dimensional electron–hole (e–h) system has been studied in quantum wells (of thicknesses $L_z = 8$, 15, and 19 nm) of undoped $In_{0.53}Ga_{0.47}As/InP$ heterostructures. This investigation was carried out at 2 K in the density range $n_{eh} = (0.4\text{–}5.0)\times 10^{12}\ \text{cm}^{-2}$ using magnetic fields $H = 0\text{–}9$ T. In a magnetic field a range of validity of the plasma approximation for the e–h system became much narrower due to enhancement of the exciton effects in a system with discrete energy levels. Nevertheless, this approximation described satisfactorily the properties of particles in a dense e–h plasma found from the Fermi level ($\varepsilon < \varepsilon_F$). An increase in the e–h plasma density altered not only the reduced cyclotron masses of electrons and holes, $\mu^{-1} = m_e^{-1} + m_h^{-1}$, but also modified strongly the dependence of μ on the quasimomentum. As n_{eh} increased to $4\times 10^{12}\ \text{cm}^{-2}$ the value of μ decreased by 20–25 and 7–10% in the range of the quasimomenta $(1.5\text{–}2.0)\times 10^{6}\ \text{cm}^{-1}$ and $(2.5\text{–}3.0)\times 10^{6}\ \text{cm}^{-1}$, respectively. An analysis was made of the contribution of the renormalization of μ related directly to the electron–electron interactions, and also of an additional contribution associated with a change in the energy spectrum of holes because the interparticle interactions resulted in an additional modification of the splitting of the heavy- and light-hole subbands when their populations were different.

1. INTRODUCTION

Quasi-two-dimensional ($2D$) electron–hole (e–h) system of high density in semiconductor structures are attracting major attention because they exhibit a number of fundamentally new physical properties and open up new opportunities for practical applications of heterostructures. Interparticle interactions in a dense e–h system renormalize both the width of the band gap E_g and the dispersion laws of electrons and holes $\varepsilon^0_{e,h}(k)$ (Refs. 1 and 2). It follows from earlier investigations[3–6] that the degree of renormalization of E_g in an e–h plasma in quantum wells and in bulk semiconductors rises monotonically as the e–h plasma density n_{eh} increases (Refs. 7 and 8). The problem of the influence of the interparticle interactions on the effective masses of carriers has been investigated less thoroughly. An e–h plasma in a bulk semiconductor is usually described using the "rigid" band shift approximation.[8] This approximation gives satisfactory results because the changes in the effective masses of carriers in a $3D$ plasma are small.[1,8] The reason is the short-range nature of the interparticle interaction because of a strong screening of the Coulomb potential in a three-dimensional plasma.[1,8]

The rigid band shift approximation had been used also in several investigations to describe the luminescence spectra of an e–h plasma in quantum wells.[3–6] However, the nature of screening of the Coulomb potential changes significantly in quasi-two-dimensional systems, so that the validity of this approximation is not self-evident. Moreover, theoretical estimates[9] show that the change in the effective masses in a dense quasi-two-dimensional e–h system may be considerable.

In the present paper we shall describe an investigation of many-particle interactions in a quasi-two-dimensional e–h plasma in quantum wells based on a determination of the behavior of the photoemission from this plasma in a magnetic field perpendicular to the quantum-well planes. In this situation the motion of charged particles in quantum wells becomes quantized, leading to a discrete energy spectrum of electrons and holes. In the plasma approximation framework the degree of renormalization of the effective masses of carriers is deduced from the change in the energy gaps between the Landau levels. However, it should be pointed out that quantization of the motion of electrons and holes also enhances the exciton effects, particularly in the case of particles in higher filled Landau levels[10,11] and, consequently, narrows the range of validity of the plasma approximation in the description of the quasi-two-dimensional e–h system.

We investigated $In_{0.53}Ga_{0.47}As/InP$ quantum wells. In structures of this kind the exciton radius is known to be quite large: $a_{ex}(3D)\sim 200$ Å, $a_{ex}(2D)\sim 100$ Å. The condition of quasi-two-dimensionality of the e–h system is thus satisfied for quantum well thicknesses less than 200 Å and the condition of a high density $r_s = (\pi a_{ex}^2 n_{eh})^{-1/2} < 1$, is obeyed in the range $n_{eh} \gtrsim 10^{12}\ \text{cm}^{-2}$. A description of these heterostructures is given in §2.

Reliable information on the parameters of an e–h system can be obtained if we ensure the homogeneity of a photoexcited e–h plasma. The conditions under which a photoexcited e–h system is highly homogeneous will be discussed in §2. The spectra of a homogeneous e–h system differ qualitatively from the luminescence spectra determined without any constraints on the propagation of an e–h plasma in the quantum well planes (§3). They are step-like and reflect the reduced density of states of recombining quasi-two-dimensional electrons and holes and, as expected, exhibit no singularities in the exciton region of the kind dis-

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cussed earlier in Ref. 12.

A quantizing magnetic field gives rise to a clear structure in the luminescence spectra of an e–h plasma, which represents allowed transitions between discrete Landau levels (§3). The results of an analysis of this spectral behavior of the splitting of the emission lines, carried out in a wide range of magnetic fields ($H = 4$–8.65 T) and e–h plasma densities [$n_{eh} = (0.4$–$5) \times 10^{12}\ \mathrm{cm}^{-2}$] in order to determine the range of validity of the plasma approximation in the description of the e–h system in a magnetic field, are given in §2. It is found that even if $n_{eh} \sim 10^{12}\ \mathrm{cm}^{-2}$, we have to allow for the exciton effects when describing the energies of the transitions between the upper partly filled Landau levels in the valence and conduction bands. On the other hand, the behavior of the lower Landau levels is in qualitative agreement with the predictions based on the plasma approximation.

An analysis of the dependence of the energies of the allowed Landau transitions on the magnetic field in the range $H = 4$–8.65 T is made in §5 using the plasma approximation. This analysis gives the dispersion relationships $E(k) = E_e(k) + E_h(k)$ for electrons (e) and holes (h) in the lowest quasi-two-dimensional subbands ($n_z = 1$) in quantum wells with $L_z = 80$, 150, and 190 Å in a wide range of the e–h plasma densities [$n_{eh} = (0.9$–$5) \times 10^{12}\ \mathrm{cm}^{-2}$], as well as the reduced cyclotron mass of carriers $\mu^{-1} = m_e^{-1} + m_h^{-1}$ (m_e and m_h are the electron and hole masses, respectively). In a neutral e–h plasma in $In_{0.53}Ga_{0.47}As/InP$ quantum wells the degree of renormalization of μ depends on the carrier quasimomentum and on the e–h plasma density. The dependence $\Delta\mu(k, n_{eh})$ can be explained allowing for the direct influence of interparticle interactions on the cyclotron masses of carriers and for an additional indirect influence of these interactions on the energy spectrum of holes because of renormalization of the splitting of the heavy- and light-hole subbands.

2. EXPERIMENTAL METHOD

The properties of a dense quasi-two-dimensional e–h plasma were investigated in undoped $In_{0.53}Ga_{0.47}As/InP$ heterostructures. Crystals of $In_{0.53}Ga_{0.47}As$ and InP had the same lattice parameters. Therefore, the structures were unstressed and the subband splitting in the quantum wells was governed solely by the size quantization effects.

Our $In_{0.53}Ga_{0.47}As/InP$ heterostructures with single InGaAs quantum wells were grown by the MOCVD (metal organic chemical vapor deposition) method.[7] The quantum well thicknesses were 8, 15, and 19 nm. The width of the exciton emission line recorded at low excitation densities when $T = 2$ K was within 4–5 meV. The spectral position of the exciton line at 2 K was shifted relative to its position in the photoexcitation spectra by 5–7 meV toward lower energies (because of the localization of excitons at quantum well inhomogeneities).

The e–h plasma was excited in quantum wells using a cw argon laser with an output power up to 2 W ($\lambda = 5145$ Å). The sample was exposed directly to a superfluid helium in a cryostat with a superconducting solenoid. The magnetic field could be varied over the range 0–8.65 T. All the measurements were carried out under conditions in which a magnetic field was perpendicular to the quantum well planes. The spectroscopic instrument was an MDR-23 monochromator with a 600 lines/mm grating (dispersion 26 Å/mm). The emission (recombination radiation) was detected with a cooled germanium detector under lock-in conditions.

Special attention was given to ensuring homogeneity of the photoexcitation of the e–h plasma in the quantum wells. First, we avoided the use of heterostructures with several quantum wells, because a distribution of photocarriers varying with the depth in the sample would have given rise to different e–h plasma densities in different quantum wells. Moreover, in the absence of confinement of photocarriers in the quantum well plane, the density in the e–h system varied strongly because of the high rate of diffusion of high-energy Fermi electrons and holes from the photoexcitation region. The most effective method for confining the propagation of carriers to the quantum well planes was a restriction of the size of these wells to that of the exciting laser beam. Our $In_xGa_{1-x}As$ crystals were characterized by an extremely low surface recombination velocity,[13] so that there were no significant restrictions on the linear dimensions of the $In_xGa_{1-x}As$ quantum wells. On the other hand, a reduction in the linear dimensions of the quantum wells to 30–40 μm made it possible to excite a fairly dense ($n_{eh} \approx 5 \times 10^{12}\ \mathrm{cm}^{-2}$, $r_s \approx 0.25$) e–h plasma using a cw argon laser with an output power below 1 W. This enabled us to change over from pulsed to quasi-cw (with pulse length $\tau > 1$–10 ms) measurements which minimized the changes in the e–h plasma density with time.

The linear dimensions of the quantum wells were limited by forming mesas of $30 \times 30\ \mu$m dimensions (Fig. 1). These mesas were produced by optical lithography and dry etching, which ensured a low surface recombination velocity.[13] These measures made it possible to excite an e–h plasma with a high spatial and temporal homogeneity in the quantum wells.

3. EMISSION FROM A TWO-DIMENSIONAL e–h PLASMA

Figure 2 shows the emission spectra of a homogeneous e–h plasma in $In_{0.53}Ga_{0.47}As/InP$ quantum wells with $L_z = 15$ nm at 2 K, obtained using various excitation densities. These spectra differ qualitatively from the spectra of the e–h system without restricting the carrier propagation in the quantum well planes. It is clear from Fig. 2 that as the excitation density W increases above 30 W/cm^2 the intensity of the

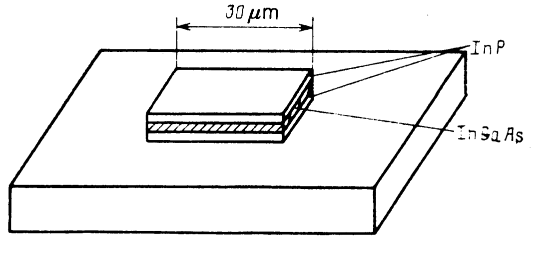

FIG. 1. Schematic diagram of a sample with a mesa used to ensure homogeneity of the e–h system in a quantum well.

exciton emission reached saturation, which was followed by a monotonic broadening of the emission line. The observed saturation of the maximum intensity was a direct consequence of the Pauli principle as applied to electrons and holes. It was evidence of practically complete filling of the states near the edge of the ground ($n_z = 1$) subband in both the conduction and valence bands. When the propagation of carriers in the quantum well planes was unrestricted, the spreading of the carriers out of the excitation zone had the effect that the saturation of the intensity in the emission spectrum of the *e–h* system was not reached even when $W > 10^6$ W/cm^2. Moreover, it was then found that, right up to the highest excitation densities, the emission peak retained its spectral position which was equal to the exciton energy.[12] This behavior of the emission spectra raised the question, discussed in the literature, of whether both the excitons and the plasma exist simultaneously in a dense quasi-two-dimensional *e–h* system.[12,13] The emission spectra of a homogeneous *e–h* system did, however, demonstrate quite convincingly that the conventional excitons are completely screened in a dense *e–h* plasma in a quantum well.

Figure 2 shows that at excitation densities exceeding 300 W/cm^2. The emission spectrum exhibited a second step demonstrating the onset of filling of the next ($n_z = 2$) subband. Its spectral position was practically identical with the position of a peak in the photoexcitation spectra representing the excitation of the $n_z = 2$ excitons. The energy gap between the first and second steps increased from 40 meV to 120 meV when the quantum well thickness was reduced from $L_z = 190$ Å to $L_z = 80$ Å.

At high excitation densities the emission line of the *e–h* plasma shifted toward lower energies. This shift was due to an increase in the exchange–correlation energy. In a three-dimensional *e–h* plasma the magnitude of the shift of the different energy bands is known to be governed primarily by the total densities of electrons and holes, and is almost independent of the distribution of the carriers between these energy bands.[8] Figure 2 shows clearly that in the case of the *e–h* plasma in the quantum wells the shift of the luminescence step for the weakly filled $n_z = 2$ subbands was considerably less than for the strongly filled $n_z = 1$ subbands. The different degrees of renormalization of the width of the band gap for the $n_z = 1$ and $n_z = 2$ subbands indicated that the self-energy component of the *e–h* system was more sensitive to the nature of the wave functions of the particles than would have been the case in a three-dimensional plasma. We made a detailed analysis of the degree of renormalization of the width of the band gap between the various subbands earlier.[7,14] In the present study we were interested in the renormalization of the effective masses.[2)] It was not possible to find the effective masses from an analysis of the luminescence line of the *e–h* plasma in zero magnetic field. The difficulty was that the line profile was described by an integral which contained not only m_e and m_h, but also a number of other parameters dependent on the carrier energy [including the *e–h* recombination probability, damping of single-particle electron and hole states (Γ_e and Γ_h), etc.[1,7,8]]. We therefore returned to measurements in a magnetic field.

Figure 3 shows the emission spectra of our *e–h* plasma in a quantum well with $L_z = 15$ nm when the sample was subjected to a field $H = 6.84$ T perpendicular to the quantum well plane. The emission spectra of the neutral *e–h* plasma with the same occupancy of the Landau levels in the conduction and valence bands were dominated by the allowed ($j_e = j_h$) transitions between the electron (j_e) and hole (j_h) Landau levels. The emission spectra of the *e–h* plasma in a magnetic field had a clear structure. Conse-

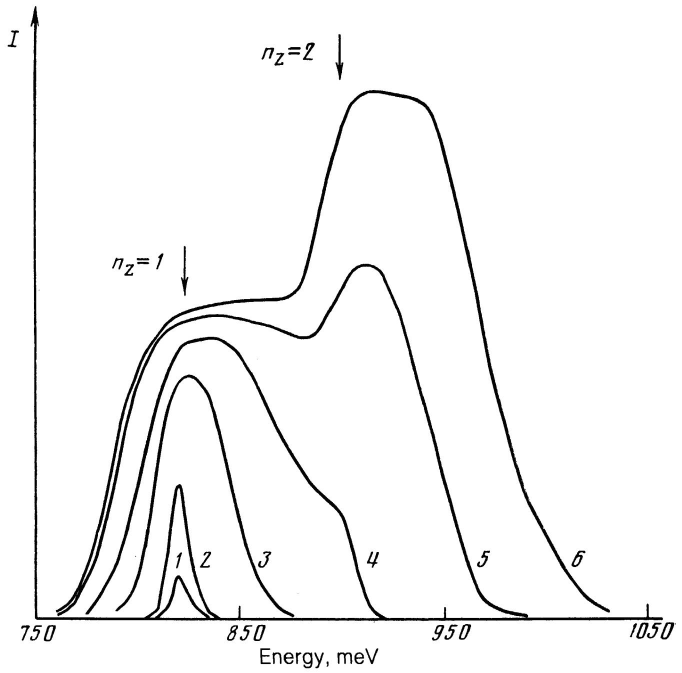

FIG. 2. Emission spectra of a homogeneous *e–h* system in $In_{0.53}Ga_{0.47}As/InP$ quantum wells with $L_z = 15$ nm at 2 K, obtained at different excitation densities W (W/cm^2): 1) 1; 2) 7; 3) 53; 4) 300; 5) 1.7×10^3; 6) 5×10^3.

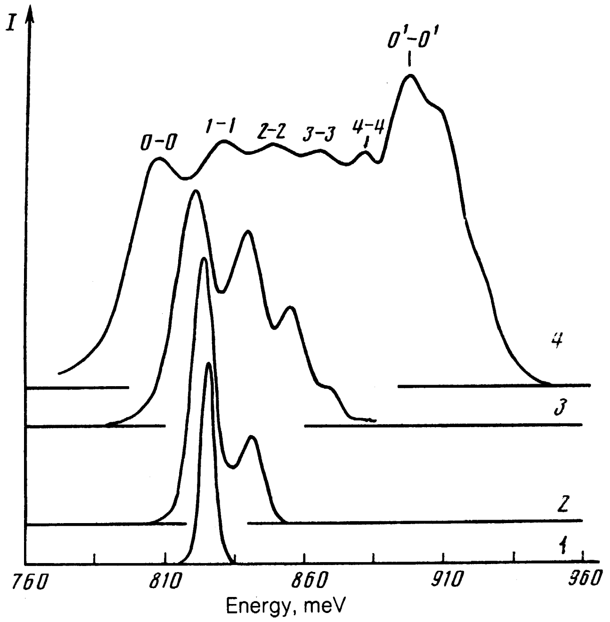

FIG. 3. Emission spectra of an *e–h* system in $In_{0.53}Ga_{0.47}As/InP$ quantum wells with $L_z = 15$ nm at 2 K in a field $H = 6.84$ T, applied at right-angles to the quantum well plane, recorded at different excitation densities W (W/cm^2): 1) 20; 2) 50; 3) 200; 4) 2000.

quently, even at high plasma densities the damping of single-particle states was relatively weak in the case of the filled Landau levels. The emission spectra of the e–h plasma were described by the following three most important parameters: the cyclotron masses, the damping of the single-particle states, and the radiative recombination probability. All these parameters were found practically independently: the cyclotron masses were deduced from the energy gaps between the Landau levels, the damping of the single-particle states was obtained from the line half-widths, and the radiative recombination probability was deduced from the intensities of the individual emission lines.

It is clear from Fig. 3 that an increase in the e–h plasma density resulted in the following main changes in the photoluminescence spectrum. First, new lines appeared and they represented allowed transitions between the higher Landau levels. Their appearance was due to successive filling of these levels on increase in n_{eh}. Second, an increase in n_{eh} caused broadening of all the emission lines because of an increase in the damping of single-particle states. Third, an increase in n_{eh} monotonically shifted all the lines toward lower energies because of an increase in the exchange–correlation energy in the dense e–h system. Finally, as demonstrated clearly in Fig. 3, the energy gaps $\Delta_{i,i-1}$ between the individual Landau levels changed considerably as the density n_{eh} increased. The changes in $\Delta_{i,i-1}$ were mainly due to renormalization of the cyclotron masses of electrons and holes because of interparticle interactions in the e–h plasma. An analysis of these changes was in fact the main purpose of the present investigation.

Using the approximation of renormalized noninteracting quasiparticles in an e–h plasma, we can obtain the dispersion laws of these quasiparticles in the same way as for free carriers in an empty energy band, i.e., on the assumption that the energy gap between the allowed Landau transitions is equal to the sum of the cyclotron frequencies of electrons and holes. However, the range of validity of the plasma approximation in the description of the e–h plasma in a magnetic field is limited because the magnetic field changes qualitatively the energy spectrum of carriers (from continuous to discrete) and thus enhances greatly the exciton effects.[10,11] We shall therefore begin by considering the conditions under which the plasma approximation can be used.

4. RESTRICTIONS ON THE USE OF THE PLASMA APPROXIMATION

In many-body theory the changes in the energies of quasiparticles associated with the interparticle interactions in an e–h plasma are described by the self-energy component $\Sigma_{e,h}[k,\varepsilon(k)]$ (Refs. 1 and 2). The quantity $\mathrm{Re}\,\Sigma_{e,h}$ describes the renormalization of the dispersion laws $\varepsilon^0_{e,h}(k)$ of noninteracting electrons and holes:

$$\varepsilon_{e,h}(k)=\varepsilon^0_{e,h}(k)+\mathrm{Re}\,\Sigma_{e,h}(k,\varepsilon), \qquad (1)$$

where $\mathrm{Im}\,\Sigma_{e,h}$ represents the damping of these states.

In the plasma approximation the dependence $\varepsilon_{e,h}(k)$ can be reconstructed from the energy positions of the Landau levels on the assumption that the Landau level with a number $j_{e(h)}$ corresponds to a quasimomentum

$$\langle k^2_{e,h}\rangle=\frac{eH}{\hbar c}(2j_{e,h}+1). \qquad (2)$$

One of the most direct warnings that the limit of validity of the plasma approximation has been crossed is the discrepancy between the parameters of an e–h plasma found using the approximations represented by Eqs. (1) and (2) in different magnetic fields. Figure 4 shows how the energies of the allowed transitions between the Landau levels in the conduction (j_e) and valence (j_h) bands depend on the e–h plasma density in a quantum well characterized by $L_z = 15$ nm, measured in fields of $H = 4.86$ and 8.65 T. In all cases the e–h plasma density was determined directly from the photolu-

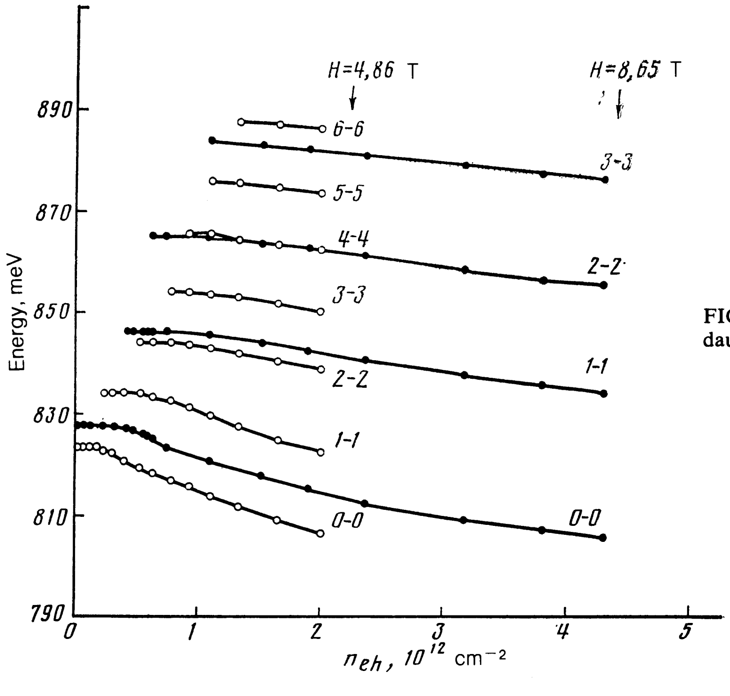

FIG. 4. Dependence of the resolved transitions between the various Landau levels on the e–h system density at 2 K in fields $H = 4.86$ and 8.65 T.

minescence spectra using the Landau level occupancy. We used the circumstance that the number of states in each Landau level was governed only by the magnetic field, in accordance with the expression $n = 2.418\times 10^{10}$ Hv [cm^{-2}], where v is the multiplicity of the spin degeneracy of a level and the magnetic field H is measured in teslas. In our case the spin splitting involving electrons ($s_z = \pm 1/2$) and holes ($j_z = \pm 3/2$) in a field $H < 10$ T was less than the half-width of the luminescence lines, so that we assumed that $v_e = v_h = 2$.

It is clear from Fig. 4 that at high plasma densities the behavior of the energies of the transitions between the filled Landau levels is in qualitative agreement with that expected for the e–h plasma for both $H = 4.86$ T and $H = 8.65$ T: the energies of the transitions decreased monotonically as the plasma density n_{eh} increased because of an increase in Re $\Sigma_{e,h}$ in such a way that the shifts of the levels with the same energy (for example, the $j = 4$ levels when the field is $H = 4.86$ T and the $j = 2$ levels when the field is $H = 8.65$ T and the density is $n_{eh} > 1.2\times 10^{12}$ cm^{-2}) are practically identical. The behavior of the energies of the transitions between the upper partly filled Landau levels differs qualitatively from that just described. It is clear from Fig. 4 that the energies of the transitions between these levels remain unchanged until they are filled completely. This behavior cannot be explained using the plasma approximation framework. It demonstrates that the exciton effects are very important in the case of the upper partly filled Landau levels.

This enhancement of the exciton effects in an e–h plasma in a magnetic field is due to two factors: a reduction in the Bohr radius on the one hand and weakening of the screening of the exciton states, because of the discrete nature of the energy spectrum, on the other. The reduction in the Bohr radius imposes limitations on the range of validity of the

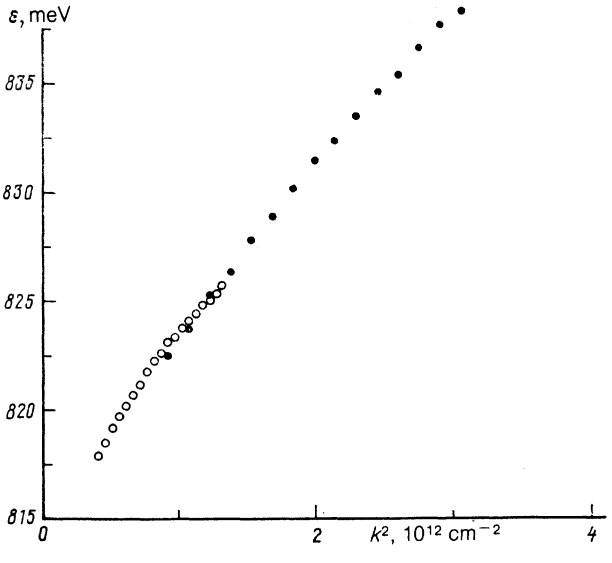

FIG. 5. Dependence of $\varepsilon(k^2) = \varepsilon_e(k^2) + \varepsilon_h(k^2)$ for quasiparticles in an e–h plasma with quantum wells characterized by $L_z = 15$ nm and $n_{eh} = 3\times 10^{11}$ cm^{-2}, reconstructed from the energies of various Landau transitions in fields $H = 3.5$–8.65 T using Eq. (2). The black dots and the open circles represent the values found respectively from the energies of the 0_e–0_h and 1_e–1_h transitions.

plasma approximation on the side of low values of n_{eh}. The influence of the second factor is more complicated.[10,11] In view of the discrete nature of the spectrum of quasi-two-dimensional particles in a magnetic field the screening of the electron–hole correlations in the e–h system when the occupancy of the Landau levels is close to an integer is extremely weak.[10,11] Therefore, the exciton effects do not disappear even at high e–h plasma densities.

The influence of the exciton effects at high values of n_{eh} is illustrated well by Fig. 5. This figure gives the dependence $\varepsilon(k^2) = \varepsilon_e(k^2) + \varepsilon_h(k^2)$ for quasiparticles when $n_{eh} = 3\times 10^{11}$ cm^{-2} in a quantum well with $L_z = 15$ nm, deduced from the energies of various transitions in fields $H = 3.5$–8.65 T using Eq. (2). The transitions between the partly filled levels are excluded from consideration. It is clear from Fig. 5 that there is a small systematic discrepancy ($\sim$1–2 meV) between the values found from the energies of the transitions between different levels. The discrepancy exceeds the limits of the experimental error. Such a discrepancy cannot be explained within the plasma approximation framework. However, as demonstrated in Fig. 5, the magnitudes of the corrections are relatively small and in the case of the states at the Landau levels separated far from the Fermi level we can ignore these corrections, at least in the first approximation.

5. CYCLOTRON MASSES OF CARRIERS IN A DENSE *e–h* SYSTEM

a. Experimental results

The luminescence spectra of an e–h plasma reflect only the allowed, $j_e = j_h$, transitions between the Landau levels. The energy gaps Δ_{ij} between the i_e–i_h and j_e–j_h luminescence lines can be used to find the reduced cyclotron mass of electrons and holes $\mu^{-1} = m_e^{-1} + m_h^{-1}$. It is not possible to separate the contributions of electrons and holes. An analysis of changes in $\Delta_{i,i-1}$ as the number i increases makes it possible to study the behavior of μ throughout the full range of energies $\varepsilon < \varepsilon_F$. The changes in the quantities $\Delta_{i,i-1}$ with increasing Landau level, observed in experiments at a fixed e–h plasma density, demonstrate the nonparabolicity of the dispersion law.

In the rigid band shift model, usually employed to describe an e–h plasma, it is assumed that the effective masses m_e and m_h and, consequently, the value of μ are all independent of n_{eh}. Since the number of states at a Landau level is also independent of n_{eh}, it follows from this model that the energy gaps $\Delta_{i,i-1}$ should remain unchanged when the e–h plasma density is increased. However, Figs. 3 and 4 demonstrate clearly that this is not true. Consequently, in describing the dispersion of carriers in a quantum well in the e–h system we must go beyond the rigid band shift model.

Figure 6 shows the dependences $\varepsilon(k^2) = \varepsilon_e(k^2) + \varepsilon_h(k^2)$ for the $n_z = 1$ subband in quantum wells with $L_z = 8$, 15, and 19 nm, found using Eq. (2), and the energies of the Landau transitions deduced from the luminescence spectra of the e–h plasma with densities 0.9×10^{12} and 1.8×10^{12} cm^{-2} in magnetic fields 4–8.65 T. In all cases the range of energies is limited from above by the energy gap between the $n_z = 1$ and $n_z = 2$ subbands. This is because in the spectra reported in fields $H < 9$ T the luminescence lines representing the Landau transitions from the $n_z = 1$ sub-

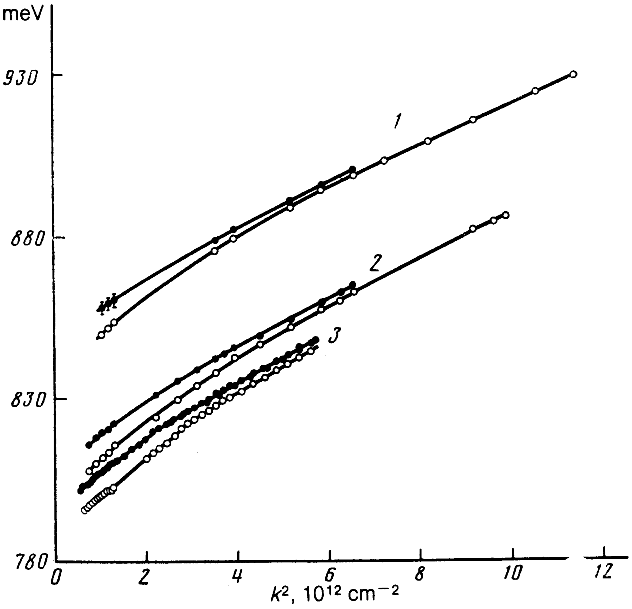

FIG. 6. Dependence of $\varepsilon(k^2) = \varepsilon_e(k^2) + \varepsilon_h(k^2)$ obtained for the $n_z = 1$ subband in quantum wells with $L_z = 8$ nm (1), 15 nm (2), and 19 nm (3), found using Eq. (2) starting from the energies of the Landau transitions in the luminescence spectra of an *e–h* plasma with $n_{eh} = 0.9\times 10^{12}$ cm^{-2} (●) and 1.8×10^{12} cm^{-2} (○).

band were not resolved against the background of the transitions from the $n_z = 2$ subband.

Figure 6 demonstrates that the $\varepsilon(k^2)$ dependence is qualitatively similar for different *e–h* plasma densities. At high values of k $\varepsilon(k^2)$ is nearly linear, which is typical of the parabolic dispersion law. However, at low values of k $\varepsilon(k^2)$ rises strongly, indicating a reduction in the effective mass near the bottom of the energy band. Variation of the quantum well thickness within the range 8–19 nm does not alter qualitatively the behavior of $\varepsilon(k^2)$. However, there are fairly large quantitative differences between the $\varepsilon(k^2)$ dependence obtained for quantum wells with different thicknesses or with different *e–h* plasma densities. We analyze these differences by considering the behavior of the cyclotron masses of the carriers.

The experimental dependence of the reduced cyclotron mass of electrons and holes on k^2 is plotted in Fig. 7 for several *e–h* plasma densities in the case of $In_{0.53}Ga_{0.47}As/InP$ quantum wells with $L_z = 15$ nm. The reduced cyclotron masses were determined from the slopes of the $\varepsilon(k^2)$ dependence. For all values of n_{eh} we found that μ depended weakly on k at high values of the quasimomenta. However, the absolute values of μ decreased by about 8% as n_{eh} increased from 0.9×10^{12} to 3×10^{12} cm^{-2}. Moreover, at all densities n_{eh} there was a strong reduction in μ on approach to the bottom of the energy band. An increase in the *e–h* plasma density shifted the range of the quasimomenta corresponding to the onset of the rapid fall of μ toward higher values of k, but the relative reduction in μ was almost independent of n_{eh} and amounted to $\approx 35\%$.

The changes in the behavior of the $\mu(k^2)$ dependence as the quantum well became narrower are plotted in Fig. 8. The $\mu(k^2)$ dependence remained qualitatively the same for all quantum wells. It is clear from Fig. 8 that a reduction in L_z tended to reduce the effective mass of carriers. This tendency was fully expected for quantum wells with nonparabolic dispersion laws of electrons and holes. The reduction in L_z increased the average quasimomentum k_z of carriers ($\sim L_z^{-1}$), as well as their total kinetic energy (including the size quantization energy $\sim L_z^{-2}$) and, consequently, increased μ (Ref. 15). The limits of the change in the effective mass of electrons as L_z decreased from 19 nm to 8 nm, obtained following Ref. 15, gave $\sim 10\%$. This was in satisfactory agreement with the observed reduction in μ in the range of large quasimomenta $k_{x,y} > 2\times 10^6$ cm^{-1}. Approximately the same change in μ was observed also at low values of k ($< 1.5\times 10^6$ cm^{-1}). In the intermediate range of the quasimomenta, where μ varied rapidly with increasing k, the dependence of μ on L_z was not so clear. In accounting for the

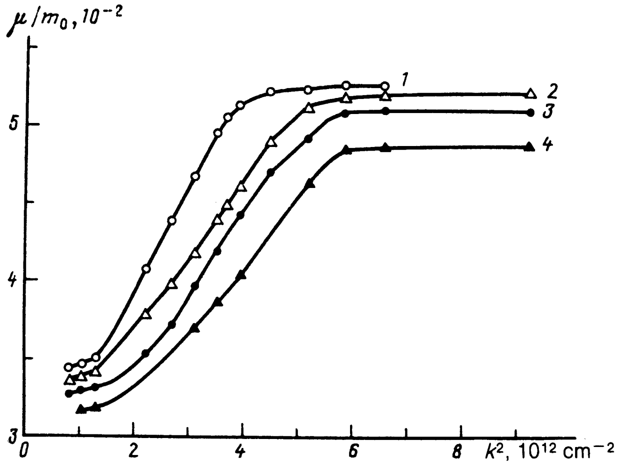

FIG. 7. Experimental dependence of the reduced cyclotron mass of electrons and holes on k^2 plotted for different *e–h* plasma densities in a $In_{0.53}Ga_{0.47}As/InP$ quantum well with $L_z = 15$ nm: $n_{eh} = 0.9\times 10^{12}$ cm^{-2} (1); 1.25×10^{12} cm^{-2} (2); 1.8×10^{12} cm^{-2} (3); and 3×10^{12} cm^{-2} (4).

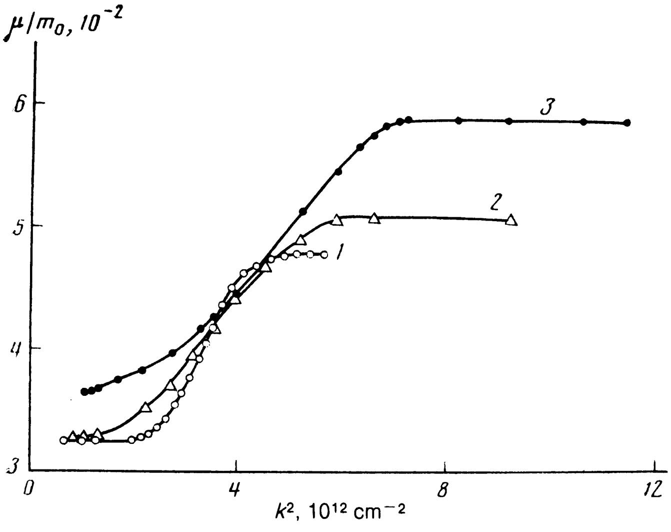

FIG. 8. Changes in $\mu(k^2)$ obtained for a *e–h* plasma with $n_{eh} = 1.8\times 10^{12}$ cm^{-2} as the quantum well thickness is reduced: $L_z = 19$ nm (1), 15 nm (2), and 8 nm (3).

behavior of μ in this range one would have to allow for additional factors which will be discussed in §5b.

Finally, in Fig. 9 we plotted in detail the dependence of μ on the density of the e–h plasma with a quantum well $L_z = 15$ nm thick for several values of the carrier quasimomenta. It is clear from Fig. 9 that the greatest changes in μ occurred in the range $n_{eh} < 2\times 10^{12}$ cm^{-2}. At higher values of n_{eh} the effective mass changed relatively little.

Summarizing the results of these experimental investigations, we noted the following. The $\varepsilon(k)$ dependence was found to be strongly nonparabolic for all the e–h plasma densities and all the quantum well thicknesses in the ranges $n_{eh} = (0.6–4)\times 10^{12}$ cm^{-2} and $L_z = 8–19$ nm. An increase in the e–h plasma density altered the absolute values of the effective masses and modified the dependence of μ on the quasimomentum; the greatest changes in the effective masses occurred near the bottom of the relevant energy band.

b. Discussion

The nonparabolicity of the dependence $\varepsilon(k)$ exhibited by quantum wells in $In_{0.53}Ga_{0.47}As/InP$ is not surprising because both the conduction and valence bands of InGaAs quantum wells are strongly nonparabolic. The effective mass of electrons in $In_{0.53}Ga_{0.47}As$ is given by the expression[16]

$$m_e(\varepsilon)=m_e(0)\,(1+3{,}2\varepsilon/E_g), \quad (3)$$

where $m_e(0) = 0.04m_0$, $E_g = 0.82$ eV, and m_0 is the mass of a free electron. It was found in Ref. 17 that this expression describes well the nonparabolicity of the effective masses in $In_{0.53}Ga_{0.47}As/InP$ quantum wells with $L_z = 10–20$ nm, if ε includes also the size quantization energy of electrons. The nonparabolicity of the hole subband in unstressed

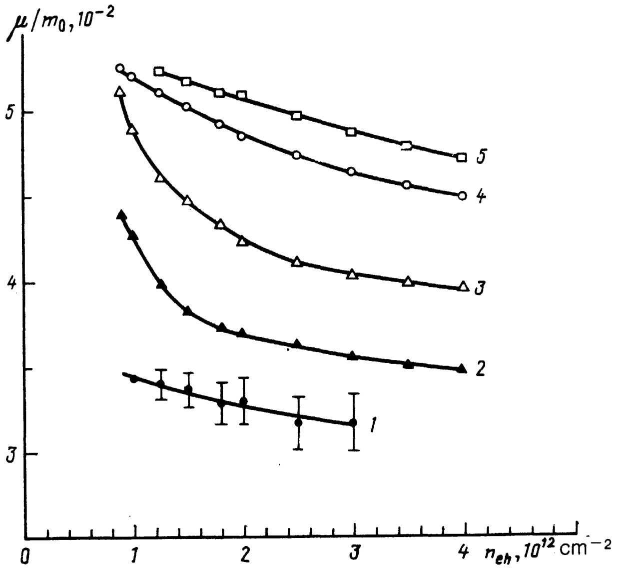

FIG. 9. Dependence of μ on the e–h plasma density in a quantum well with $L_z = 15$ nm obtained for several values of the carrier quasimomenta: $k = 1.0\times 10^6$ cm^{-1} (1); 1.64×10^6 cm^{-1} (2); 2.0×10^6 cm^{-1} (3); 2.28×10^6 cm^{-1} (4); 2.83×10^6 cm^{-1} (5).

$In_{0.53}Ga_{0.47}As$ quantum wells is considerably greater. The fourfold-degenerate valence band is split into two in quantum wells and these are called, in accordance with the ratio of the effective masses along the quantization (z) axis, the heavy-hole subband (which is the main band with $j_z = \pm 3/2$) and the light-hole subband (split off band with $j_z = \pm 1/2$). The splitting ΔE_v between these subbands increases as L_z decreases.

The opposite inequality obeyed by the effective masses of the heavy m_{hh} and light m_{lh} holes in the quantum well plane ($m_{hh(x,y)} \ll m_{lh(x,y)}$) leads to crossing of the heavy- and light-hole terms. However, this interaction between the terms results in their repulsion near the region of the expected crossing and, consequently, it is responsible for the strong nonparabolicity in the energy range $\varepsilon\sim\Delta E_v$. Estimates of $m_{hh(x,y)}$ in the range $\varepsilon \ll \Delta E_v$, obtained using the Luttinger parameters of the valence band of $In_{0.53}Ga_{0.47}As$, give $m_{hh} \sim 0.1m_0$. If $\varepsilon \gg \Delta E_v$ the value of $m_{hh(x,y)}$ lies within the range $(0.5–1)m_0$ (Ref. 17). In view of the large difference between the effective masses of electrons and holes, the contribution of the nonparabolicity of holes to the dependence $\mu(k)$ is greatly reduced. Nevertheless, it follows from the reported values of m_e and $m_{hh(x,y)}$ that for both $\varepsilon \ll \Delta E_v$ and $\varepsilon \gg \Delta E_v$ the influence of the valence band nonparabolicity on the dependence of μ on k is quite strong and can account qualitatively for the observed steep $\mu(k)$ dependence observed in the range $k\approx 1.5\times 10^6$ cm^{-1}: at high values of k where $m_e \ll m_{hh(x,y)}$ the reduced mass is $\mu\approx m_e \approx 0.05m_0$, whereas at low values of k, when $m_{hh(x,y)}$, it falls to $(2–3)m_e$ and the value of μ decreases to $(2/3 - 3/4)m_e \approx 0.03m_0$.

We now consider the changes in μ with increasing e–h plasma density. It is clear from Fig. 9 that there are three ranges of quasimomenta characterized by different behavior of $\mu(n_{eh})$. At high values of k the contribution of holes is small and we have $\mu\approx m_e$. In this range the variation of μ as a function of n_{eh} is gradual and the change amounts to 8–10%. A similar dependence of μ on n_{eh} is observed also at very low values of k, when the contribution of holes is important, but the nonparabolicity of the heavy-hole subband is too weak because of the influence of the split-off light-hole subband.

The change in μ at low and high values of the quasimomenta is mainly due to the fact that the interparticle interactions occurring in the quasi-two-dimensional e–h plasma lead to a significant dependence of the self-energy component Re Σ on the quasimomentum. The change in μ is of the order of 10% and it agrees with the results of a theoretical estimate of the change in the effective mass averaged over all the energies $\varepsilon < \varepsilon_F$ in Ref. 9. In the intermediate range of quasimomenta, characterized by $k\approx 2\times 10^6$ cm^{-1}, the relative change in μ reaches 25–30% and the greatest effect on μ occurs for $n_{eh} < 1.6\times 10^{12}$ cm^{-2}. Such major changes in the cyclotron mass near the bottom of the energy band may be due to the influence of the electron–phonon and electron–plasmon interactions.[10] Calculations reported in Ref. 10 for an e–h plasma in GaAs quantum wells in the absence of a magnetic field showed that the influence of these interactions on the intensity of states is quite strong near the bottom of the band when $n_{eh} \sim 10^{12}$ cm^{-2}. However, we wish to draw attention also to one other possibility of a change in μ near the bottom of the band associated with renormalization, in an e–h plasma, of the splitting of the light- and heavy-hole subbands.

A comparison of Figs. 7 and 9 shows that the strong change in μ as a function of n_{eh} occurs in the same range of quasimomenta as a strong increase in μ with k, associated with an increase in the hole mass due to the interaction between the light- and heavy-hole subbands. We can therefore assume that in this range of quasimomenta an additional mechanism influences the change as n_{eh} increases, which affects only m_h. This may be the difference between the shifts of the edges of the heavy- and light-hole subbands in the e–h plasma because of different populations of these subbands. For example, investigations of the renormalization of E_g in quantum wells demonstrate[6,7] that, in contrast to a bulk plasma, the influence of renormalization on each subband in a 2D system is governed mainly by the population of this subband and not by the total e–h plasma density. We can therefore expect that an increase in the concentration of the e–h pairs in the heavy-hole subband will increase the splitting between it and the light-hole subband and, consequently, will reduce the nonparabolicity of the valence band in the range of small quasimomenta. This mechanism describes qualitatively the observed behavior of μ in the range $k \approx (1.5\text{–}2) \cdot 10^{-6}$ cm^{-1}. However, further calculations are needed in any quantitative comparison.

We can also estimate how much the renormalization of the splitting between the heavy- and light-hole subbands ΔE_v influences the effective mass in a dense e–h plasma by comparing the behavior of the cyclotron mass in the e–h plasma in $In_xGa_{1-x}As$ quantum wells characterized by different initial splittings of the valence band. For example, in the case of stressed $In_xGa_{1-x}As/GaAs$ quantum wells characterized by $\Delta E_v \sim 70$–100 meV the nonparabolicity of the energy spectrum of holes is slight in the range of quasimomenta of interest to us. Consequently, renormalization of the value of ΔE_v for such quantum wells should not alter significantly the mass m_h. It should be mentioned in this connection that earlier we investigated the behavior of μ in an e–h plasma in stressed $In_{0.2}Ga_{0.8}As/GaAs$ quantum wells with selective doping. This doping ensured an equilibrium electron density of $n_e^0 \sim 10^{12}$ cm^{-2} in the quantum wells.[18] It was found that the behavior of μ in a charged e–h plasma in these quantum wells exhibited none of the singularities of the kind observed for a neutral e–h plasma in $In_{0.53}Ga_{0.47}As/InP$ quantum wells. Since the properties of neutral and charged plasmas could be different,[11] it would be desirable to carry out measurements on undoped quantum wells characterized by a weak nonparabolicity of the valence band.

6. CONCLUSIONS

Intensive investigations of the optical properties of a dense quasi-two-dimensional e–h system in quantum wells are stimulated largely by the hope that new many-particle effects can be found in low-dimensional systems. We concentrated our attention mainly on the problem of renormalization of the effective masses in a dense neutral e–h system in $In_{0.53}Ga_{0.47}As/InP$ quantum wells over a wide range of carrier densities. In view of the need to use a large number of fitting parameters to describe the luminescence (or gain) spectra, we were unable to find the values of $m_{e,h}$ from measurements in a zero magnetic field when the spectra had a weak structure. The application of a magnetic field made it possible to obtain information from the experimental spectra. However, the magnetic field altered qualitatively the energy spectrum from quasicontinuous to discrete and thus changed significantly the electron–electron and electron–phonon interactions; and it also enhanced the exciton effects in the e–h system. Consequently, the range of validity of the plasma approximation used to describe the e–h system became narrower and some of the fine effects of the influence of the electron–phonon interaction on the density of the states in the e–h plasma in $H = 0$ could change qualitatively in a magnetic field.[10,11] Nevertheless, our analysis of the luminescence spectra of the investigated e–h system carried out over a wide range of densities ($r_s > 0.2$) and magnetic fields $H < 9$ T showed that the use of the plasma model was possible only as the first approximation of the description of the e–h system in relatively weak magnetic fields and in a limited range of quasiparticle energies in the bottom of the energy band.

An analysis of the energies of the transitions between the Landau levels in our dense e–h system showed that an increase in the plasma density altered not only the cyclotron mass of quasiparticles, but also modified strongly the dependence of μ on the quasimomentum. The renormalization of μ in InGaAs/InP quantum wells was due to two mechanisms: a) modification of the E–k dispersion law of particles because of the electron–electron (electron–hole) and electron–phonon interactions; b) a change in the energy spectrum of holes because the interparticle interactions give rise to additional changes in the splitting of the heavy- and light-hole subbands (caused by their different occupations).

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