values found above for T_1 and T_2 and with the value $|\chi^{(3)}(\Delta\nu\approx 1\ \mathrm{cm}^{-1})|$, we find $\sigma_m \approx 3.7\times10^{-15}$ cm^2 and $N_a \approx 2.8\times10^{17}$ cm^{-3}. Assuming that the aggregation in the solution is complete, we find $n \lesssim N/N_a \approx 6$. The estimates of σ_m and n found here on the basis of the four-wave mixing data are close to the values $\sigma_m \approx 2.3\times10^{-15}$ cm^2 and $n\approx 4$ found in Ref. 4 through a study of exciton annihilation in PIC solutions.

In conclusion we wish to stress that the parameters T_2, σ_m, and n which we have been discussing here pertain to the "optical" aggregates which are responsible for the formation of the exciton band and the high optical nonlinearity of this medium. In contrast with "optical" aggregates, "chemical" aggregates can (according to Ref. 11) consist of tens of thousands of molecules, constituting a system of weakly bound "optical" aggregates capable of a mutual annihilation upon excitation.

[1)] We wish to thank T. V. Veselova for measuring T_1.

[2)] The comparison of the values of $|\chi^{(3)}|$ in solutions of *J* aggregates and in the solutions of the triphenylmethane dyes studied in Ref. 7 was carried out for a common optical density of the samples.

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Excitons and deexcitons in a neutral 2D magnetoplasma with an integer filling of Landau levels: Experiment and theory

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A magnetooptic method has been used to study the excitation of a dense, neutral, low-temperature electron–hole magnetoplasma in InGaAs/InP quantum wells with an integer filling of Landau levels. It is demonstrated that stable excitons and deexcitons exist at the Fermi level.

1. A quasi-two-dimensional (2D) electron gas with a high carrier mobility in a strong magnetic field has some unique properties which stem from the dependence of the electron–electron (*e–e*) correlations on the filling factor ν. In particular, *e–e* correlations lead to a fractional quantum Hall effect and a Wigner crystallization. In a neutral electron–hole (*e–h*) magnetoplasma, the *e–e* correlations are accompanied by *e–h* correlations, which should lead to some new effects.[1–3] So far, there have been only a few experimental studies[4–6] of a quasisteady *e–h* magnetoplasma, because of difficulties in realizing a uniform *e–h* system in a quantum well.

We showed in Ref. 4 that it is possible to resolve this problem to a large extent by using undoped InGaAs/InP quantum wells, in which there are essentially no dark carriers. Because of the low rate of surface recombination in such structures, quantum wells of small dimensions ($\sim 50\ \mu$m) can be used to produce a magnetoplasma which is uniform in space and time and thus to achieve a steady-state photoexcitation. It thus becomes possible to experimentally study correlation effects in a neutral magnetoplasma with a pronounced population inversion and with equal fillings of Landau levels in the conduction band and the valence band. It was shown in Ref. 6 that the *e–e* and *e–h* correlations are strong in the upper filled Landau level even if the average distance between carriers is shorter than the magnetic length, $l = (\hbar c/eH)^{1/2}$, where H is the magnetic field.

A theory was derived in Ref. 7 for elementary excitations of an inverted magnetoplasma near integer values of ν in terms of excitons and deexcitons. Deexcitons are similar to excitons but exist only in an inverted system. They are elementary "deexcitations" in the sense that they have a momentum $\vec{k}$, but their energy is negative ($E_k^d < 0$), and the creation of deexcitons moves the system toward an equilibrium state. That formations of this type might arise had been pointed out previously.[2] At integer values of ν, it is legitimate to introduce excitons and deexcitons under the conditon $\epsilon_C/\hbar w_C \ll 1$, where hw_C is the sum of the cyclotron energies of the electron and the hole, $\epsilon_C = e^2/kl$ is the Coulomb energy, and κ is the dielectric constant (see the discussion below regarding some additional restrictions). The radiative recombination of an e–h pair generates a deexciton with a momentum $\vec{k}=0$ in such a system. The emission thus has a narrow-band spectrum, with a photon energy $\hbar w = |E^d_{k=0}|$.

In this letter we are reporting a theoretical and experimental study of exciton and deexciton elementary excitations in a magnetoplasma with an integer ν. We make use of the emissioin spectra of a neutral magnetoplasma and the photoexcitation of magnetoexcitons with empty bands.

2. We studied InGaAs quantum wells with a thickness $L = 15$ nm. Undoped $InP/In_{0.53}Ga_{0.47}As/InP$ heterostructures with single quantum wells were grown by metalorganic chemical vapor deposition.[8] The nonequilibrium carriers were excited by an argon laser ($\lambda = 5145$ Å). After the emission from the sample was passed through a grating monochromator (600-lines/mm), it was detected by a cooled Ge detector. The samples were immersed in superfluid helium in a cryostat with a superconducting solenoid ($H < 8.7$ T).

Particular attention was paid to the uniformity of the e–h plasma. For this purpose we restricted the propagation of the e–h plasma in the plane of the quantum well by fabricating small mesas ($50 \times 50\ \mu m^2$ in size) and by using a laser beam of slightly larger diameter (100 μm). To excite the plasma, with a density $n_{eh} \leqslant 1.5 \times 10^{12}$ cm^{-2}, we used a 0.1-W cw laser. As a result, n_{eh} remained constant in time.

3. Figure 1 shows emission spectra of a magnetoplasma with Landau-level filling factors close to the integer values $\nu/2 \simeq 0$, 1, 2, and 3. The magnetic field was $H = 8.65$ T, and the temperature $T = 2$ K. Since the spin splitting of the Landau levels was much smaller than the cyclotron frequency, states with different spins are filled essentially simultaneously. With regard to interband excitons (not magnetoplasmons!), the two spin subsystems behave independently. The presence of carriers with opposite spins does not change the positions of the exciton levels. Accordingly, we will be using a filling factor $\nu/2$, and we will say that the filling is "integer" if $\nu/2$ is an integer. The 0–0, 1–1, and 2–2 emission lines correspond to allowed ($j_e = j_h$) transitions between Landau levels in the conduction band (j_e) and the valence band (j_h).

At a low excitation density ($\nu/2 \simeq 0.05$) the emission line corresponds to the emission of magnetoexcitons from the zeroth Landau level. At $\nu/2 \simeq 1$, two lines appear simultaneously in the spectrum: the 0–0 line and a weak 1–1 line. The appearance of carriers at the first Landau level, even before the zeroth level is completely filled, is attributed to the finite energy relaxation time of the photoexcited carriers. It can be seen from Fig. 1 that the smearing of the distribution function is slight and that at $\nu/2 \simeq m$ the $m-1$ level is almost filled, while the m level is almost empty. Under these conditions the radiative recombination of an e–h pair with the slightly filled m level corresponds to the annihilation of an exciton, while the annihilation of a pair with the nearly filled $m-1$ Landau level corresponds to the creation of a deexciton (in level $m-1$) in such a system. It follows from the theory of Ref. 7 that the emission has a narrow band in both cases (Fig. 1). It can also be seen from Fig. 1 that the spectral positions of the m–m line are the same in the cases of the nearly empty ($\nu/2 = m$) and nearly filled ($\nu/2 = m+1$) mth Landau level. This situation corresponds to an equa-

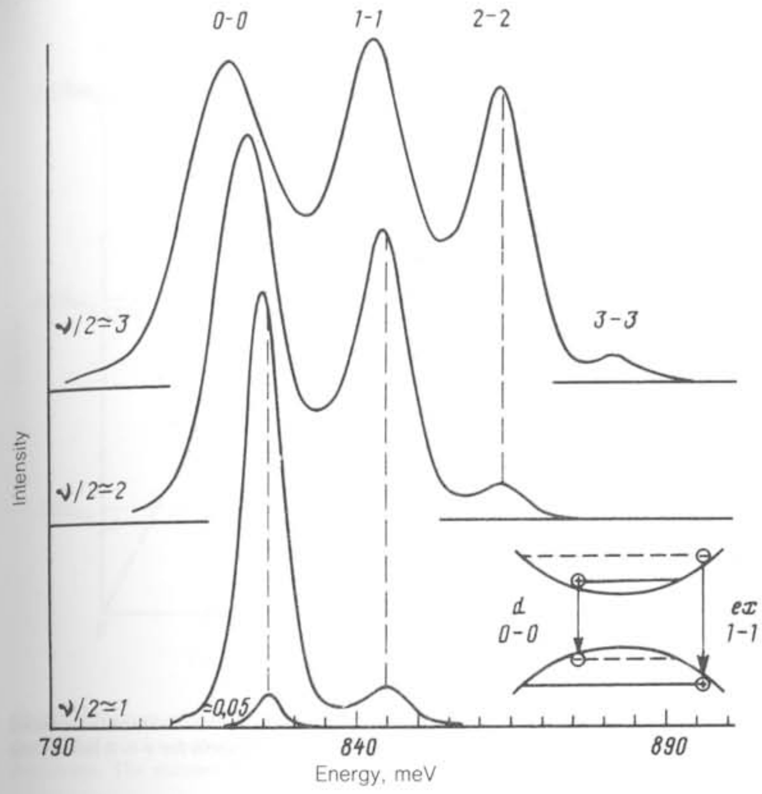

FIG. 1. Emission spectra at filling factors close to integer values. The dashed lines connect the peaks of bands which correspond to stable excitons and deexcitons. The inset shows exciton (ex) and deexciton (d) transitions at $\nu/2 = 1$ (the solid lines represent filled levels, and the dashed lines vacant ones).

lity of the energies of the exciton, $E^{ex}_m(\nu/2 = m)$, and of the deexciton, $|E^d_m(\nu/2 = m+1)|$. As we will see in §4, the theory predicts that these energies will be equal.

Figure 2 shows comprehensive experimental data on the spectral positions of the 0–0, 1–1, and 2–2 lines for $\nu/2 \simeq 0$, 1, 2, and 3. Also shown here are the energies of the 1–1 and 2–2 magnetoexcitons in the vacant band as found from the exciton photoexcitation spectra. It can be seen from Fig. 2 that the energies of the m–m transitions at $\nu/2 = m$ (the beginning of the filling) and $\nu/2 = m+1$ (complete filling of the m level) are the same. Beside these pairs of points, to their left and right the transition energy falls off rapidly with increasing ν. For example, the value of $E^{ex}_1(\nu = 0)$ is 14 ± 3 meV greater than $E^{ex}_1(\nu/2 = 1)$, while $E^{ex}_2(\nu = 0)$ is 22 ± 3 meV greater than $E^{ex}_2(\nu/2 = 2)$.

4. Working from the theory of Ref. 7, we can write explicit expressions for the exciton and the deexcitation which correspond to the direct allowed m–m transition for the case in which there are $N + 1 = \nu/2$ inverted levels $(0,1,...,N)$. For an exciton $(m > N)$ we have

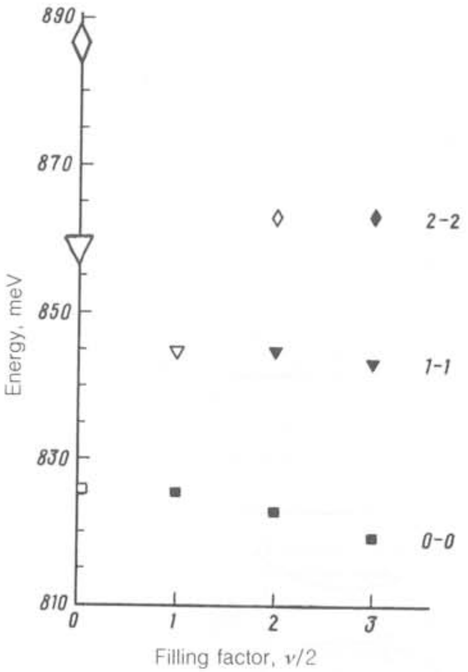

FIG. 2. Positions of the first three exciton transitions (open symbols) and of the deexciton transitions (filled symbols) for integer filling factors (experimental data). The larger size of the symbols for $\nu = 0$ reflects the experimental error.

$$E^{ex}_{m,\vec{k}=0}(N) = E^0_m + \epsilon_m(\vec{k}=0) + E^{HF}_m(N), \tag{1}$$

and for a deexciton ($m<N$) we have

$$E^{d}_{m,\vec{k}=0}(N) = -E^0_m + \epsilon_m(\vec{k}=0) - E^{HF}_m(N). \tag{2}$$

The energy E^0_m includes the gap E_g and the Landau quantization energy in both bands. The term

$$\epsilon_m(\vec{k}) = -\epsilon_C \int d\vec{q} V(\vec{q}) w_{mm}(\vec{q}^2) e^{i\vec{k}\vec{q}}/(2\pi)^2 \tag{3}$$

is the correlation energy of an electron and a hole which are bound in an exciton. The momentum $\vec{q}$ is expressed in units of l^{-1}; $V(\vec{q}) = 2\pi/\vec{q}$ is a Fourier component of the Coulomb potential. In the Hartree–Fock exchange energy

$$E^{HF}_m(N) = -2\epsilon_C \sum_{s=0}^{N} \int d\vec{q} V(\vec{q}) w_{sm}(\vec{q}^2)/(2\pi)^2 \tag{4}$$

each term corresponds to an interaction between an electron in level m and all the electrons in level s (plus the corresponding contribution for holes in the valence

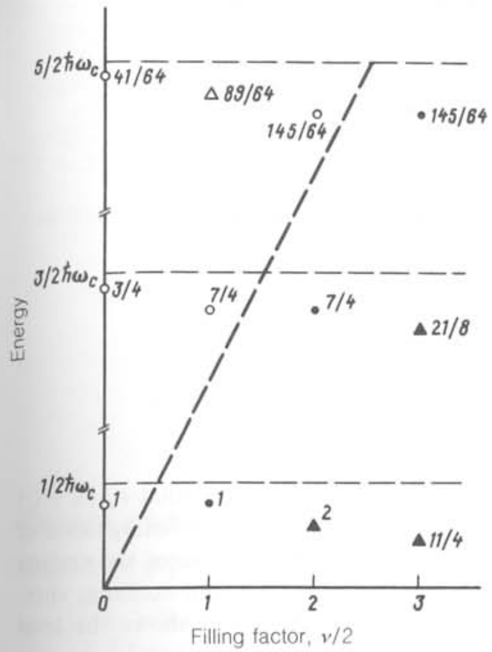

FIG. 3. Positions of the first three transitions (theoretical). Circles—Stable particles; open circles—excitons; filled circles—deexcitons; triangles—unstable particles; open triangles—excitons; filled triangles—deexcitons. The number beside a symbol gives the Coulomb energy in units of $(\pi/2)^{1/2}\epsilon_C$.

band). The functions $w_{sm} > 0$ are expressed in terms of Laguerre polynomials. Since $\epsilon_m(\vec{k}=0)$ appears with the same sign in (1) and (2), while E^{HF}_m appears with opposite signs, the electron–hole symmetry is not totally trivial, because of the renormalization of E_g.

These equations are valid under the condition $\epsilon_C \ll \hbar\omega_c$ in three cases in which the particles are stable: (1) for excitons with arbitrary m in the absence of free carriers [in this case the third term in (1) is not present]; (2) for excitons with $m = N+1$; (3) for deexcitons with $m = N$. In the latter two cases, the transitions occur between the Landau levels which are closest to the Fermi level, i.e., at the forward Burshteĭn edge. These states are shown by the open circles (for excitons) and the filled circles (for deexcitons) in Fig. 3. No claim is made that the theory gives a reliable description of any states other than these. In other cases, the situation is greatly complicated by the strong interaction of the excitons (deexcitons) with the multiparticle continuum, which leads to processes in which these entities decay.

In this theory, the energies of the excitons and deexcitons in Landau level N are the same:

$$E^{ex}_{N,\vec{k}=0}(N-1) = |E^{d}_{N,\vec{k}=0}(N)|. \tag{5}$$

TABLE I. Differences between the energies of exciton and deexciton transitions (in meV).

Landau level	δ_{01}	δ_{12}	δ_{02}	δ_{23}
0	0 (0,5)	-	-	-
1	16 (14)	0 (< 0,3)	-	-
2	-	-	26 (22)	0 (< 0,3)
0	-	16 (2,5)	-	12 (3,5)
1	-	-	-	14 (1,5)

Shown here are theoretical values and (in parentheses) experimental values of $\delta_{ij} = E(\nu/2 = i) - E(\nu/2 = j)$ for the corresponding Landau levels. The upper half of the table corresponds to stable particles, and the lower half to unstable particles.

Equation (5) follows from Eqs. (1)–(4), as can be seen by noting that $\epsilon_N(\vec{k}=0)$ agrees precisely with the $s = N$ term in the sum in (4) (without the preceding factor of 2), cancels out the deexciton energy in it, and (in contrast) introduces the exciton energy in it. Since the cancellation occurs between the correlation and exchange energies, Eq. (5) is a direct reflection of correlation effects. Figure 3 shows the level positions calculated from Eqs. (1)–(4) for $m, N \leqslant 2$.

5. The experimental ν dependence of the positions of the emission lines (Fig. 2) is compared with the theoretical dependence (Fig. 3) in Table I. We see that experiment and theory lead to identical exciton and deexciton energies in each of the Landau levels. Both theory and experiment show that the levels in a magnetoplasma are substantially lower than the magnetoexciton levels in an empty quantum well. The quantitative agreement of experiment and theory can be judged completely satisfactory for the stable particles, since (1) the condition of a strong magnetic field is just barely satisfied at $H \sim 8.5$ T ($\epsilon_C/\hbar\omega_c \simeq 0.5$) and (2) the valence band in InGaAs has a complex structure. For the unstable deexcitons, the agreement of theory and experiment is only qualitative. The theory predicts an overly large lowering of the transition energy with increasing ν. This result is apparently evidence of a strong interaction of unstable deexcitons with a multiparticle continuum. The theory is thus supported convincingly for the stable particles, while for the unstable particles it requires generalization to incorporate virtual and real decays of deexcitons.

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Direct observation of the compression of a transport current by an alternating magnetic field

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A collapse of the transport current in an alternating magnetic field has been observed directly through measurement of the spatial distribution of the magnetic field. The transition of a superconductor to a resistive state under these conditions has been studied.

The critical-state model proposed by Bean about thirty years ago[1] is now widely used to describe the electromagnetic properties of high-T_c superconducting samples. The application of this model to conventional ("cold") superconductors has resulted in a successful explanation of the irreversibility of the magnetization curve of hard superconductors and has made it possible to calculate the loss in a low-frequency magnetic field. However, there has been no detailed analysis of the very simple situation in which the external magnetic field H is directed along the transport current $\vec{I}$, and there is a force-free configuration.

An analysis of the case $\vec{H} \parallel \vec{I}$ (($I < I_c$, where I_c is the critical current) carried out in Ref. 2 showed that a compression of a transport current (a collapse) by an alternating external magnetic field can be observed in this geometry. As a result, the region through which the transport current flows shifts toward the axis of the sample, and the current reaches the center of the sample at a certain field amplitude H^*. Under these conditions the transport current density rises to the critical value, and an electric field arises in the interior of the sample as a result of the dissipative flow of the current I. The current I^* at which the dissipation arises was measured as a function of the field

